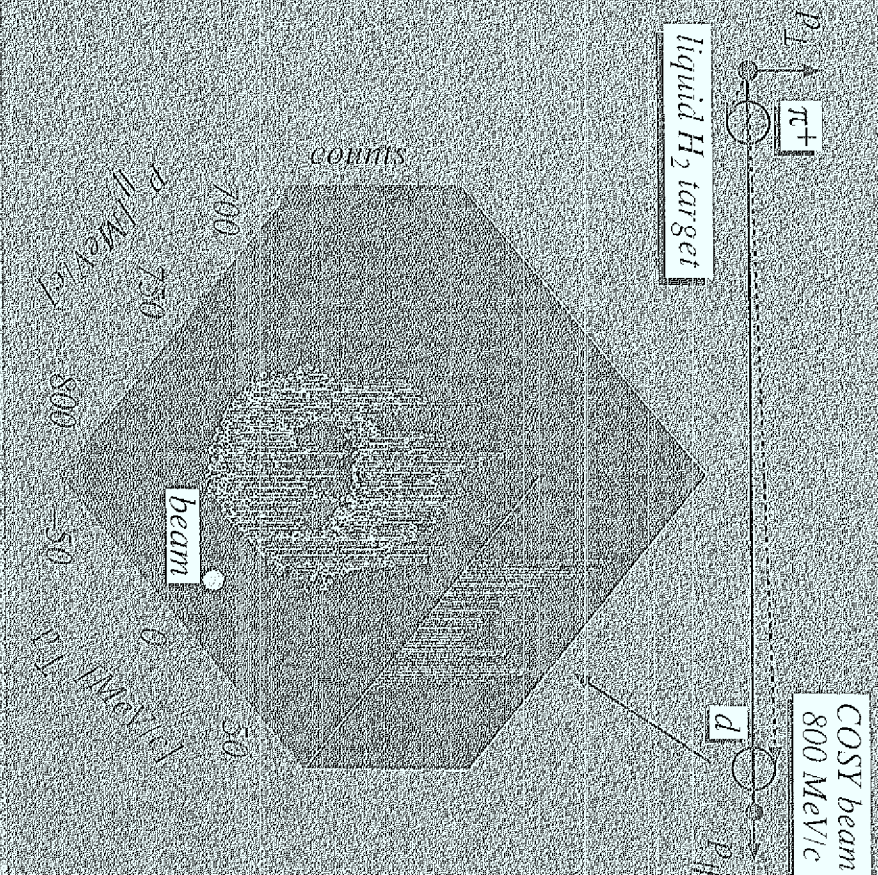
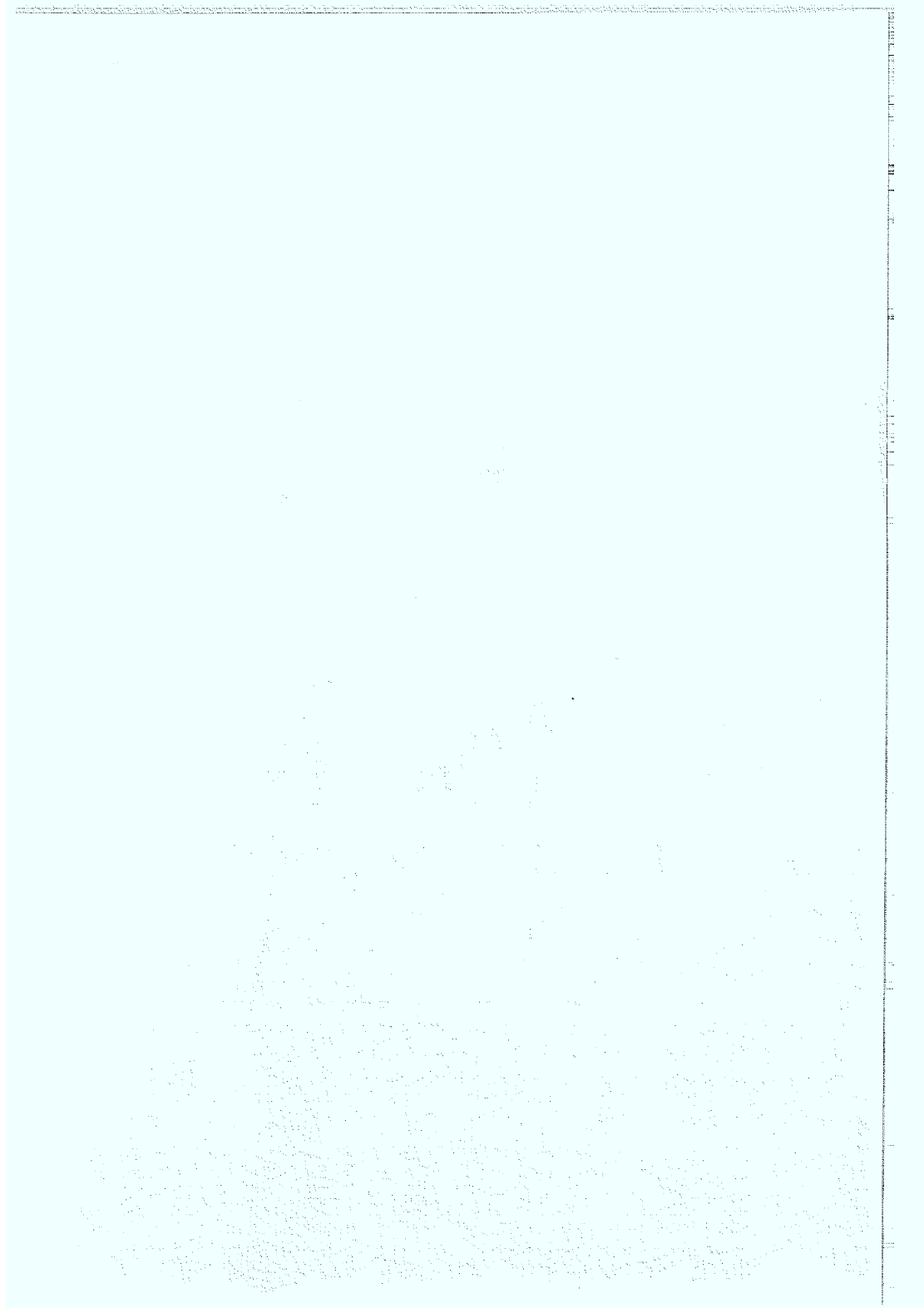
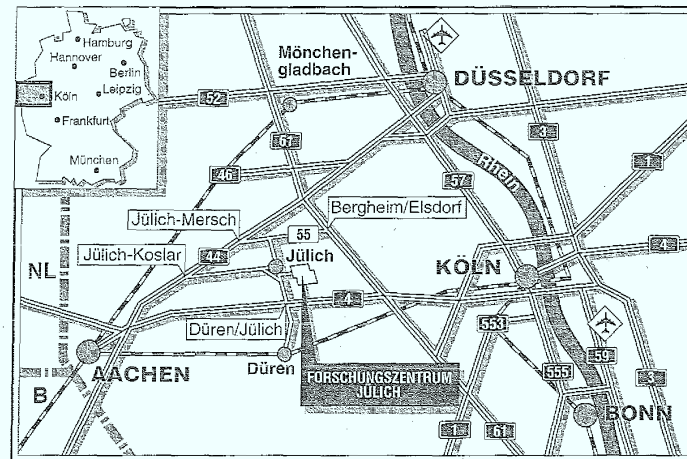


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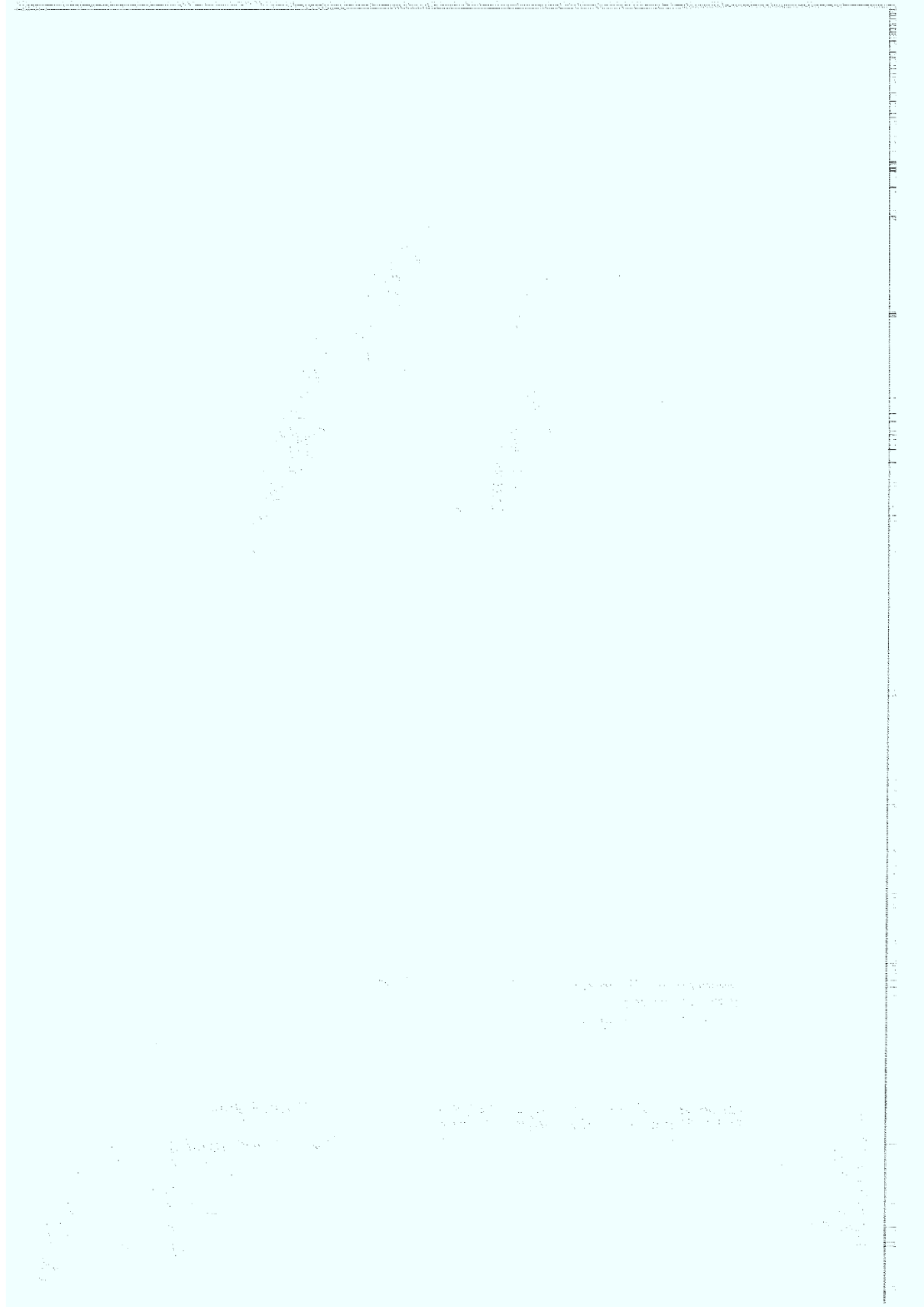
ANNUAL REPORT 1994





Berichte des Forschungszentrums Jülich ; 3035
 ISSN 0944-2952
 Institut für Kernphysik Jüli-3035

Zu beziehen durch: Forschungszentrum Jülich GmbH · Zentralbibliothek
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Annual Report 1994

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Measurement of the reaction $p + p \rightarrow \pi^+ + d$ very close to the threshold.

In the upper part of the figure the momentum vectors of the four reaction particles are shown, where $p_{\parallel}$ is the direction of the beam and $p_{\perp}$ the one perpendicular to that. The target proton at rest and the beam proton at 800 MeV/c are shown as full dots. The kinematical loci of the emerging two reaction products are ellipses as shown. A typical event is indicated by the dotted arrows. The lower part of the figure shows a density distribution of deuterons in such a polar diagram measured with the magnet spectrometer BIG KARL. Already these raw data show the elliptical distribution of the deuteron momentum vectors. The spectrometer has nearly 100 % acceptance for the deuteron momenta. The projection on the parallel momentum axis corresponds to the differential cross section in the centre of mass system. For details see contribution of the GEM collaboration (page 36).

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P r e f a c e

The past year was a very successful one for both our machine and also our physics program.

Sustained efforts in accelerator development at COSY have been rewarded with many highlights. The most outstanding of these was the successful acceleration of the proton beam up to the design momentum of 3300 MeV/c. This enabled the EDDA-collaboration to take precise data in pp-scattering in a continuous energy range from 0.5 to 2.5 GeV. The remarkable quality of the data is also testimony for the high spacial stability of the proton beam achieved during the acceleration process. Beams with various energies were delivered to all external experiment areas. In accordance with the needs of the experiments, the number of extracted particles was raised significantly to 10^8 per spill and the extraction period lengthened beyond two seconds improving the data quality substantially. Another achievement was the reduction in beam radius on the targets of the experiments to about 1 mm. This is of crucial importance for the reconstructions of vertices.

Preparation of the 0^0 -Facility and its detector system, which will be used for the study of the subthreshold K^+ -production in proton-nucleus reactions and the proton-induced deuteron break-up, has advanced considerably in 1994. The special beam tube for the installation of COSY-13 has been implemented in the COSY-ring.

With COSY having achieved its design energy, the internal beam experiment EDDA using a large solid angle scintillator hodoscope has entered its data production phase. A fiber target intercepting the circulating beam during acceleration. A perfect luminosity control and precise kinematic triggering has allowed COSY to measure the excitation functions of the spin averaged cross section of elastic proton-proton scattering over almost the full angular range.

The spectrometer BIG KARL was used by the GEM- and MOMO-Collaborations for their first successful runs to study pion production in different reactions. The MOMO-vertex wall behind the target chamber, which is used to detect the coplanar pions, was completed with 672 independent scintillating fibers. The collaboration was able to identify the first events of the exclusive $pd \rightarrow {}^3\text{He} \pi^+ \pi^-$ reaction to study meson-meson and meson-nucleon interactions.

The GEM-Collaboration studied the reaction $pp \rightarrow d\pi^+$ close to threshold. They were able to clearly display the unique kinematic signature of this reaction in the focal plane of the spectrometer. This success was of special importance because the remarkable results demonstrated the capability of the new ion optical entrance system to resolve both the momentum and also the scattering angles a prerequisite for many other experiments.

The Time-of-Flight spectrometer (TOF) went into operation in May. The most simple version has already allowed the simultaneous measurement and reconstruction of various two- and three-particle channels in proton-proton reactions.

The development of a polarized COSY beam has made big progress: a polarized atomic beam source has been installed at the injector cyclotron to take up operation.

The EUROBALL project has been formally initiated by the signature of a memorandum of understanding by the funding agencies of six European countries. The prototype of the CLUSTER detector has been successfully completed in collaboration with the Cologne group.

The theoretical work is closely connected with the experimental program. The central issue is to connect the physics at COSY with the research at other laboratories. For example, a method was suggested to determine the pion structure function down to very small x which has been implemented at HERA.

A large fraction of our successes was made possible by the accomplishments of our engineers and technicians, the IKP service groups, and the general infrastructure of the KFA and last but not least by our numerous Diplomanden and Doktoranden.

Josef Speth

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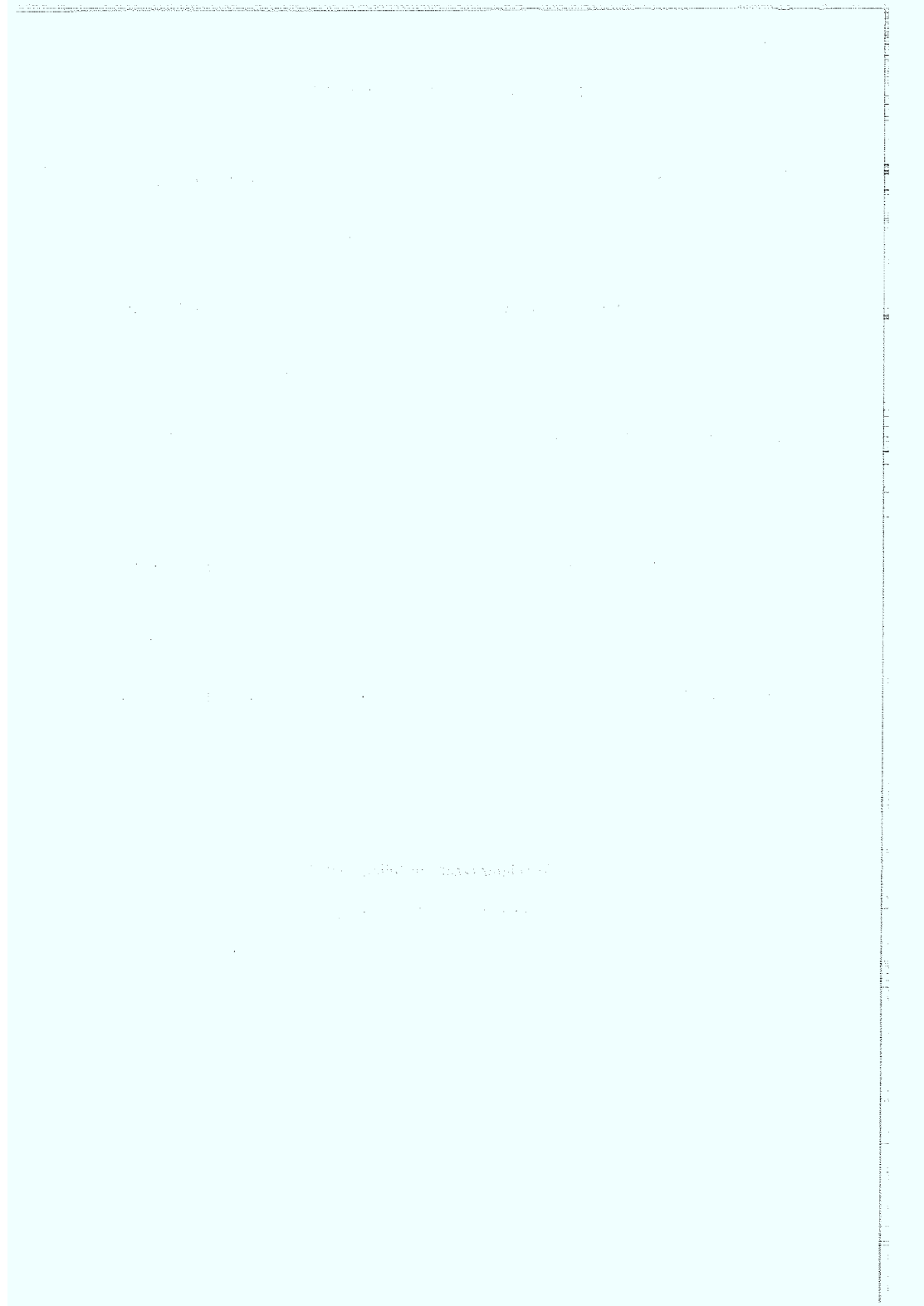
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I. Experimental Nuclear Physics

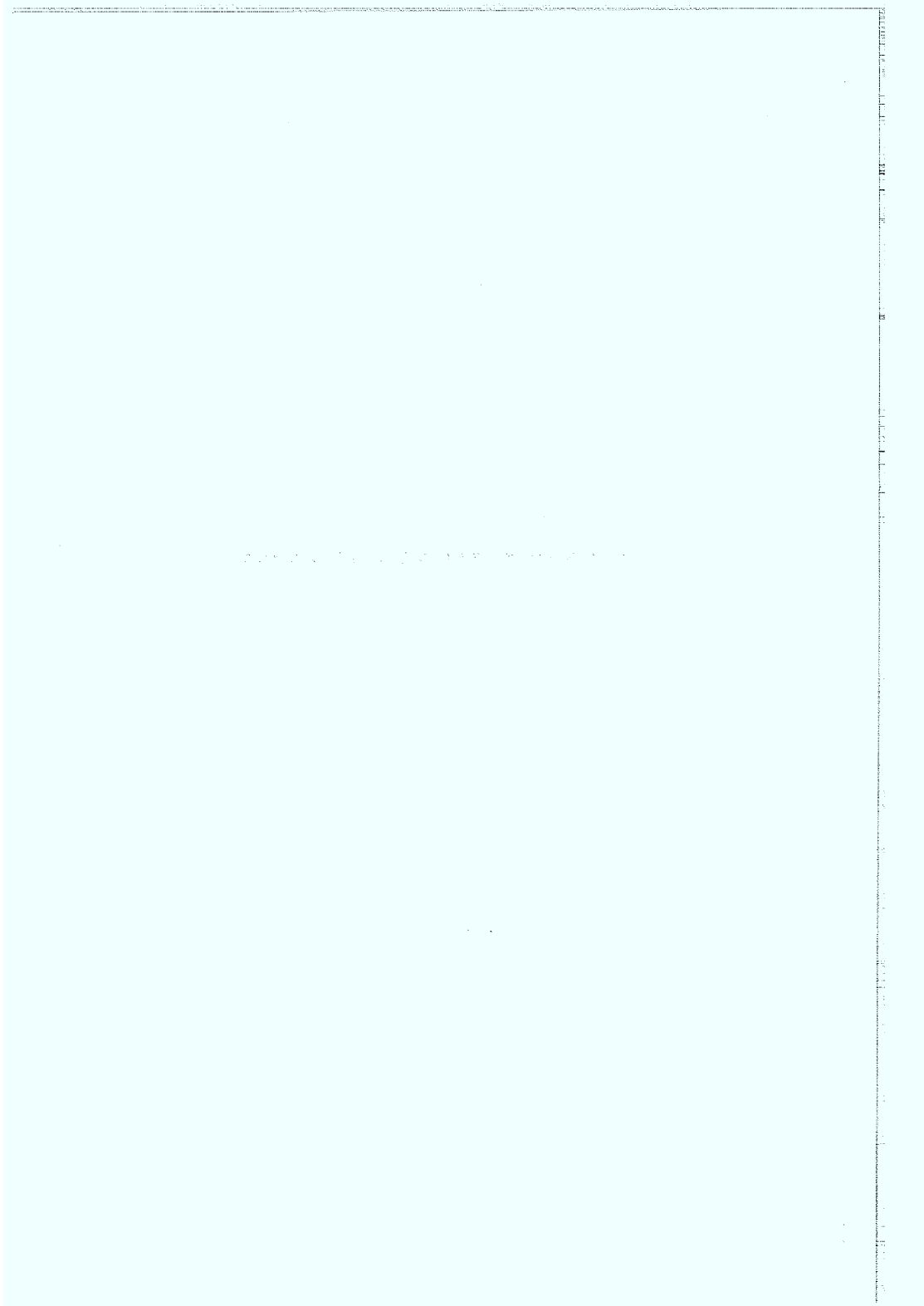
1. MEDIUM ENERGY PHYSICS

Experimental Section

General Procedure

The general procedure for the synthesis of the polymers was as follows. A solution of the monomer (1.0 mmol) and the initiator (0.1 mmol) in the solvent (10 mL) was placed in a round-bottomed flask equipped with a magnetic stirrer. The mixture was stirred at 60 °C for 24 h. The reaction was quenched by the addition of methanol. The resulting polymer was isolated by precipitation into methanol and dried under vacuum at 40 °C for 24 h. The yield of the polymer was determined by gravimetric analysis. The molecular weight of the polymer was determined by gel permeation chromatography (GPC) using polystyrene as a standard. The GPC was performed on a Waters apparatus with a Styragel HR5E column and a refractive index detector. The mobile phase was THF at a flow rate of 1.0 mL/min. The oven temperature was 35 °C. The data were analyzed using the GPC software package. The inherent viscosity of the polymer was determined by measuring the flow time of a solution of the polymer in THF (0.5 wt %) at 30 °C using an Ubbelohde viscometer. The inherent viscosity was calculated using the equation: $[\eta] = (t - t_0) / (c \times t_0)$, where $[\eta]$ is the inherent viscosity, t is the flow time of the solution, t_0 is the flow time of the solvent, and c is the concentration of the solution.

1.1 EXPERIMENTS AT COSY



Identifying $pp \rightarrow d\pi^+$, $pp \rightarrow pn\pi^+$, and $pp \rightarrow pp\pi^0$ with a First Measurement with the Time of Flight Spectrometer at COSY

E. Roderburg
for the COSY-TOF-Collaboration

In May 1994 a first measurement with the time-of-flight detector and a COSY-beam of 1 GeV/c protons was performed. The central detector "Quirl" [1] was mounted at a distance of 70 cm to the liquid hydrogen target (Fig. 1).

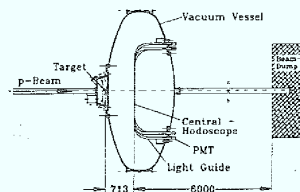


Figure 1: Arrangement of the first measurement with the time-of-flight spectrometer.

This distance provides that both charged particles of the reactions $pp \rightarrow d\pi^+$, $pp \rightarrow pn\pi^+$, and $pp \rightarrow pp\pi^0$ can be detected in the central detector.

The beam was adjusted at the target point by evaluating the spot location and size with a grid of semiconductor strips, which have a pitch of only 0.2 mm. The beam profile is shown in Fig. 2 in a horizontal and vertical cut. This beam profile is small enough to pass through the cylindrical LH_2 target, which has a diameter of 6 mm and a length of 5 mm.

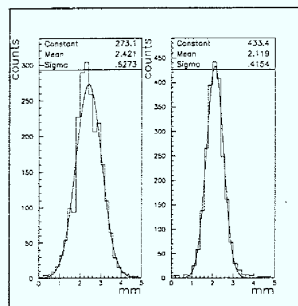


Figure 2: Beamprofile, left: horizontal, right: vertical.

The events are analyzed by the information of three scintillators for each of the two tracks. Each track is recorded by three time measurements with TDC's and three energy loss measurements with ADC's, this is a 6 fold measurement of the velocity of the particle. The direction of the track is evaluated by a coincidence of the three layers of the Quirl as shown for a measured event in Fig. 3.

The "two particle" events $pp \rightarrow d\pi^+$ have a signature which is given by the coplanarity of the beam particle with the scattered particles, the difference of arrival time of the π^+ and the deuteron, and the difference of energy loss of both particles. The difference of time of flight of between pion and deuteron is shown in Fig. 4.

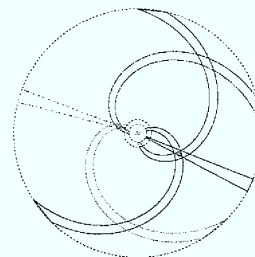


Figure 3: The scintillator elements which have been hit by a two track event are plotted.

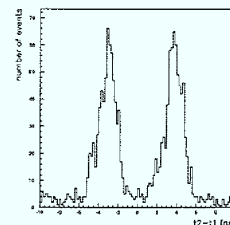


Figure 4: Difference of time of flight for deuteron π^+ events.

Assuming the masses of the deuteron and the pion, the full kinematics of the process can be calculated. The resulting momenta vectors have to point to the ellipses, which reflect the momentum and energy conservation. (Fig. 5). The remaining events, which show no coplanarity, can result from $pp \rightarrow pn\pi^+$, and $pp \rightarrow pp\pi^0$ (compared to these reactions the bremsstrahlung $pp \rightarrow ppy$ is suppressed in this momentum range by a factor of about 1000).

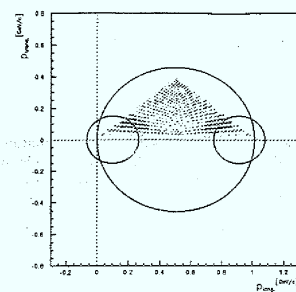


Figure 5: Measured endpoints of momentum vectors, the small ellipses are the calculated position of the deuteron and the pion momentum vector endpoint respectively.

These events are analyzed by assuming the masses of track 1 and track 2 to be a) pp, b) $p\pi^+$, c) π^+p . With the measured velocities and directions of both tracks and the assumption of the masses, the mass of the missing neutral particle is calculated. For each event the "missing masses" are calculated with the assumptions a), b), and c). Fig. 6 shows the resulting squared missing mass distributions. Two vertical lines indicate the area where the neutral particle is assumed to be the π^0 or the neutron. The width of the area is due to the finite detector resolution. The hatched distributions show the remaining events, when the events which can be interpreted in the other distributions as neutral particles are cut out, in case of Fig. 1 the upper cuts in Fig. 2 and 3 were chosen to be $9 \text{ GeV}^2/c^4$ [2].

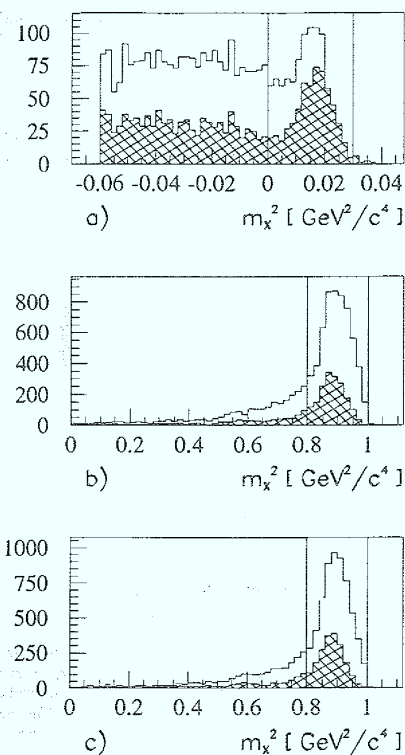


Figure 6: Squared missing mass distributions for the mass assumption a) pp, b) $p\pi^+$, c) π^+p .

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Estimate of the Total Cross Section of the Proton-Proton-Bremsstrahlung near Pion Threshold

E. Kuhlmann¹, H. Müller², B. Naumann², L. Naumann²
for the COSY-TOF collaboration

The COSY-TOF collaboration plans the investigation of the proton-proton-Bremsstrahlung (pp γ) with a sufficiently high precision to study off-shell effects of the nucleon-nucleon scattering near the pion threshold [1]. To understand the influence of the background reactions on the investigated nuclear process it is important to know the order of magnitude of their total cross sections. The aim of this study was to estimate the total cross section of pp γ at proton beam energies of about 300 MeV. By Rothe et al. [2] a total cross section for pp γ of $\sigma_{\text{tot}} = (0.7 \pm 0.15) \mu\text{b}$ was obtained at 204 MeV for a photon energy of $E_\gamma \geq 35$ MeV. At TRIUMF [3] and Bloomington IUCF [4] differential cross sections had been measured at 280 MeV and 294 MeV, respectively. The TRIUMF data were obtained with the help of a coplanar spectrometer, where all three scattered particles had been detected. The IUCF data were measured with a noncoplanar spectrometer geometry. Both outgoing protons were constrained to angles from 4.8° to 12° and kinetic energies between 20 and 120 MeV.

The Bremsstrahlung had been selected by a missing mass procedure. Authors of the IUCF data noted a good agreement between the pure phase space distribution and the shape of the data [4]. We generated phase space distributed events for the IUCF kinematic conditions taking into account a scattering angle straggling of 1.2° , a momentum resolution of 8 % and the quadrupole character of the radiation. For the three angular distributions the calculated differential cross sections are shown in Fig. 1. They were normalized with an averaged total cross section of $\sigma_{\text{tot}} = 9 \mu\text{b}$. The averaged error of this value has been obtained to 3 μb . The result of similar calculations for the TRIUMF data at proton scattering angles about 12° is in a good agreement with the one above. The estimate of the total cross section has been carried out in the framework of the program AFKINE [5] based on the General Simulation program FOWL of the CERN Program Library.

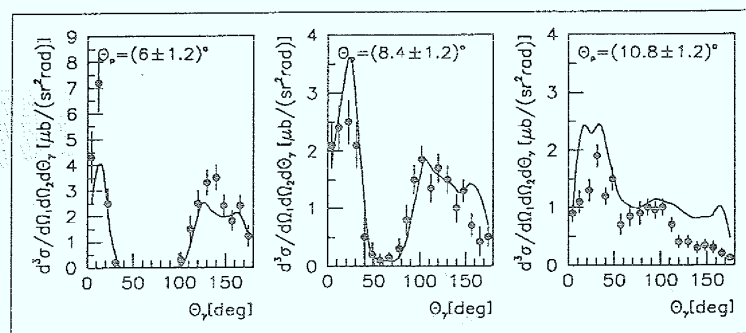


Fig. 1: Differential cross section for pp γ in the laboratory system. The data correspond to [4] and the curves to a phase space distribution, normalized with the averaged total cross section $\sigma_{\text{tot}} = 9 \mu\text{b}$.

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² Forschungszentrum Rossendorf, Institut für Kern- und Hadronenphysik

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On the way to Hyperon Production at the COSY Time-of-Flight Spectrometer

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A proton beam was first transported into the area of the time-of-flight spectrometer at COSY in may 1994. As fig.1 shows the beam with a momentum of 1 GeV/c could be focussed on a spot of about 1 mm diameter, monitored by different arrays of silicon microstrip detectors and scintillating fibres.

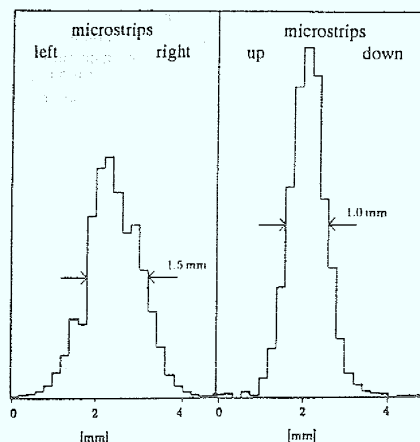


Fig.1: beam profile at the focus of the TOF-beamline, measured with silicon microstrip detectors (resolution 200 μ m)

After the beam preparation the "Erlangen Startdetector" in combination with the "Jülich Quirl" was used for first test measurements towards the associated hyperonproduction in the reaction $pp \rightarrow K^* \Lambda p$. In fig.2 the scheme of the startdetector system consisting of two segmented layers of scintillators, a ring microstrip detector and a hodoscope of scintillating fibres is shown together with a $pp \rightarrow K^* \Lambda p$ Monte Carlo event with subsequent decay $\Lambda \rightarrow p\pi^-$.

Due to the relatively low momentum available up to now for the external COSY beam, the threshold of the $K^* \Lambda$ -production at 2.34 GeV/c could not be reached. However the performance of the detector system could be tested in detail especially by investigating the elastic scattering $pp \rightarrow pp$ and the pion production $pp \rightarrow \pi^+ d$. The results of these analysis confirmed the design parameters of all detector components. In fig.3 a typical elastic event is shown in a θ - ϕ -projection. The two colinear tracks can be clearly recognized by the "crossing points" of the different detector elements.

Fig.4 shows the $\Delta\theta$ - $\Delta\phi$ -distribution for a sample of two track events where only the information of the fibre hodoscope is used. The sharp peak corresponds to the elastic scattering on hydrogen. The small background is mainly due to reactions on the carbon content of the used CH_2 -target.

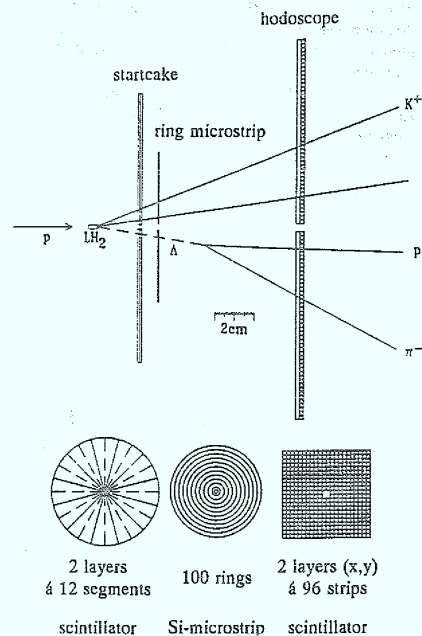


Fig.2: "Erlangen startdetector": detector setup and Monte Carlo event $pp \rightarrow K^* \Lambda p \rightarrow K^* p \pi^- p$

An important feature of the startdetector system is its time resolution. Using the two segmented layers of scintillators in the case of one track events an effective value of about 450 ps (FWHM) was obtained; for events with two tracks as in the Λ -production a resolution of about 300 ps (FWHM) and less can be extrapolated.

The most important point for the observation of the hyperon production is the reconstruction of primary and secondary vertices. Using the information of all detectors of the Erlangen system and the Quirl one gets a vertex resolution of 1-2 mm (FWHM). Hodoscope and quirl information are sufficient to achieve a resolution in vertex reconstruction of about 5 mm.

This ensures the possibility to reconstruct the delayed Λ -decay vertex.

During the test runs it was also shown that the trigger for the hyperon production $pp \rightarrow K^* \Lambda p$, based on the increase of charged tracks in the Λ -decay, yields a sufficiently strong suppression of background events.

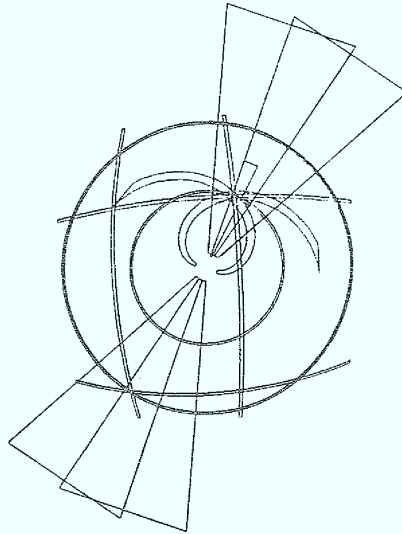


Fig.3: eventplot in theta-phi-projection, elastic scattering

Moreover extensive Monte Carlo investigations have clearly shown that the existing setup of the Erlangen Startdetector system together with the Quirl (distance of about 1 m) covers the full phase space of the reaction $pp \rightarrow K^* \Lambda p$ from the threshold up to 3.3 GeV/c with high efficiency.

Altogether the test measurement at COSY shows that the existing apparatus consisting of the Erlangen detector system together with the Jülich Quirl satisfies all requirements for successful hyperon production runs.

In September 1994 COSY reached its design momentum of 3.3 GeV/c. The TOF-collaboration is now waiting for an extracted beam with momentum above the Λ -production threshold.

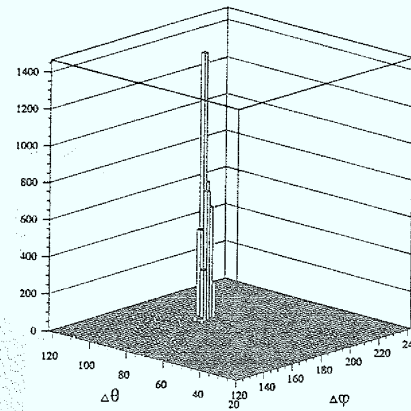


Fig.4: reconstructed two-prong events as a function of relative angles

A UV-Laser based Test System for the COSY-TOF-Spectrometer

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For online-monitoring and calibration purposes we are developing a UV-laser based test system [1]. The UV-light is led and distributed to scintillator counters by optical fibres, where light pulses similar to those originating from charged particles are produced. To obtain a uniform distribution of propagation time and light intensity in all detector channels we constructed a dedicated reflector head to couple the fibres to the scintillation counters (fig.1).

It consists of a support to fix the fibre, an aperture and a mirror (alu mylar foil) to reflect the light towards the photomultiplier. By adjusting the mirror it is possible to regulate the light intensity. A setup of 32 fibres was connected to the counters of the start detector and tested during a test measurement in december 1994. The results we achieved are shown in figures 2 and 3.

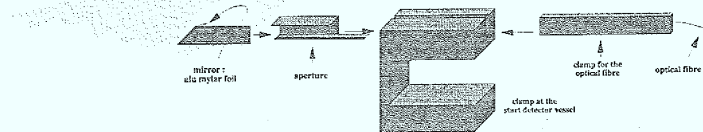


Figure 1: Schematic drawing of the reflector head

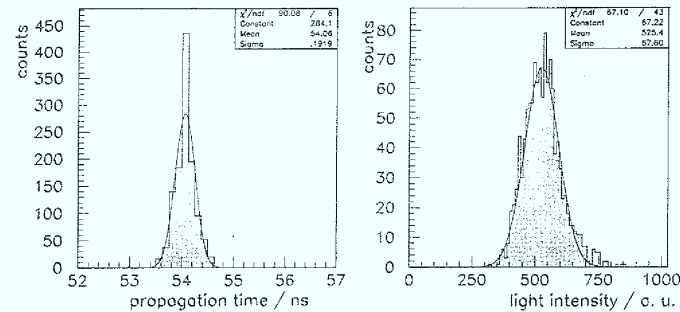


Figure 2: Typical TDC and QDC histograms of one channel. The time resolution is in the order of the intrinsic one of the photomultiplier XP2020.

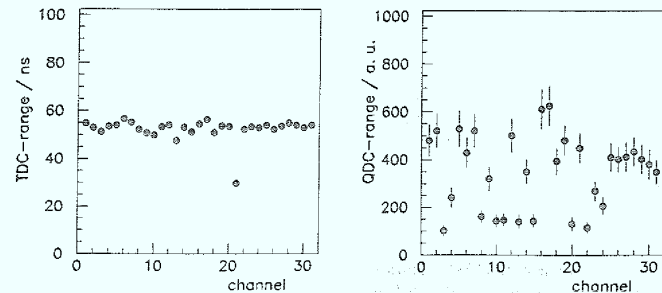


Figure 3: Time spread and light output of all 32 channels of the start detector.

Reference:

[1] P. Hermanowski, Diplomarbeit, Bochum 1993

R. Loke, Diplomarbeit, Bochum 1994

P. Hermanowski et al., AR, FZ Rossendorf 1992 and 1993

This work has been supported by the BMFT and KFA

Test facility for methodical investigations of long scintillator strips for use in the COSY-TOF-Spectrometer

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for the COSY-TOF collaboration

In next years a step by step extension of the COSY-Time-of-Flight spectrometer [1] is planned towards the final version of a 4π detector.

The existing component of the endcap (central hodoscope) and the component presently under construction (ring-like hodoscope) consist of three layers of scintillating elements (one layer of wedge-like elements and two layers of right- and left-twisted Archimedian spirals correspondingly). Position determination is achieved by arranging the elements such as to form a grid of triangular pixels.

An analogous configuration for the barrel sections (one layer of straight scintillator strips, two layers right- and left-bend strips correspondingly) turned out to be problematic because of the light pulse attenuation length in the bend elements being too small and because of technical difficulties in manufacturing the four meters long bend scintillator strips.

Therefore there was the necessity of looking for an alternative concept avoiding the problems mentioned on one hand and maintaining the good spatial resolution properties determined by the geometrical pixel structure of the original version on the other hand.

In the new concept special care was taken to solve the problem of on-line position determination.

In a first step towards the realization of the new concept a methodical measurement was performed at CERN [2] to test two-sided read out of long straight scintillating strips using minimally ionizing particles. The position

information of the detected particle was deduced from the signal time difference between the two ends.

A time resolution of $\sigma = 0.45$ ns and a corresponding spatial resolution of $\sigma = 5$ cm were obtained.

The on-line position measurement was carried out by a newly developed real time to digital converter (RTDC) [3] with conversion times in the range of a few nanoseconds.

From the test measurements at CERN one can draw the preliminary conclusion that the new concept could be an alternative for the original one.

The new concept would mean that the three layer arrangement can be replaced by a single layer consisting of straight elements only.

The final decision, however, on whether to realize the new concept or not requires the solution of quite a number of additional problems such as

1. Technical realization of 90° -read out
2. Check of the suitability of the existing light guide vacuum flanges
3. Vacuum behaviour of detector materials
4. Compensation of mechanical expansion of scintillator material caused by temperature increase
5. Ageing processes of scintillators under vacuum conditions

To investigate these problems a test facility was designed and manufactured (fig. 1).

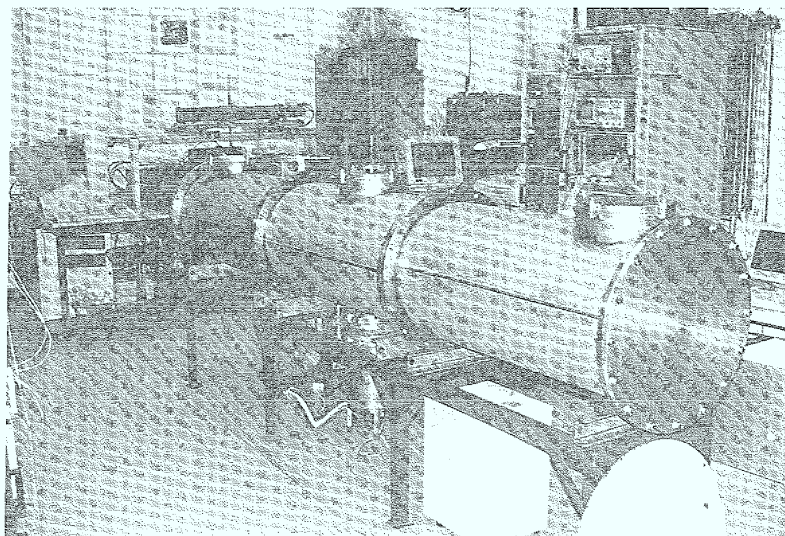


Fig. 1: Test facility for the investigation of long scintillator strips

Its main component is a vacuum vessel which consists of three segments with total length of 3300 mm and inner diameter of 500 mm.

Each segment has three radial flange connections (K250) for attaching the light guides.

Different combinations of the segments in addition to different flange positions allow to adapt the apparatus to scintillators up to a length of 3280 mm.

The configuration shown in fig.1 is adjusted to test scintillators of maximum length.

For the investigation of the vacuum behaviour of the detectors the vessel can be evacuated down to high vacuum values of 10^{-8} mbar by means of a turbomolecular pump of type TURBOVAC 450 (Leybold).

To generate fine vacuum the central fine vacuum pipe system of the laboratory is used. To perform ageing investigations of the detector materials pressure values in the range of 10^{-1} to 10^{-4} mbar must be kept constantly for long time periods (weeks).

The pressure control in this case is carried out by a feed back system using pressure sensors (Pirani, Penning), a computer programmable control system (type OMRON CQM1) and an adjustable magnetic valve.

The computer control of the system allows to automatically perform the following modes of operation: fine vacuum, high vacuum, pressure control, venting, standby and error checks.

The control of the vacuum system is carried out menu-driven with a touch screen display (type OMRON NT20).

The control programm comprises different options of diagnosis such as long term recording of measured vacuum data and automatic measurement of leaking and degassing rates.

To determine the influence of temperature changes on the detector elements a possibility is provided to heat the vacuum vessel.

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TU Dresden

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- [2] P. Michel et al., Annual Report 1993, Institut für Kern- und Hadronenphysik, FZ Rossendorf 1993; Annual Report 1993, IKP, Jül-2879
- [3] F. Gabriel and P. Michel, FZR-30, FZ Rossendorf

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The first beam time at the COSY Time-of-Flight spectrometer TOF in spring 1994 with a 455 MeV proton beam was used mainly to perform functional tests of the various components of the spectrometer. In this first test measurement the TOF spectrometer was used in a configuration consisting of the central-endcap-hodoscope (stop detector) and (alternatively) one of the two versions of the start detector [1],[2]. Here a short account is presented on the results of methodical investigations of the Rossendorf start detector MARS. The configuration used in the measurements is shown schematically in Fig.1.

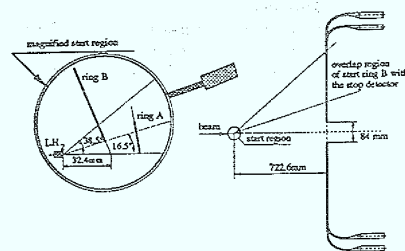


Fig. 1: Schematic view of the configuration used in the first measurement

In this configuration the complete B-ring and four elements of the A-ring of the start detector were used. The four segments of the A-ring were arranged in a cross-like geometry. They served for monitoring the primary beam.

From beam diagnostics as well as data analysis it turned out that the (up to now uncooled) beam was not of the desired quality. Data analysis revealed that only 6% of all hits detected can be assigned to pp-reactions. All other hits had to be discarded since there was no geometrical link between target position, B-ring and the central hodoscope for these hits. This means that the particles registered originate from beam halo rather than nuclear reactions.

Since TOF is a time-of-flight spectrometer one of the crucial properties of the spectrometer is the time resolution. In Fig.2a a TDC spectrum is presented showing the time of flight between a start detector segment and a certain pixel of the central hodoscope (timing information is extracted from the wedge-like element). The TDC spectrum was generated in data analysis on the basis of the parasitic beam halo particles because these particles have a small energy spread. The time resolution obtained under these conditions turned out to be 1.6 ns (FWHM). This time resolution is composed of both the start and the stop detector resolution. The arrangement used in the test measurement did not allow to separate the time resolution of the start detector

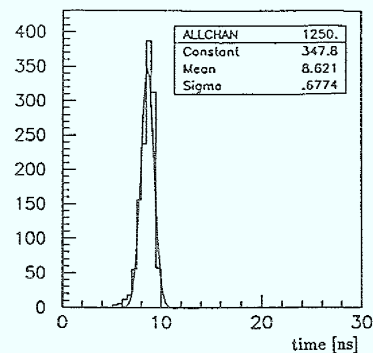


Fig. 2a: Start-stop time-of-flight resolution

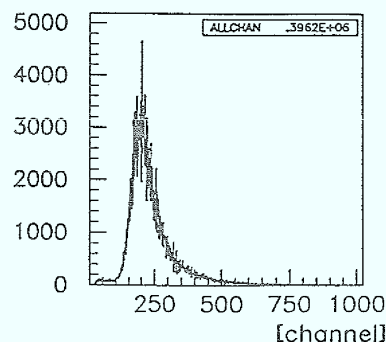


Fig. 2b: Start detector ADC spectrum

from the total resolution. In fact there was no necessity to do this since the time resolution of the start detector had been measured accurately under clean conditions in an earlier experiment [1].

As concerns the energy resolution of the start detector (Fig.2b) so it is quite satisfactory with signals well above noise.

In a next step the geometrical adjustment of the start detector was investigated. Neighbouring segments of the start detector B-ring are designed to have small overlap to assure 100% efficiency if it is operated as a veto detector [5]. Fig.3 illustrates a central projection of the overlap regions of the B-ring measured with particles arising from the target point. The signature chosen for this projection was two hits ($M_B = 2$) in neighbouring segments of ring B. The obvious left-right

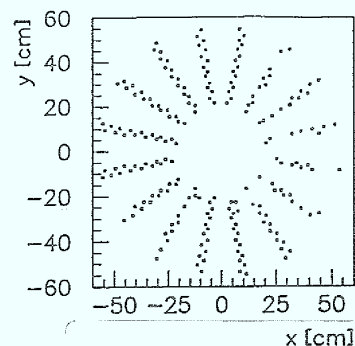


Fig. 3: Central projection of the overlap regions of the start ring on the stop detector from the target point

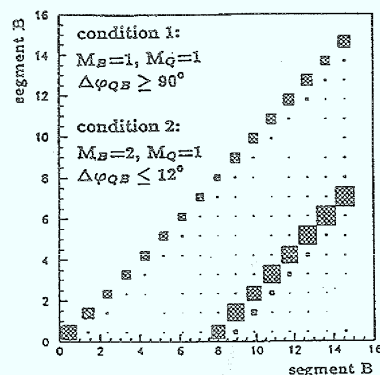


Fig. 4: Coincidence matrix of start detector

asymmetry in Fig.3 is caused by beam imperfections (asymmetric beam halo).

The results of the data analysis of the measurement are represented in form of a correlation matrix (coincidence matrix) for the start detector segments (Fig.4). This kind of representation was introduced in connection with an earlier test of the start detector MARS [5]. In Fig.4 all those hits are represented of the total bulk of data which fulfill the 2 conditions indicated in the picture. Condition 1 says: 2 particles are registered (one hit in the B-ring, one hit in the central hodoscope) with relative azimuthal angle $\Delta\varphi_{BQ} > 90^\circ$. Particles fulfilling condition 1 correspond to the entries on the main diagonal in Fig.4. Part of the hits on the main diagonal (with signature $\Delta\varphi_{BQ} \approx 180^\circ$, $\Delta\Theta < 16.5^\circ$) are due to elastic pp-scattering.

The hits of the coincidence matrix outside the main diagonal correspond to condition 2. According to this condition two particles are detected in ring B with one of these particles registered additionally in the central hodoscope. The main feature of the coincidence matrix is the straight line starting at the point (0,8). This

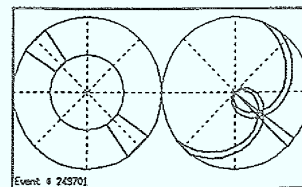


Fig. 5: Hit pattern of an elastic event

line corresponds to elastic pp-scattering. It is obvious from Fig.4 that elastic scattering can easily be separated from background.

In Fig.5 as an example the hit pattern of an elastic event is shown (2 hits in the start detector B-ring, 1 hit in the stop detector).

For data analysis an adapted version of the Bochum TOF data analysing program system was used which was developed by members of the Bochum-TOF-group on the basis of a program package from M. Gasthuber (University of Hamburg).

This work has been supported by the BMFT and the Forschungszentrum Jülich.

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Computer Controlled Vacuum System for the COSY-TOF-Spectrometer

R. Klein, P. Michel¹, K. Möller², Ch. Nake, B. Naumann¹, L. Naumann³, A. Schamlott¹ and A. Schülke¹
for the COSY-TOF collaboration

To minimize background production, energy loss and multiscattering the start and stop detector components of the COSY-TOF-spectrometer [1] are placed in a modular vacuum vessel with a maximum volume of 60 m³. Vacuum pressure less than 10⁻² mbar is necessary for precise nuclear physics experiments. The installation of the LH₂-target [2] in each of the two versions of the start detector [3,4] requires a sufficiently low vacuum pressure in the order of 10⁻⁵ mbar around the cold target cell to suppress the growth of residual gas condensates at the target windows [5]. Such condensates would falsify the original composition of the target material. The spectrometer vessel operates at fine vacuum conditions (10⁻² mbar). Between the Rossendorf start detector and the vessel a Mylar foil is installed to separate the high vacuum and the fine vacuum regions. To reduce background from the primary beam the

The layout of the COSY-TOF-spectrometer vacuum system is shown in fig.1.

To make the control system independent of operator mistakes or system crashes a memory programmable controller - system OMRON [6] - was chosen. It consists of a CQM1 central processing unit and a lot of binary and analog I/O-units. A pseudo graphical touch screen panel (OMRON NT20) is used as a modern user interface. Controller unit, touch screen panel, vacuum measuring device and contactors for the pumps are housed in a 19 inch crate. The electropneumatic valves are controlled by 24 V DC.

At changing the experimental set-up it is necessary to separate the control and the vacuum system. Therefore all pumps except the turbo molecular pump and most of the valves are placed on a mobil platform and the control system is mounted in a mobil rack in the TOF

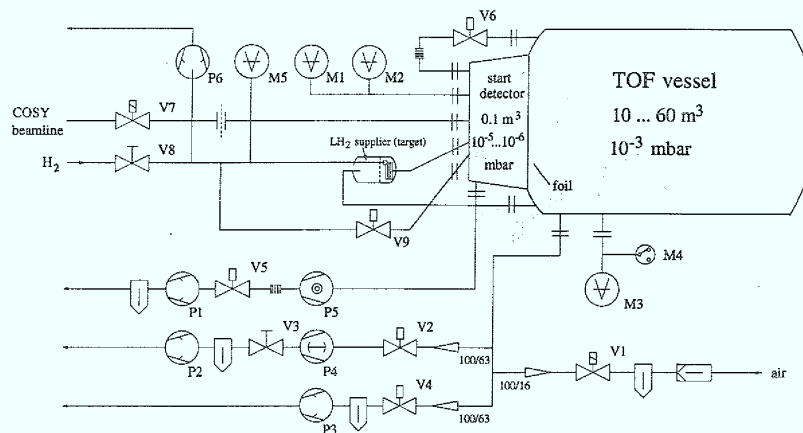


Fig.1: Schematic layout of the COSY-TOF-Spectrometer vacuum system
V - valve; P - pump; M - vacuum gauge;

Mylar foil was chosen as thin as possible (20 µm). To avoid limitation of the acceptance of the spectrometer a large area foil is necessary (ø600 mm). In the normal working regime of the spectrometer the pressure difference between the start detector and the vessel is small and no visible deformation of the foil occurs. But if in the case of technical trouble the pressure difference rises to a few millibars the foil is deformed substantially. This can lead to serious damage of the sensitive start detector elements (hollow light guides, target cell etc.). For this reason a bypass valve (V6) is provided to limit the pressure difference between the start detector chamber (M2) and the vessel region (M3). To automate vacuum equipment as a whole a computer controlled system was designed and installed.

area.

For remote check of the vacuum system while the beam is active it is possible to display the system state, measuring values, recording informations and errors via a V24-line on a PC in the TOF counting room.

To avoid mishandling by unauthorized persons a remote control of the system from this room was not installed intentionally.

The control of the vacuum system is carried out menu-driven with the touch screen panel. The main menu entries *FineVac*, *HighVac* and *Venting* make it possible to evacuate as well as to ventilate the vacuum system automatically. In these modes of operation a complex error supervision is active which controls the following conditions:

- pressure difference between the two recipients
- limit positions of the valves
- response signals of the pumps (e.g. overload)
- trouble information of the vacuum measuring device
- availability of compressed air
- 24 V DC power supply

All these errors cause closing valves (except bypass valve), opening the bypass valve and switching off the pumps.

If the high vacuum is reached an enable signal is generated to allow filling the LH₂-target. For venting a signal "LH₂-target empty" is required, moreover a signal for COSY-Control is generated. In the Service mode each valve and each pump can be steered separately. The system state and the pressure values can be displayed. The error supervision is suppressed. This mode is protected by a keycode, so an unintentional operating is impossible. In all operating modes the state of the vacuum system can be displayed in a vacuum scheme. It is also possible to store and display all pressure values over the time (up to a time scale of a few days - Recorder, see fig. 2) and to show pressure gradients for instance to measure leaking rates or the pump performance.

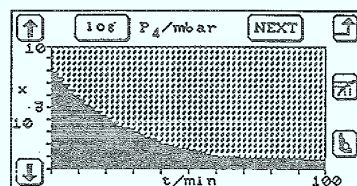


Fig. 2: Recorder information of vessel pressure (displayed on the NT20-panel)

The described vacuum control system was tested successfully during a two weeks COSY-TOF run in 1994. It is planned to integrate the LH₂-target handling into the control system in near future.

This work has been supported by the BMFT and the Forschungszentrum Jülich.

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Residual Gas Analysis of the Liquid Hydrogen Target in the Start Detector Device of the COSY-TOF-Spectrometer

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The installation of the LH₂ target [1] in the start detector vessel [2] of the COSY-TOF-Spectrometer requires a sufficiently low vacuum pressure, to suppress the growth of residual gas condensates at the target windows. In a first step the components of the residual gas inside the start detector vessel were analyzed after a long pumping duration, when the total pressure reached the asymptotic limit of approximately $3 \cdot 10^{-6}$ mbar. This measurement was carried out with the Transpector Gas Analysis System from Leybold Inficon Inc. [3]. The mass spectrum of the residual gases is shown in Fig. 1 as a pattern of peaks on a plot of ion current as a function of ion mass to charge ratio. To calibrate the mass to charge ratio, nitrogen gas was additionally filled in the vacuum vessel. The dashed line in the mass spectrum shows a growth of peaks in the channels for nitrogen. The partial pressures of the major residual gases were calculated from the mass spectrum and are shown in Table 1.

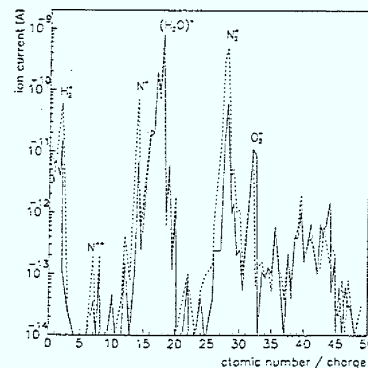


Figure 1: Mass spectrum of residual gases in the start detector device [4].

Table 1: The partial pressure of residual gases calculated from the spectrum in Fig. 1.

Residual gas molecules	partial pressure [mbar]
H ₂ O	$6.8 \cdot 10^{-7}$
N ₂	$1.3 \cdot 10^{-6}$
O ₂	$3.1 \cdot 10^{-7}$

In a second step the contamination of the windows of the liquid hydrogen target with the main residual gas components was investigated. The target cell filled with liquid hydrogen at 16K and 200 mbar was mounted in the start

detector vessel with a total pressure of $3 \cdot 10^{-6}$ mbar. Water, nitrogen and oxygen were frozen out on the cold target surface and an ice layer became visible. The contamination of the hydrogen target increases with the exposure time for $3 \cdot 10^{-6}$ mbar as shown in Fig. 2. The time duration to eliminate the contamination of nitrogen and oxygen by target heating is just 5 minutes and for water 1 hour.

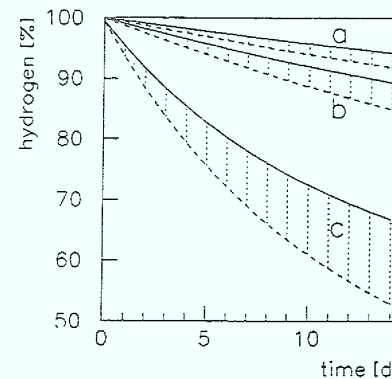


Figure 2: Calculated percentage of hydrogen atoms to all atoms in the target region for various target length: a - 10 mm, b - 5 mm and c - 1 mm at total pressure $3 \cdot 10^{-6}$ mbar for window foil thickness of 2 μ m; solid lines: all hydrogen atoms were considered; dashed lines: only the well defined number of hydrogen atoms from the target volume and the foils are taken into account [4].

Because proton-proton cross section measurements with high accuracy at the COSY-TOF-Spectrometer require well defined hydrogen targets of high purity, the target contamination must be suppressed. There are some possibilities to do that:

- to improve the vacuum ($p \leq 10^{-7}$ mbar);
- to eliminate the ice layer by heating ($T \geq 170$ K).

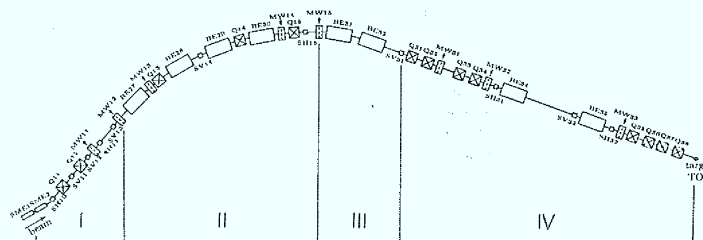
¹ FZ Rossendorf, Dresden

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If the waist of the beam entering the last four quadrupoles is far away (typically more than 50 m). In this case the acceptance is limited by the beam tube in the vertically deflecting dipoles BE34 and BE35.

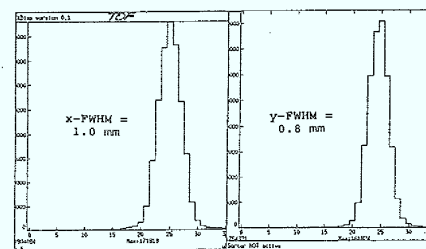


A two steerer method was applied to adjust the beam in position and angle near the entrance of the beam line. In this case the magnetic septum is used as a first steerer to adjust the beam in the middle of the second steerer SH10, which is imaged (in the optical sense) on a position sensitive device. So the center of the beam in SH10 becomes directly visible. Then the angle is adjusted by the second steerer. The beam of COSY deviates 8 mrad from the axis of the beam line, depending somewhat on the extraction conditions. Calibration of the quadrupoles is described in [2]. In the thin lens approximation one has

with $a = (0.0108 \pm 0.0002) [(\text{GeV}/c)/(\text{m} \cdot \text{A})]$.

The two dipole section (III) can be adjusted by turning the beam around a point on the position monitor MW31. This is the case if $dI(MDBE31)/dI(MDBE32) = -0.657$ is in good agreement with the inverse distance ratio of the dipoles to the monitor $r = 0.64$. The vacuum chambers near the dipoles BE31 - BE32 are limiting the acceptance.

The matrix elements calculated with the program TRANSPORT [3] are in agreement with the current settings to about 1 %. A dataset containing all current settings is listed in file ex941212.ppic in the COSY control computer. Emittance measurement of the beam from COSY show that the transport of the total beam can not be expected. The beam path brings a beam of appropriate extension to TOF, but it is not yet optimized in all details. An example for the size of beam spot on the target in TOF is shown in Fig. 2.



References

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- ** FZ Rossendorf
- *** Uni Bochum

Reactions with strangeness non-conserving Λ -production would allow the observation of a pure weak process not via interference effects but in a clean direct way. It has been proposed [1] to measure the strangeness violating weak Λ -production at COSY-TOF in the reaction $dp \rightarrow \Lambda pp$.

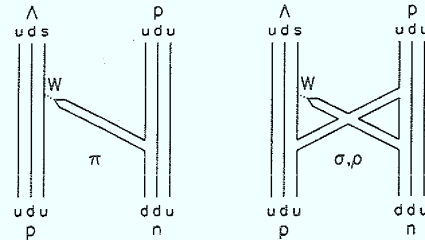


Figure 1: Quark line diagrams of weak Λ -production

Hypernuclei decay experiments are somehow the inverse. They have shown that a simple description of the underlying decay process is not sufficient: Λ decaying into nucleon and pion, with the pion interacting with a neighbouring nucleon. This simple model incorporates a spin and isospin dependence caused by the pion exchange, a dependence which is not seen in the experimental data. Additional processes have to be taken into account, for instance σ -, ρ -, ω -, or Kaon-exchange [2]. In order to obtain any experimental data about these decay modes which obviously cannot be seen in a free Λ -decay it would be favourable to have a reaction where this weak process could be seen isolated and not inside a nucleus. A solution seems to be the weak strangeness production in nucleon-nucleon scattering, i.e. Λ -production below any strangeness conserving energies.

The reaction has a very low cross-section. The branching ratio can be estimated from the ratio of typical time constants for hadronic and weak interaction to be $10^{-23}s/10^{-10}s \sim 10^{-13}$ (see also [2]). It is obvious that for a measurement an extremely clean environment is needed. This might be obtained by two conditions: first, a reaction which runs at an energy below any threshold of strangeness conserving strong Λ -production, second, the Λ -decay which is the main signature for strangeness identification must happen in vacuum in order to avoid background reactions which could fake a Λ -decay. The first condition is fulfilled by the reaction $dp \rightarrow \Lambda pp$. Its threshold of 1.5406 GeV/c is far below strong Λ -production even on heavy nuclei (1.726 GeV/c). The advantage of measuring d on p instead of $pd \rightarrow \Lambda pp$ is the fact that the emission angles of all particles (except the decay pion) are below 35° in the whole measurable momentum range. The second point is fulfilled by using the TOF-detector with its internal vacuum.

We want to begin these studies with the same setup like for $pp \rightarrow pK^+\Lambda$, a reaction which is already included in the experimental program of COSY for 1995 [3]. Here,

the experiment consists of the liquid hydrogen target with a small start detector behind it. The following rectangular scintillator hodoscope in combination with the outer detector elements 'quirl', 'ring', and 'barrel' (all 3-layer scintillator hodoscopes) will allow a precise track reconstruction and time-of-flight measurement. A limit of $< 10^{-9}$ for the branching ratio should be rather easy to reach like this.

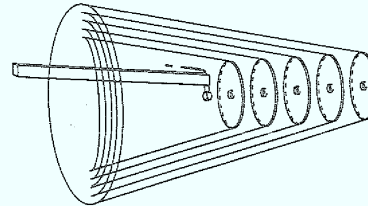


Figure 2: Target and decay spectrometer

For higher sensitivity we plan to improve the detectors step by step and finally use the experimental setup proposed for hypertriton lifetime measurement [4]. Following the liquid hydrogen target a stack of silicon counters (a decay spectrometer) with a 'quirl'-like surface structure is used for track and therefore decay vertex determination. These silicon disks with ~ 70 mm diameter will be mechanically and electrically supported by cones made of copper plated capton foils. A spherical scintillator hodoscope covering 2π in laboratory system surrounds the inner detector section. In conjunction with the outer scintillator hodoscopes it allows for precise track reconstruction and time-of-flight measurement also for decay pions. An additional energy detector behind the 'quirl' will improve energy resolution.

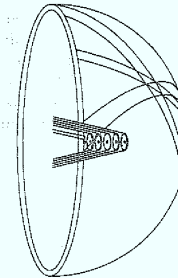


Figure 3: Spherical scintillator hodoscope containing target section

References:

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Monte Carlo Studies for Recent COSY-TOF Experiment Proposals

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Several experiments in the field of strangeness production have been proposed to be performed in a modified version of the time of flight spectrometer TOF at COSY /1,2/.

Monte Carlo studies have been made to verify the feasibility of these experiments regarding acceptance, resolution, and separation against competitive reaction patterns. The following reactions were objects of the study:

1. $dp \rightarrow pp\Lambda$
2. $pd \rightarrow K^+ {}^3\text{H}_\Lambda$

Both cases are characterized by the occurrence of delayed particle decay, so that the detector is triggered by particles of different type and origin. Therefore the detector must be enabled to distinguish between different particle types and origins.

To find the origin of a particle, its track must be detected in at least two points, and the directional information of adjacent particles has to be used in a vertex reconstruction. For particles coming from the primary collision the vertex must be located within the target. A start detector near the target can provide the directional information. Decay products on the other hand may be generated outside of the start detector or may escape between the layers of the suggested start detector, which has been designed as a stack of Si-disks, having a whirl-like readout structure. Thus, for decay products an additional detector component, designed as a 3 layer whirl-like scintillator hodoscope, has been introduced to provide the start signal for the time of flight measurement and spatial information for the track and vertex reconstruction.

Recognition of different particle types requires measurement of the complete four-momentum of the particles. The time-of-flight method only provides information about β . The simultaneous measurement of the energy loss allows to distinguish between particle types, but significance depends strongly on the quality of both measurements. A scintillator calorimeter, stopping the particles in forward direction gives an absolute measure of the particle energy with a design resolution of 4%.

In our studies we used the MC4 program. MC4 integrates a set of Monte Carlo routines for the simulation of

- primary collisions,
- particle tracking in standard geometries,
- and interaction with matter with an application dependent set of routines used to
- describe non-standard geometries,
- describe the technical behaviour of the detector as
 - light attenuation in scintillators,
 - thresholds,
 - etc,

generating data similar to the data recorded by a data acquisition system in a real experiment

- a set of reconstruction routines to be tested in the MC environment.

Thus, MC4 can be used to compare reconstructed events with true events (coming from MC), to study the quality of reconstruction methods and the influence of technical parameters of the detector.

Looking at the reaction $dp \rightarrow pp\Lambda$ we introduced the geometry of the new detector elements into the application code used for the standard TOF design. Figures 1 and 2 show a schematic view of the two different setups which have been compared.

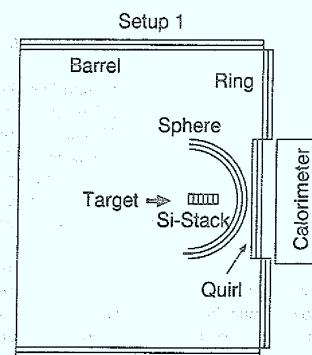


Fig. 1: Experiment setup 1 with scintillator calorimeter

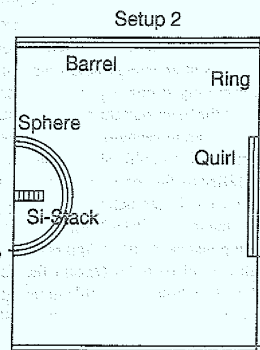


Fig. 2: Experiment setup 2 (time-of-flight setup)

Setup 1 covers a greater solid angle. The measurement of protons coming from the primary collision is fairly good, because of the energy resolution of 4% in the scintillator calorimeter. The energies of decay products are rather approximately measured because they are measured with the time-of-flight method, and the flight distances are rather short.

Setup 2 gives better results for the time-of-flight measurements but the backward angles are not covered at all and the calorimeter is removed because of its small solid angle.

Figure 3 shows the missing mass reconstruction of the Λ particles alone with the acceptance rates for the reconstruction of the production and the decay term. The overall efficiency of the system is determined by the weakest part of the reconstruction. Figure 4 shows the same for setup 2.

As mentioned above, some assumptions have been made for the behaviour of the non-existent detector components:

- The calorimeter resolution has been set to 4%.
- The time resolution of the spherical scintillator hodoscope has been neglected.

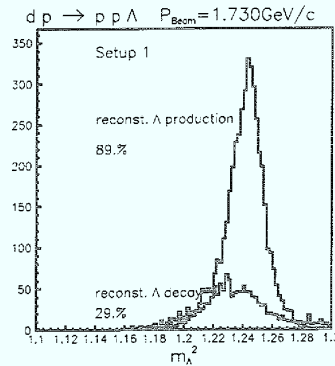


Fig. 3: Missing mass reconstruction of Λ in the setup 1-experiment

On the other hand track and vertex reconstruction can still be improved using fitting routines after a careful study of the typical errors in the new components and by taking into account special cases like decay products stopped in the calorimeter or the sphere for which a much better energy reconstruction can be achieved.

Looking at the reaction $pd \rightarrow K^+ {}^3\text{H}_\Lambda$ (hypertriton production) we used the geometry outlined in setup 2. The final state particle ${}^3\text{H}_\Lambda$ decays in two modes of unknown fraction:

1. ${}^3\text{H}_\Lambda \rightarrow pd \pi^-$
2. ${}^3\text{H}_\Lambda \rightarrow {}^3\text{H} \pi^-$

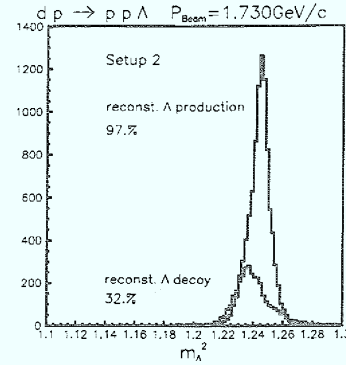


Fig. 4: Missing mass reconstruction of Λ in the setup 2-experiment

The decay modes were tested in separate runs. Thus ${}^3\text{H}_\Lambda$ can be reconstructed using

- the K^+ and the beam momentum
- the measurement of the final state particles of decay mode 1
- the measurement of the final state particles of decay mode 2

Figures 5 and 6 show the results of the missing mass reconstruction.

In both cases the K^+ can be used to improve the reconstruction resolution. The acceptance is a function of the decay channel reconstruction, because only the decay channel provides a unique pattern to recognize the reaction.

Figure 7 illustrates the capability of MC4 to compare reconstruction results with true events. Shown here are differences of the reconstructed ${}^3\text{H}_\Lambda$ 4-momentum and the true Monte Carlo values.

References:

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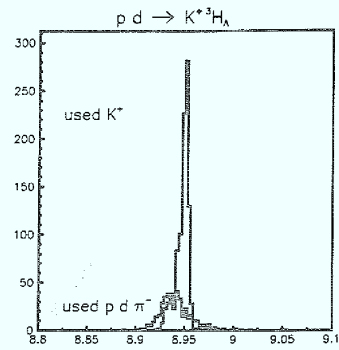


Fig. 5: Missing mass reconstruction of ${}^3\text{He}$ in the setup 2-experiment. The different used reconstruction modes are shown.

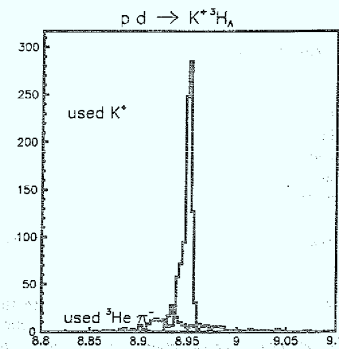


Fig. 6: Missing mass reconstruction of ${}^3\text{He}$ in the setup 2-experiment. The different used reconstruction modes are shown.

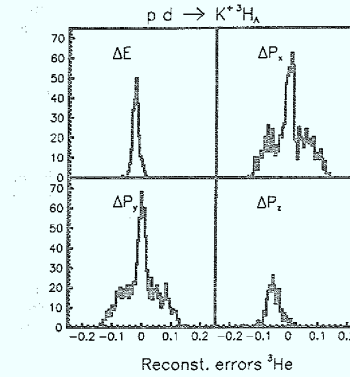


Fig. 7: Differences of the „measured“ and calculated 4-momentum vector components of the ${}^3\text{He}$ final state particle

The EDDA Experiment: Proton-Proton Elastic Scattering Excitation Functions at Intermediate Energies [1]

J. Bisplinghoff^a, F. Hinterberger^a and W. Scobel^b for the EDDA collaboration [2]

The EDDA experiment is designed to provide a high precision measurement of proton-proton elastic scattering excitation functions ranging from 0.5 to 2.5 GeV of (lab) incident kinetic energy. It utilizes the proton beam of the cooler synchrotron COSY operated by KFA Jülich.

EDDA is meant to make a significant contribution to the global proton-proton phase shift analysis, particularly in the energy domain above 1.2 GeV, where existing data is sparse. It will also test the energy range in which meson exchange models are a valid approximation for the nucleon-nucleon interaction. Furthermore, EDDA is highly sensitive to the existence - or non-existence - of dibaryonic resonances with any appreciable coupling to the elastic channel. Hence the acronym: Elastic Dibaryons Dead or Alive.

Both for phase shift analysis as such, and its use in identifying dibaryonic resonances, it is mandatory to measure not only spin-averaged cross sections, but also spin observables such as analyzing powers, spin correlation coefficients etc.. This requires a polarized COSY beam as well as a polarized EDDA target. A polarized ion source has been built and fitted [3] to the COSY injector cyclotron, and the development of the polarized COSY beam may commence. The polarized EDDA target is a polarized atomic beam device now under construction [4]. With reasonably achievable atomic beam densities in mind, EDDA was conceived as an experiment using the internal (recirculating) COSY beam to obtain sufficient luminosities. Spin-averaged cross sections are measured with a CH_2 fiber target [5]. Data collection proceeds during synchrotron acceleration in a multi-pass technique, so that a full excitation function is measured in each accelerator cycle. Statistical accuracy is obtained by averaging over many cycles. Elastic pp events are identified by kinematic coincidence. The luminosity is monitored as a function of the COSY beam energy by means of the electrons knocked out of the target. An integral measurement of the secondary electron current is performed as well as detection of high energy δ -electrons emerging at 40° using PIN-diodes.

The EDDA detector [6] is shown in Fig. 1 in a schematical fashion. It consists of two cylindrical double layers covering 30° to 150° in $\theta_{\text{c.m.}}$ for the elastic pp channel and some 85% of the full solid angle. The inner layers are composed of scintillating fibers which are helically wound in opposing directions. The outer layers consist of scintillator bars which are running parallel to the beam axis and are surrounded by scintillator semi-rings. The scintillator cross sections are designed so that each particle traversing the outer layer produces a signal in two neighbouring bars and rings. Analysis of the fractional light output may then be used to determine the angle, at which the particle traverses the outer layers, much more

accurately (by a factor of 5) than would be possible on the basis of detector granularity alone. Combined with the spatial resolution of the 2.5 mm diameter scintillating fibers, this provides for vertex reconstruction to a precision of some very few millimeters.

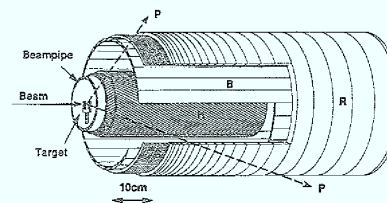


Fig. 1 The EDDA-Detector (not to scale): Target (CH_2 fiber suspended from fork, later polarized atomic beam), Beam pipe (2 mm Al in sensitive region), H Scintillating Fibers wound in opposing helicities, B Scintillator bars, F + R Scintillator semi-rings

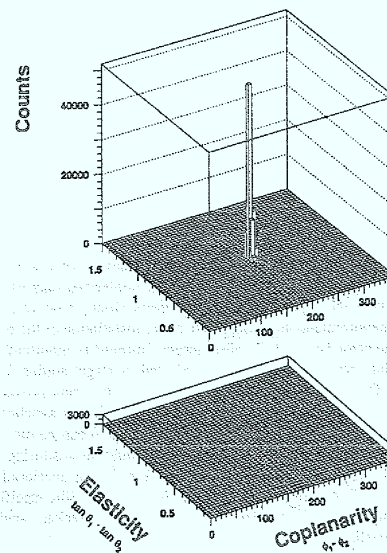


Fig. 2 Kinematic distribution of EDDA events as obtained with a CH_2 (upper part) and with a carbon (lower part) fiber target ($T_p = 793 \text{ MeV}$)

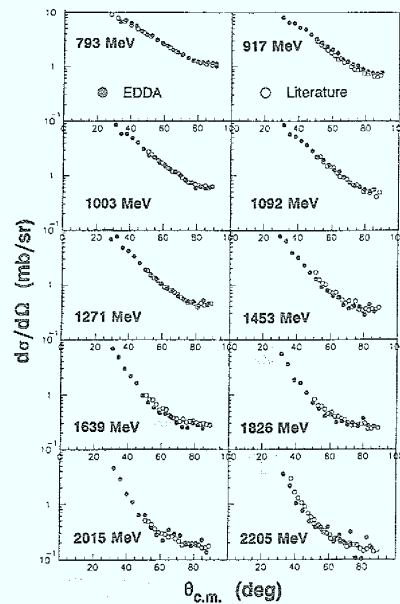


Fig. 3 Angular distributions of elastic proton-proton scattering events as obtained (preliminary, analysis incomplete) by EDDA and as available from the literature for the bombarding energies (lab) indicated. The EDDA data were normalized as a function of bombarding energy using the δ -electrons recorded in PIN diodes as luminosity signal.

Vertex reconstruction is essential, when a polarized atomic beam target is used. Unpolarized data can be taken with the outer detector layers alone, however, and measurements have begun in 1994. Installation of the outer detector layers, CH_2 fiber target, luminosity monitors and data acquisition [7] was completed at target station TP2. COSY reached its design energy, and the beam proved to be nicely stable, so that data collection during synchrotron acceleration poses no problems. The beam can be tuned to roughly compensate small angle Coulomb scattering off the fiber target by adiabatic cooling during acceleration. Detector performance was shown to be within specifications, and luminosity monitoring was checked against existing data [8] obtained with external targets and "beam counting".

Fig. 2 illustrates the identification of elastic proton-proton events by kinematic coincidence. The events in the conspicuous peak (upper part of Fig. 2) show the kin-

ematics of elastic proton-proton scattering, coplanarity with the beam and the characteristic combination of polar angles. Quasielastic background due to the carbon content of the CH_2 fiber target, on the other hand, results in a wide and rather structureless distribution (lower part of Fig. 2) that can be measured separately using a carbon fiber target and subtracted out.

Fig. 3 shows some of the elastic scattering angular distributions measured during an EDDA run in the fall of 1994 at the beam energies indicated and employing the multi-pass technique. Due to still insufficient statistics, the data was sorted into 20 MeV wide bombarding energy bins. It was normalized to the luminosity signal obtained from the PIN diode monitor system. Normalizing to the secondary electron current gives essentially the same result. Fig. 3 compares our data to existing cross sections [8] measured with external targets and careful "beam counting" for luminosity determination. The agreement found proves the EDDA luminosity monitors to be accurate and reliable. The angular resolution is about 2° (center of mass).

EDDA is now ready to take unpolarized data at full scale in order to reach the statistics required for 1% accuracy in 2 MeV wide beam energy bins, the design goal. This will constitute the center of effort in 1995, along with preparing to measure spin observables, i.e. installing the inner detector layers and building the polarized atomic beam target.

[1] Supported by BMFT under contract 06 BN 664 I(6)

[2] The EDDA collaboration:

J. Bisplinghoff^a, R. Bollmann^b, P. Cloth^c, O. Diehl^a, F. Dohrmann^b, V. Drücke^c, S. Eisenhardt^b, H.P. Engelhardt^a, J. Ernst^a, P.D. Eversheim^a, D. Filges^c, M. Gasthuber^b, R. Gebel^a, J. Greife^b, A. Groß^b, R. Gross-Hardt^a, S. Heider^a, A. Heine^a, F. Hinterberger^a, R. Jahn^a, M. Jeske^a, U. Lahr^a, R. Langkau^b, R. Maier^c, R. Maschuw^a, T. Mayer-Kuckuk^a, G. Mertler^a, B. Metsch^c, F. Mosel^a, M. Müller^b, M. Münstermann^b, H. Petry^c, D. Prasuhn^c, M. Radtke^b, H. Rohdte^a, D. Rosendaal^a, U. Roß^a, P. von Rossen^c, H. Scheid^a, H. Paetz gen. Schieck^d, N. Schirm^b, M. Schulz-Rojahn^a, F. Schwandt^a, W. Sobel^b, G. Sterzenbach^c, H.J. Trelle^a, A. Wellinghausen^b, W. Wiedmann^a, K. Woller^b, R. Ziegler^a

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A laser based monitoring system for the EDDA experiment

S. Eisenhardt¹, J. Greiff¹, M. Müller¹ and M. Radtke¹
for the EDDA Collaboration

The outer layers of the EDDA detector represent a total of 176 channels with photo-multiplier (PM) readout. This amount calls for a computer based, versatile PM gain monitoring and control system that can be used during the cyclic operation of the internal target when EDDA is taking data as well as in between and in preparation for beam times at COSY. The high quality of the COSY beam

schematically in Fig. 1. A N_2 -laser (VSL-337 from LSI Laser Scientific) is triggered with typically 16 Hz. The UV pulses ($\lambda = 337$ nm) are converted into visible light ($\lambda \approx 425$ nm) either by means of plastic scintillator rods or with a dye laser module. A small fraction of the N_2 -pulses goes into the " N_2 -monitor" of the UV intensity (photosensor H5783 of HAMAMATSU) that also triggers the data acquisition.

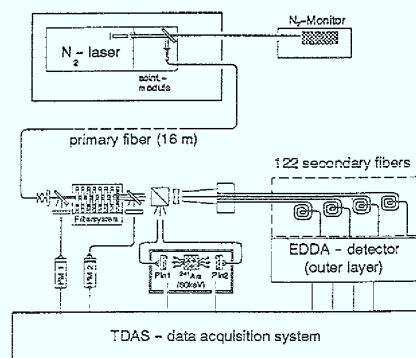


Fig. 1: Schematic layout of the laser based monitoring system.

and the count rates resulting from its interaction with the internal target exclude radiation damage of the scintillator components. Therefore it was decided not to include the scintillation process in the detector survey and to operate with light pulses that correspond in duration and spectral shape to scintillation pulses.

The final layout of the system [1] is shown

The 425nm-light pulses are then transported with a primary fiber to a reference and distribution box below the EDDA detector. In the box, the light passes through a system of filters and several beam splitters before it homogeneously illuminates the bundle of secondary fibers connected, near their PMs, to scintillator elements of the EDDA detector. The altogether seven different grey filters can be independently moved by CAMAC control and allow to adjust the light intensity quantitatively within a wide range.

There are several submonitors that control the correct operation of the monitor system. The fractions of primary light illuminating PM1 and PM2 allow to survey the intensity of the light source and the attenuation due to the filter system, and, in combination with the " N_2 -monitor", the conversion of the UV to visible light. The light pulses applied to the pin diodes Pin 1 and Pin 2 facilitate a pulse height calibration in terms of the photopeak of a ^{241}Am source.

The homogeneity of the light distribution is controlled by means of two fibers feeding

secondary light to PM1 and PM2 with an optical delay of 150 ns in order to allow a discrimination against the preceding signals. Several of these components have been adopted from the light pulsing system of TAPS [2].

So far, the following tasks have been successfully performed:

- (i) Test of the linear PM response by means of different filter settings (Fig.2)

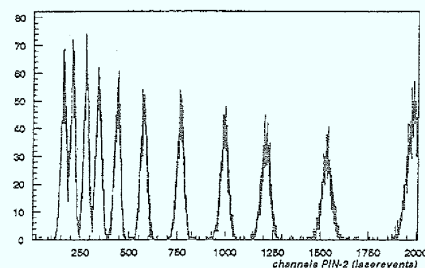


Fig. 2: Pulse height spectra obtained with different filter settings from a reference PIN-diode.

- (ii) Online control of the electronics setup before data taking. As an example shows Fig. 3 an ADC-malfunction.

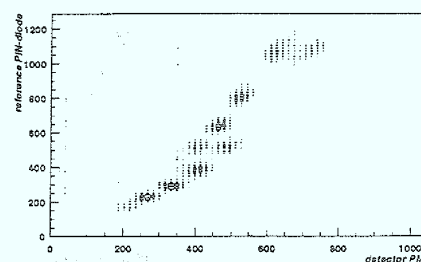


Fig. 3: Correlation plot of laser monitor pulse heights read from Pin 2 and one of the EDDA scintillators showing an irregularity in the ADC of the latter.

- (iii) Monitoring the gain of the PMs in the EDDA detector as a function of time

by division of the integrated charges of their anode signals by the corresponding PIN diode signals on an event-by-event basis. Fig. 4 shows gain drifts with respect to gain values obtained 24h earlier.

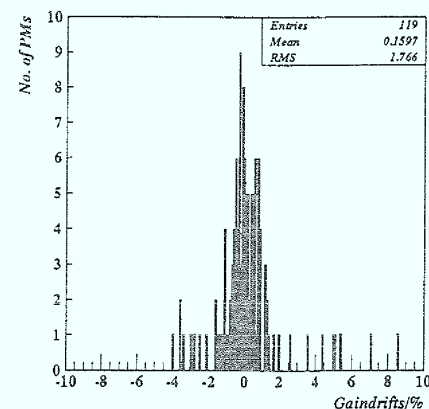


Fig. 4: Long term gain control. Relative gain-drifts of PMs measured after 24h of operation.

- (iv) Detection of short term gain drifts in cyclic operation of the EDDA detector, when data is taken with systematically varying count rates (measurement during ramping, influence of beam life time).

Additional applications like automatic online-control-functions are under development.

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Scintillating fiber elements for the outer layer of the EDDA detector

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for the EDDA Collaboration

The scintillator semirings comprising the outer layer of the EDDA detector (cf. Fig. 1 of the contribution J. Bisplinghoff et al. to this report) have to serve two purposes, namely (i) provide fast triggering on elastic pp scattering and (ii) allow a precise determination of the scattering angle Θ_{lab} . In order to comply with (i), the semirings are given a width $\Delta\Theta_{cm} \approx 5^\circ$. For (ii) the semirings should be designed in width such that each proton originating from the target traverses two adjacent rings; the point of incidence may then be reconstructed offline from the ratio of their light yields. This is the "light splitting scheme" [1].

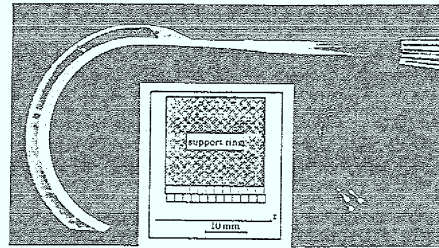


Fig. 1: Semiring section composed of 24 scintillating fibers bent with $R = 170$ mm and arranged as shown in the inset. Also seen are the three connections to the light pulser system.

Both requirements are compatible and lead to rings that are the smaller the higher the scattering angle Θ_{lab} is. For $\Theta_{lab} \geq 52^\circ$ the resulting width of a semiring drops below 14 mm; this causes tremendous problems in manufacturing, mounting and performance [2,3]. Therefore we replaced the first 9 semirings for the interval $52^\circ \leq \Theta_{lab} \leq 72^\circ$ by two layers of each 36 scintillating fibers BCF-12 of quadratic

cross section 2×2 mm. The individual fiber covers an average interval $\Delta\Theta_{lab} \approx 0.5^\circ$ and 2×12 of them are combined in one assembly corresponding to 3 semirings.

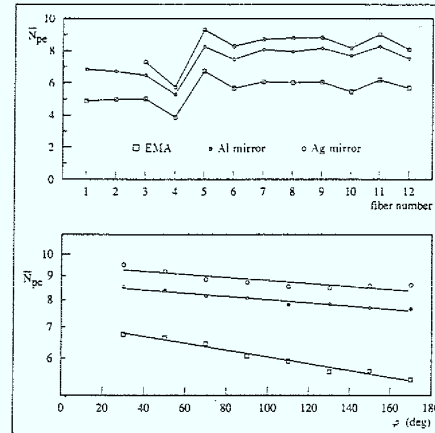


Fig. 2: Top: Average number $\bar{N}_{pe}$ of photoelectrons vs. fiber number for different designs of the far fiber ends. Bottom: Attenuation of light along a fiber; $\Delta\phi = 90^\circ$ corresponds to $\Delta s = 27$ cm.

This assembly is shown in Fig. 1. One can see, that the upper layer is shifted backward (by 0.4 mm) with respect to the lower one in order to avoid cracks allowing protons to escape after crossing less than 2 mm of scintillator material. The fibers are mounted into a supporting plexiglas ring and optically decoupled by a coating of extramural absorber (EMA). The two fibers on top of each other are combined for read out into one pixel of the 16 anode photomultiplier R 4760. Twelve of the

pixels are operated this way, among them are three, that are also connected to fibers of the laser monitoring system (cf. contribution S. Eischenhardt et al.). Before being installed in the EDDA experiment, the assemblies have been tested thoroughly with minimum ionizing electrons from a ^{90}Sr source, because the fraction of elastic scattering events involving the fiber-based semirings ranges from 75% (at $T_p \approx 800$ MeV) to 78% (at $T_p \approx 2205$ MeV) [4].

Fig. 2 shows for one semiring the average number $\bar{N}_{pe}$ of photoelectrons being released by the energy $\Delta E_e \approx 360$ keV deposited in a fiber. The results obtained with fibers with EMA coating at the end were considered insufficient because a minimum value $\bar{N}_{pe} \approx 6$ is estimated to be the lower limit for safe detection. The light yield was improved by reflecting the light at the far fiber end with a mirror made of aluminium foil of $R \approx 90\%$ reflectivity, and even more effectively, with a hard coated silver mirror ($R > 99.8\%$). At the same time the response along the fiber was improved such that the light loss from the far ($\varphi = 180^\circ$) to the near ($\varphi = 0^\circ$) end was reduced to typically less than 15% (Fig. 2b).

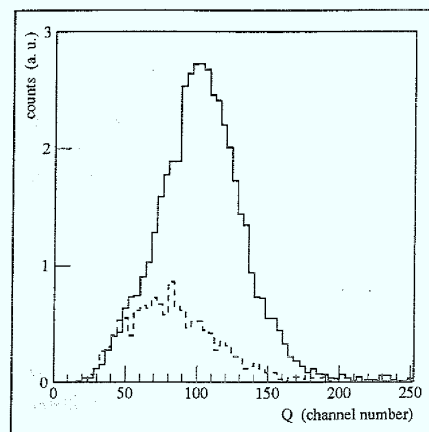


Fig. 3: Pulse height spectra obtained from a fiber homogeneously irradiated with 220 and 450 MeV protons.

After installation into the EDDA experiment the fibers were further tested with energetic protons. Fig. 3 shows the response of a fi-

ber pair positioned at $\Theta_{lab} = 53.5^\circ$ to $E_p = 220$ MeV ($E_p = 450$ MeV) protons from scattering of $T_p = 793$ MeV ($T_p = 2205$ MeV) projectiles depositing $\Delta E_p = 1.0$ MeV ($\Delta E_p = 0.75$ MeV). It can be seen that the fibers produce signals above any low energy cut off and are well suited for high energy proton detection. Meanwhile the fiberring section is routinely used in EDDA measurements, and Fig. 4 shows as an example the pp angular distribution (raw data, carbon background from the CH_2 target subtracted) normalized to data from [5]. The data was taken in November, 1994, and their analysis is in progress.

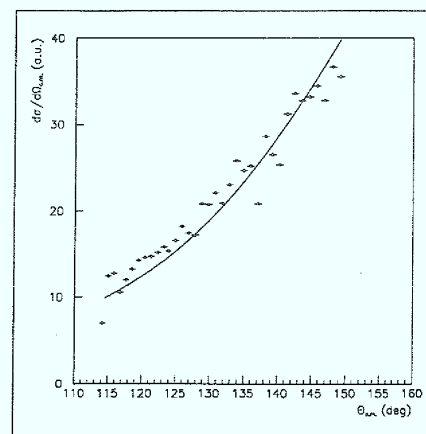


Fig. 4: Angular distribution of elastic pp scattering at $T_p = 793$ MeV taken with the fiber semirings.

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Monte-Carlo Simulations of the EDDA-Detector

J. Ernst, R. Groß-Hardt, U. Roß
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In the EDDA experiment, described in the previous contribution, some questions arise that may only be answered using the help of Monte Carlo (MC) simulations. MC techniques are especially useful in

- determining the acceptance of the detector,
- testing event reconstruction software,
- estimating the origin and importance of background reactions,
- helping to understand the shape of δ -electron spectra which are used for monitoring beam intensity and target thickness.

There exist several MC codes with which the simulation of EDDA should be possible in principle. However, in the near future the inner part of the EDDA detector will be added consisting of four layers of helically wound fibers. This geometry is very difficult to handle with certain MC codes – e.g. the GEANT code used in prior simulations of EDDA. Further, all MC codes known to us are written for use in high energy physics. Hence, we were looking for a more flexible code that allowed the implementation of medium energy physics and the helix geometry respectively. We found the C++-class-library GISMO [1] from SLAC to be the best choice. Its inherent advantages result from the thorough use of object oriented programming techniques. In less than a year several tools and options were developed, especially target event generators matching different requirements:

- **TwoParticles:** Two particles are generated in the CM system and transformed into the lab system. It was applied to simulate elastic pp scattering with cross sections following Simon et

al.[2], and for δ -electron production according to Rosenbluth's formula [3].

- **MICRES:** This event generator models the non-elastic hadronic interactions of mesons (π^0 , $\pi^\pm$, and $K^\pm$), nucleons and anti-nucleons up to $E_{kin} \leq 4.5 \text{ GeV}$ with nuclei $A \geq 4$. It is based on Bertini's intra-nuclear-cascade model (INC code of [4]) but describes the hadron nucleon collisions inside the nucleus as formation and decay of hadronic resonances according to the HADRIN code [5]. It also allows for evaporation and fission of the residual nucleus [4]. MICRES was developed by D. Theis [6] and was implemented into GISMO by U. Roß [7] for both the use as target event generator and for simulating hadronic reactions during the tracking phase. It is also available as a stand-alone version [8].
- **FOWL:** This well known MC code provides the 4-momenta of N particles equally distributed in phase space.
- **GunGenerator:** This generator produces single particles uniformly distributed in a user specified range of solid angle and according to one of three different energy distributions. This allows for simulation of radioactive β - and γ -sources.

Electromagnetic and hadronic reactions during the tracking of particles through different media are performed by EGS4 and MICRES respectively. For other hadrons than those mentioned above only energy loss and angular straggling is calculated.

In order to facilitate the use of GISMO for non-experts an interfacing utility was developed allowing specific calculations by fixing a couple of parameters in an initialising

file. The simulation provides the type of event (multiplicity, type of particles), the energy loss in the hit detector elements and the mean position where the hit occurred. Each MC event can be stored in its internal format or (planned) in a history tape format. By a converter programme the energy loss information is translated into ADC channels and transformed into the same format as used in the experiment (TDAS format). In addition, the hit pattern is specified in such a way that the experimental trigger conditions can be applied (coplanar and kinematical trigger).

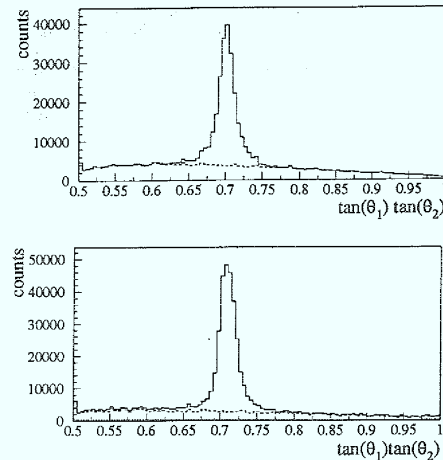


Figure 1: Comparison of the experimental (upper) to simulated (lower) distribution of coplanar events from elastic pp -scattering on a ultra-thin CH_2 target ($4 \mu m$) with a thin aluminium coating.

Fig. 1 compares the experimental distribution of coplanar events from elastic pp -scattering on an ultra-thin CH_2 target ($4 \mu m$) with a thin aluminium coating. The versatile use of GISMO is demonstrated also in Fig. 2 where a simulated δ -electron spectrum rather well reproduces the experimental one measured with PIN diodes which were inserted in the Al beam pipe at 45° to the target. Additionally, by a light ray tracing programme three-dimensional views of the EDDA-detector could be produced in colour.

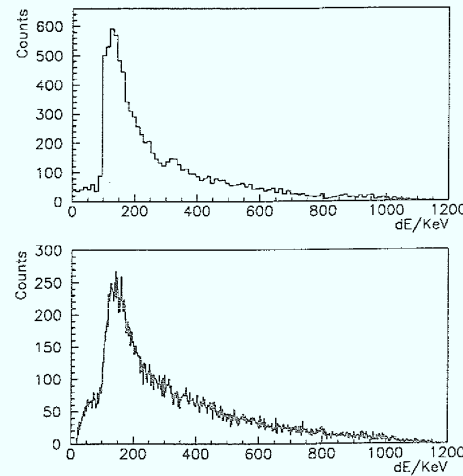


Figure 2: Simulated (upper) δ -electron spectrum as compared to the experimental (lower) one measured with PIN diodes at 45° to the target.

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Development of Event Reconstruction Programs for the EDDA Experiment
based on Monte Carlo Simulation of the Detector Response

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Introduction

After a successful period of construction and testing of the EDDA detector in the COSY ring, development of data analysis became necessary providing full scale resolution of the proton-proton elastic excitation function by optimization of the event reconstruction, rather than utilization of the EDDA granularity only, which ignores the full detector capability. The EDDA /1/ detector components consist of two layers measuring polar and azimuthal angles of the final event particles each. The outer layer comprises ring shaped scintillators to determine the polar angle and bar shaped ones for the azimuth. The scintillators are overlapping (from the view of the target) and thus, provide solid angle resolution higher than due to their size only, i.e. from the relationship of the overlapping signals information may be deduced to enhance the EDDA resolution power by an estimated factor of 5. This, however, requires evaluation of this somewhat complex relationship. The inner layer of EDDA consists of layers of helically wound scintillating fibers, two right and two left wound, providing combined measurement of the polar and azimuthal angles, and are also following the overlapping principle.

According to construction and experimental schedules of EDDA, development of the necessary off-line software was started for experiments with the fiber target and the outer layer of EDDA alone, the configuration in which data have already been taken for spin averaged cross-sections.

Development of sophisticated reconstruction software requires a detector signal simulation (Monte Carlo) system of adequate capability. The Jülich MC4 /2/ Monte Carlo program, especially designed for detector development, and to incorporate event reconstruction algorithms, is used for this purpose.

One of the basic features of this Monte Carlo system is the possibility of precise geometrical description of the detector. The BSG /3/ option for geometrical description in MC4 is used. Although the inner layer of EDDA is not considered by the time being in the reconstruction program, it is already fully available in the geometry part of MC4. The helix fiber layers are treated by a special program section dedicated to EDDA solely. Figures 1 and 2 show cross sections of the BSG geometry in its present status.

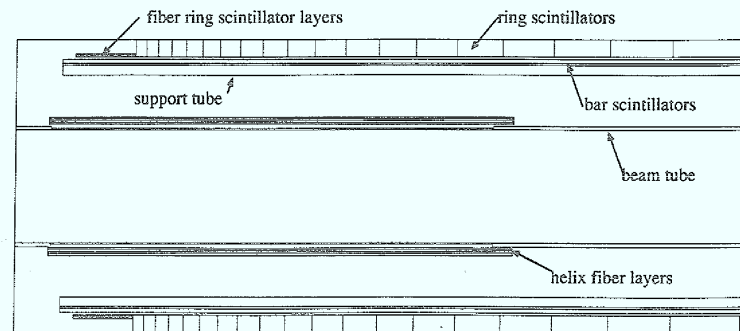


Fig.1 EDDA BSG z,x cross section

Short Description of the Reconstruction Method

For elastic proton scattering at known beam momentum the CM scattering angle θ_{cm} is the only free parameter. All other quantities describing the final state of the collision can be written as functions of θ_{cm} .

During start of run elastic events are produced in a loop over discrete z values, where z is the z coordinate of a final state particle entering the shell of polar detector elements.

These final state particles are fed into the tracking algorithm of MC4. Whenever they cause a signal in an active detector layer, the visible energy, the time of impact, and the detector element identification are recorded together with the description of the final state particles. The locations of the first and last event pattern involving a detector element are stored in a description area assigned to that element.

Running the same procedure with a sample of events per loop cycle with all straggling and secondary collision modules etc. turned on, results in a set of prestored mean

patterns for a given z value and typical errors for the detector elements. Since these estimates can only be as good as the assumed error functions, they have been neglected so far.

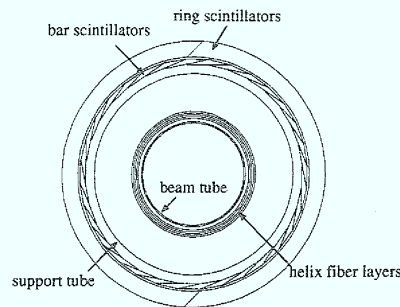


Fig.2 EDDA BSG -x,y cross section

After start of run, MC4 produces primary events using INC (Bertini intranuclear cascade model) or PSM (phase space model), then tracks through the geometry applying slow down (with straggling), small angle scattering and secondary nuclear events. All of this features can be controlled by user input.

Passing through the active detector elements, particles deposit energy which is converted into visible energy. This visible energy is then converted into detector signals taking into account light attenuation and time dependence in a semi-empirical way.

The first steps of reconstruction are multiplicity checks. For the azimuthal detector layer 4 elements have to trigger, grouped in 2 clusters of consecutive elements. The clusters numbered 16 and 32 are neglected, because they are not uniquely associated with the polar structures.

Up to 3 consecutive elements contribute to clusters in HRU and HRD, the two half ring groups. Only one cluster may be found per detector module. If HRU or HRD are not triggered at all, the half ring fiber layers take their role. In these detector modules a cluster is formed by up to 4 elements grouped in 2 layers with fibers in neighbourhood.

Multiplicity checks are accompanied by filling histograms with the element and cluster trigger rates.

The next reconstruction step is testing of the kinematics.

Coplanarity is checked based on element numbers triggered in the azimuthal layer. The kinematical trigger is checked after reconstruction of both θ values.

Estimations of the visible-energy to signal conversion functions are used in the reconstruction.

The two azimuthal clusters are used to estimate the ϕ angles of the final state particles. Therefore, the mean ϕ 's of the matching detector elements are averaged, weighted with their signals. The values of ϕ are then used to compute the distance of the point of impact to the photomultiplier tube of the polar ring structures. Attenuation of light can then be estimated.

Finally the most likely prestored pattern has to be selected. For each polar cluster in HRU and HRFU pointers had been stored into the list of elastic event patterns, recorded at start of run. In a loop over this list each pattern is compared to the current measure. Therefore, all signals in the detector modules have been added (squared and weighted with their typical errors). For each hit in the prestored pattern, the difference to the measure (again squared and weighted) replaces the contribution coming from the measure alone. This procedure assures that all hits in the measure and all hits in the prestored pattern are taken into account with their correct weights. The pattern with the smallest sum of errors is taken to be the best guess in reconstruction.

The results of the reconstruction are histograms including the trigger rates of detector elements and clusters. A n-tupel relates the Monte Carlo record of the event to these results.

As examples Figure 3 shows the distribution of the scattering angle in the CM system. They correspond to three separate runs producing clean samples of elastic, single π production and double π production events at a beam energy of 2.5 GeV. The concurrent events are suppressed fairly good. The acceptance of the elastic sample is as expected.

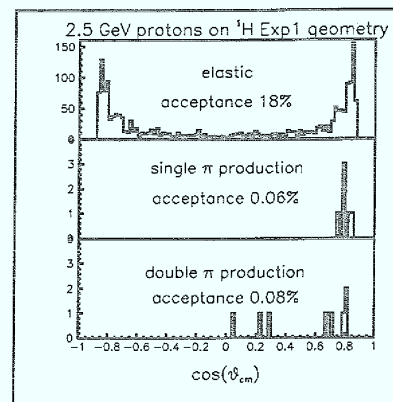


Figure 3

Figure 4 shows corresponding data for a 1 GeV beam. Here double π production does not occur.

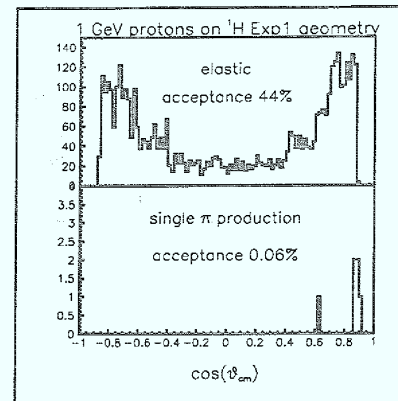


Figure 4

References:

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First Results from MOMO

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The MOMO experiment focusses on near threshold meson production via the reactions $pd \rightarrow {}^3\text{He} \pi^+ \pi^-$ and $pd \rightarrow {}^3\text{He} K^+ K^-$. It takes advantage of the high quality of the cooled external COSY beam and the existing spectrometer BIG KARL. The setup consists of a high granularity scintillating fibers meson detector near the target with a ± 45 deg. opening angle, and the spectrometer, which is used for ${}^3\text{He}$ -identification. The large solid angle and high resolution of this detection method will yield precision data on the low energy ($T < 50 \text{ MeV}$) meson-meson interaction and probe into questions like meson-nucleon resonances and KK-molecule.

The MOMO vertex dector consists of 672 scintillating fibers (round, 2.5mm diameter) arranged in three planes tilted 60 deg. versus each other. Each plane is subdivided into two identical modules. The fibers are read out by 16-fold photomultipliers (Hamamatsu R4760) followed by individually adjustable level adapters for each channel, which match the signals for the utilized PCOS III data aquisition system. All six modules are installed on location and the detector is fully operational. The efficiency was measured to be better than 99% for minimal ionising particles.

The MOMO scattering chamber houses the 4mm LD2 target as well as a remotely steerable ladder for beam viewers and solid targets. The 5 mm thin Al front end of the chamber faces the vertex detector on the outside and keeps straggeling of the mesons at a small level.

The BIG KARL spectrometer is also fully operational. Its status is described in detail elsewhere in this annual report.

Several test runs were successfully performed during the year including calibration of the setup via the reaction $pd \rightarrow {}^3\text{He} \pi^0$. The quality of the extracted COSY - beam increased constantly over the year. In two recent runs first measurements of two pion production via the reaction $pd \rightarrow {}^3\text{He} \pi^+ \pi^-$ were performed. The intensity of the 1150 MeV/c proton beam was some 10^7 /s with an almost 50% duty cycle (~6 s every 13 s). Figure 1 shows the vertex wall hit pattern of the first two pion events observed at COSY. Some 400 events could be unambiguously identified, which is demonstrated in figure 2. It depicts the coplanarity of the double hits in respect to the total meson momentum axis, which is defined by the beam and the ${}^3\text{He}$ momenta. Data analysis is in progress.

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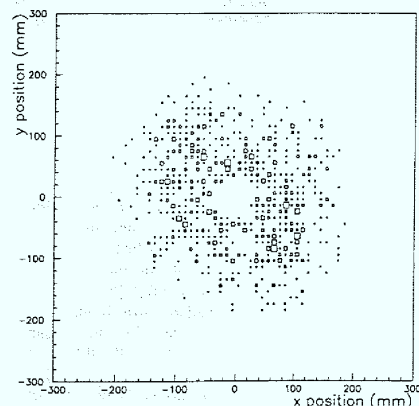


Fig. 1: Two - pion hits on the vertex wall

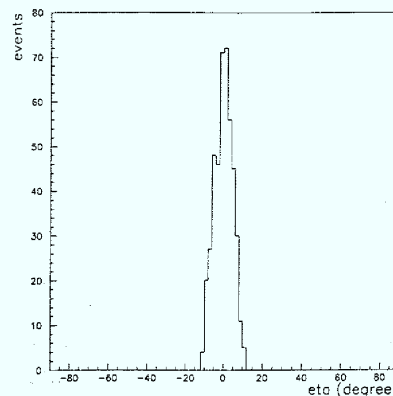


Fig. 2: Coplanarity of the $\pi^+ \pi^-$ - events

Status of GEM

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The previous year has seen a continuous progress of GEM. The hardware and software described in Ref. [1] has been further developed and has been tested during first runs with the COSY beam. In the first part we will discuss the status of the GERmanium wall. Then we will present the results achieved with the COSY beam employing only the magnetic spectrometer BIG KARL and its new detection devices.



Figure 1: The „quirl“ crystal. During operation the crystal is cooled through its mounting with liquid nitrogen. Note the read-out on printed Kapton foils, the structure on the surface and the hole in the center.

In its final form the Germanium wall will consist of four Germanium detectors: a “quirl” type tracking detector followed by three calorimeter detectors. The quirl has 200 Archimedes’ spirals which are arranged in an anticlockwise and clockwise direction on the front and rear side, respectively. The calorimeter detectors will each have 32 wedges on their front side. The first quirl built showed only a modest energy resolution, not sufficient for the required resolution to study pionic atoms. The reason for this was too shallow groves between neighbouring spirals. Alternative parameters in the implantation process were obtained leading to good separation between two adjacent spirals with a groove depth of 2 to 3 μm . This leads to smaller leakage currents and thus to an improved energy

resolution. New fabricated quirl detectors have been thoroughly tested and it was found that all spirals produce a signal output. Their energy resolution is typically 4 keV at 1332 keV energy. In Fig. 1 a quirl detector is shown with its read-out ribbons made of Kapton.

Further progress includes also the production of the last calorimeter crystal. As it has to cover the same solid angle as the ones closer to the target it has to be rather large: the

crystal has a 51 mm radius and a thickness of 15.5 mm. It was successfully transformed into a planar detector with an annular hole by adding 32 wedges to one surface. The typical energy resolution achieved was better than 5 keV for 1332 keV γ 's. Further details of the Germanium detectors and the cryostat are given in Ref. [2]. The two detectors presently available already allow some meson production study experiments close to threshold. However, for the pionic atom experiment the full detector will be required as is discussed in Ref. [3]. In order to perform such measurements a dedicated holder for such a subsystem in the cryostat was designed and built. In order to process high counting rates resulting mainly from elastic scattering, the linear shaping amplifiers were

modified to 10000 counts/second yielding 320000 counts/second for the calorimeter.

The performance of the detector set up in the experimental area was studied by mounting the cryostat with a single built-in Germanium detector in front of the first quadrupole of the magnetic spectrometer. Studies of the response of the built-in detector showed a satisfactory performance which deteriorated as soon as the quadrupole magnets of BIG KARL were switched on. This serious effect is being investigated.

The modified magnetic spectrometer together with its new focal plane detector was successfully tested with the COSY beam. The key question to be answered was whether one can extract the emission angles (polar and azimuthal angle) of reaction products out of the target from tracks measured in the focal plane of the spectrometer at a distance of 15.6 m away from the target position. Therefore the reaction $p+p \rightarrow \pi^+ + d$ was chosen: at a beam momentum close to threshold. The beam momentum of 800 MeV/c is only 11.15 MeV/c above threshold. This fact dominates the underlying kinematics shown in Fig. 2. The kinematical loci of the emerging reaction products are shown in a polar diagram as function of the momentum parallel to the beam axis and to the one perpendicular to it. The longitudinal momentum of the deuterons are approximately twenty times larger than the transverse momenta. This yields to very small emission angles in the laboratory system. Therefore, one needs very high precision in order to disentangle the centre of mass angular distribution. The complete laboratory distribution of the deuterons just fits completely in the acceptance of the magnetic spectrometer BIG KARL. The momentum of the beam protons is only slightly larger than the momenta of the deuterons, as can be seen from Fig. 2. This leads to a background problem. Nevertheless, this background can be suppressed completely. First the deuterons were measured in coincidence with pions. The latter were measured closely behind the target with an annular scintillator ring.

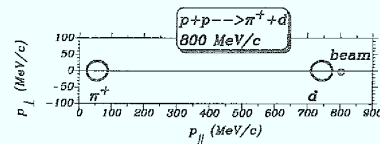


Figure 2: Kinematical loci of π^+ and deuterons from the present reaction as function of their momentum parallel to the beam direction ($p_{\parallel}$) and perpendicular to it ($p_{\perp}$). Also indicated in the figure is the beam momentum.

This detector was also used as a stop detector for a time-of-flight (TOF) measurement of the deuterons. This flight time is twice as long as the time of the protons in the same momentum acceptance of the spectrometer. In the focal plane the tracks were measured with two stacks of multiwire drift chambers (Ref. [1]). The start signal for the drift-time measurement came from a coincidence between scintillators in three layers of the trigger hodoscope (see Fig. 3).

Protons were suppressed against deuterons puls height discrimination by the first layer. A further rejection of protons was achieved by using the fourth hodoscope layer

as a veto detector located behind aluminium absorbers thick enough to stop the deuterons.

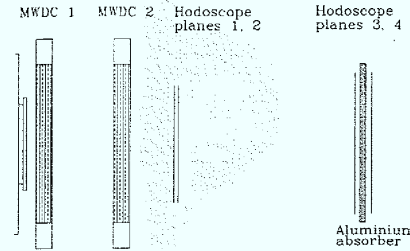


Figure 3: Schematic diagram representing the detectors in the focal plane of the magnetic spectrometer. Two stacks of multiwire drift chambers to allow for x- and y-measurements are followed by a trigger hodoscope consisting of four layers of scintillator. The first three layers operate in a threefold coincidence while the fourth layer behind the aluminium absorber is used as a veto counter.

In Fig. 4 the TOF spectra are shown for a measurement resulting from a full and from an empty target cell. The target cell (Ref. [4]) was 4 mm thick and 5 mm in diameter. Its windows were made out of 1.5 μ m thick mylar. Beam halo events were suppressed by a scintillator ring with a 3 mm diameter hole in front of the target. The beam diameter was found to be smaller than 1 mm.

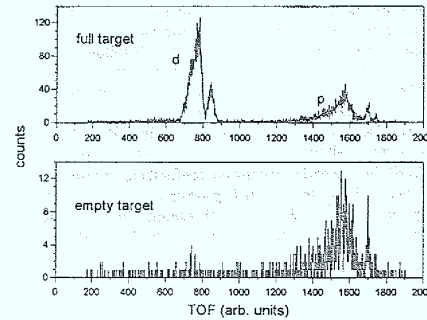


Figure 4: TOF spectra as obtained from the second layer of the hodoscope (start) and a delayed stop from the pion ring-detector. The spectrum including deuterons is from the full target cell measurement while the other one was taken with an empty target cell.

In Fig. 5 the reconstructed kinematical loci of the detected deuterons are shown in a polar diagram. For this reconstruction only theoretical transport coefficients were used. Figure 5a shows the raw data with constraints as explained. An additional geometrical cut was introduced which excluded all those events which lie outside the horizontal acceptance of the last vacuum chamber of the spectrometer. When an additional gate is set on the deuterons in the TOF spectrum (Fig. 4) the clean data sample shown in Fig. 5b is obtained. The measured data can be projected to the parallel momentum axis. This is

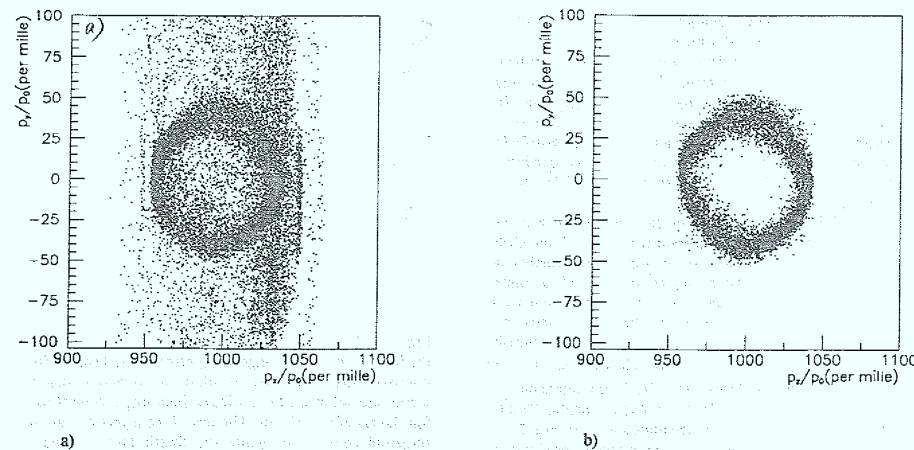


Figure 5: Polar diagram of the measured coincident events as function of the longitudinal and transverse momenta. Figure a) shows the on-line raw data. Fig. b) shows data with the off-line requirement of being deuterons in the TOF spectrum and applying a geometrical constraint on the trajectories.

shown in Fig. 6 which is linearly related to the centre of mass cosine. Also shown is a preliminary Monte Carlo calculation employing the code GEANT. It follows almost the data thus indicating the dominance of the s-wave strength in the present reaction. A very preliminary analysis of the corresponding angular distribution in the centre-of-mass system for a subset of the data yield an anisotropy as shown in Fig. 7. This quantity a_2/a_0 is the ratio of the Legendre expansion coefficients. Because this number is finite, there is still p-wave contribution. These data are much closer to threshold than previous measurements (see Fig. 7), made possible by the high quality COSY beam. The result is compared in Fig. 7 to those obtained from earlier measurements (Ref. [5]) and to a theoretical prediction [6].

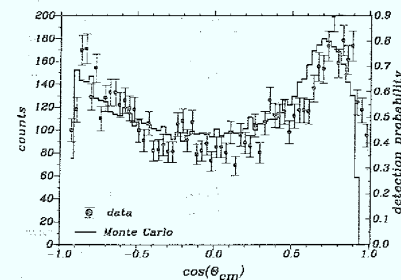


Figure 6: Projection of the data from Fig. 5b onto the longitudinal momentum axis. The data are shown as dots with error bars (left ordinate). Also shown is a Monte Carlo calculation taking target thickness finite beam diameter and acceptance cuts into account (histogram, right ordinate). The abscissa is transformed into the emission angle of the centre-of-mass system.

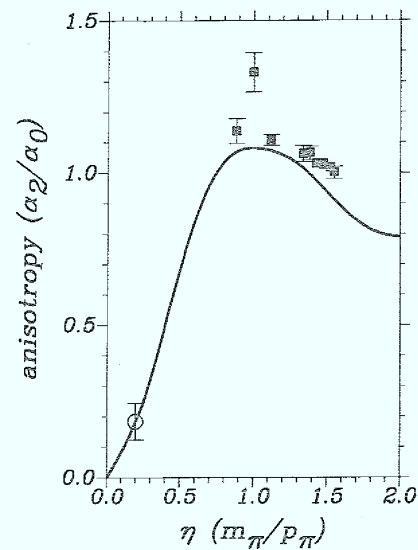


Figure 7: Anisotropy deduced from a subset of the data is compared with data taken from Ref. [4] as function of the pion momentum ($\eta = p_\pi/m_\pi$). Also shown is a calculation performed by Niskanen (Ref. [6]).

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Status of the COSY-11 Installation

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COSY-11 is an internal experimental setup (see Fig. 1) for threshold measurements of the production of mesons or mesonic pairs [1,2] which is presently installed in the COSY-ring. The experiments will be performed using a cluster target in front of a COSY dipole acting as a magnetic separator for the ejectiles, which will be registered with several tracking and timing detectors.

elements the minimum ionizing peak was well separated from the noise.

The data acquisition system for COSY-11 is working. It was used for the cosmic measurements

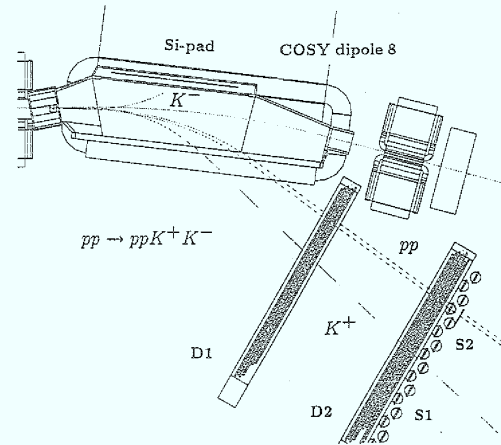


Figure 1: Sketch of the COSY-11 installation at a COSY dipole. For the positively charged particles drift chambers D1, D2 and scintillators S1, S2 and S3 are used. S3 is positioned 10 m apart from S2 and not shown in the figure. Si-pad detectors combined with a long scintillator bar was mounted in the dipole gap for the detection of negatively charged particles.

The standard COSY vacuum chamber in dipole no. 8 was replaced by a special chamber with large windows for an undisturbed transport of the reaction products to the detectors. The main window is covered by a 300 µm thick foil of carbon fibres with a 30 µm Al-layer on the inner side to guarantee UHV conditions in the ring.

The cluster target and all detectors are ready for the installation which will take place end of February 1995.

The target was mounted outside of the ring and extensive tests have been made which have demonstrated the stable operation of the system. All detectors have been tested and calibrated with cosmic radiation i.e. minimum ionizing particles. For these tests the drift chambers and the scintillation counters S1, S2 and S3 were positioned close together so that cosmes are able to pass through all units simultaneously. For the calibration of the scintillation detectors the particle tracks were used which were extracted from the drift chamber data. The function and energy resolution of the silicon pad detector was checked independently of the other systems with cosmes. In all Si-

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Status of the 0° Facility

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The preparations of the 0° Facility and its detector system for the study of the subthreshold K^+ production in proton-nucleon reactions and of the proton-induced deuteron break-up have been advanced considerably in 1994. The work focused mainly on the following aspects:

- Theoretical studies were concerned with the details of the production of fast deuterons in proton-nucleus collisions, be it in quasi-free processes, through coalescence in a late state of the reaction or in the first step of the two-step K^+ production [1]. The knowledge of the momentum spectra of those fast deuterons originating from the different processes is important for the study of the correlated K^+ -d production at the 0° Facility. Since K^+ mesons are produced together with associated Λ -hyperons, the probability of capture of hyperons with creation of hypernuclei was theoretically determined and cross sections were calculated [2].

Calculations on the subthreshold K^+ production have been continued [3]; they led to the conclusion that experimental studies of this interesting process should be possible at the 0° Facility.

Polarization observables for the deuteron break-up with polarized protons and deuterons have been calculated under the kinematic conditions of the 0° Facility [4]. All the observables manifest the $\Delta(1232)$ baryon excitation in the intermediate states as a dominant mechanism of the backward proton production with momenta close to the kinematic limit.

A new possibility for an experimental study of proton-nucleus interaction accompanied by high energy-momentum transfers to few nucleons of the target nucleus has been considered. It concerns the measurement of the two-particle differential cross sections for the emission of a fast proton and a light fragment at small forward angles. The Rossendorf collision model [5] was used to evaluate the cross

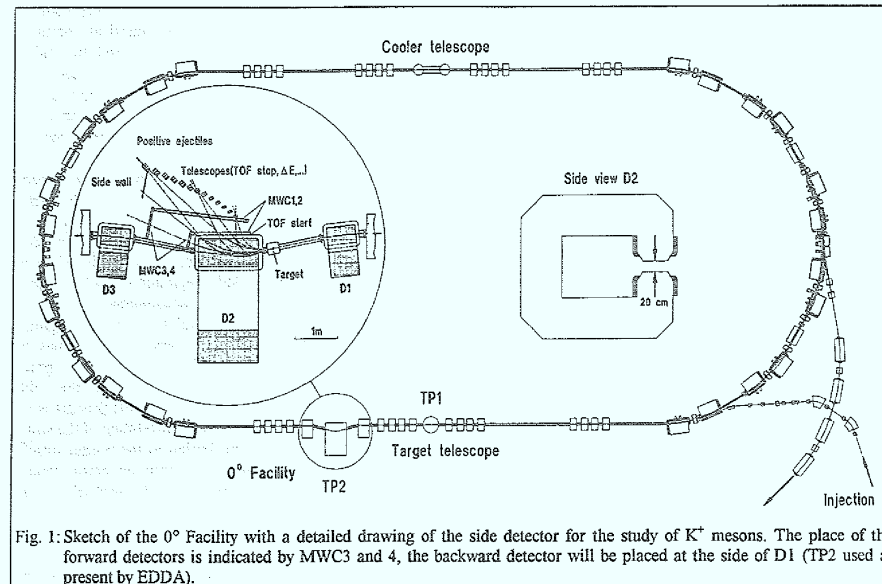


Fig. 1: Sketch of the 0° Facility with a detailed drawing of the side detector for the study of K^+ mesons. The place of the forward detectors is indicated by MWC3 and 4, the backward detector will be placed at the side of D1 (TP2 used at present by EDDA).

sections for the simplest process of that kind, pd-interaction resulting in forward emission of a fast deuteron and a proton [6].

The subthreshold K^+ meson production in proton-deuteron collisions has been studied with use of a light front formalism for description of the relativistic wave function of the deuteron. The First Collision and Two Step Models have been considered. Results of calculations of the dependence of the total cross section on the kinetic energy of the incident proton and momentum- and angular spectra of the K^+ meson show a behavior similar to that of proton-nucleus collisions.

- The design of the 0° Facility has been reviewed on March 18, 1994, by a committee of experts. The concept was approved in all its main aspects. It has been suggested to consider the target place TP1 in the centre of the straight section of COSY as a site for the installation of the facility. This alternative to the original concept of using TP2, see Fig. 1, is presently under study including both the 1π and 2π modes of operation of the telescopes.

- The realization of the facility has been pushed forward through the completion of the concept studies of the power supplies for the magnets. An available test device of COSY will be used for D1 and D3. For D2 a dedicated power supply has to be built. Calls for tenders have been sent out.

- The design of the vacuum chambers has been pursued. It is most important to work out a concept with minimal eddy currents since they reduce the dipole field during ramping and induce a sextupole-field component which would be disturbing for the circulating beam. Therefore the top and bottom walls have to be as thin as possible taking into consideration the necessary mechanical stability. Likewise, the side flanges which carry the window foils must be made in such a way that their frames are very thin without loss of the necessary sealing efficiency. This can be achieved through the use of the massive pole pieces of the magnets for picking up the forces resulting from a special chamber design. It is planned to use this technique for the chambers both in D1 and D2.

An additional advantage of thin-walled chambers is that the gap widths of the magnets can be used with more than 95 % efficiency. The remaining effect of eddy currents on the magnetic field strengths during ramping (with $\Delta B/\Delta t \leq 1$ T/s) will be compensated by a conveniently chosen current ramp of the power supplies of the 0° Facility. Here it should be noted that the 0° Facility does not affect the ramping rate when COSY serves other users since its magnets will then not be in operation.

- The time dependence of reaction rates, which are expected in the K^+ -experiment, was studied with the help of simulation calculations. For a 2 mm wide carbon foil of $100 \mu\text{g}/\text{cm}^2$, irradiated by $1 \cdot 10^{10}$ circulating protons of 1 GeV at the internal target place TP1 or TP2, we obtain a luminosity of 1 to $3 \cdot 10^{33}/\text{s} \cdot \text{cm}^2$, which can be made

constant during a period of ≈ 1 s or more, through proper displacement of the target during the irradiation [7].

- The development of the detector systems for experiments on the subthreshold K^+ production and the proton-induced deuteron break-up has been continued. In order to make sure that the detector system fulfills the specifications for the identification of the K^+ mesons in presence of an extremely intense background of protons and pions, test measurements of two telescopes have been performed at the K^+ beam of the ITEP, Moscow. These studies [8] showed that, indeed, a „background“ suppression factor of better than 10^5 can be achieved through a discrimination via time-of-flight (TOF) and energy loss. This is in good agreement with the results of simulation calculations [9].

Based on this finding the GEANT calculations have been extended in order to complete the detector concept for the measurement of the momentum spectra of K^+ mesons in the subthreshold production at nuclei. The construction of the detector telescopes has begun [10]. Special attention is paid on the shielding of the start detectors of the TOF setup against the fringe field of D2. Monitor systems for the scintillation counters are being developed. The construction of Cherenkov counters for the K^+ -detector telescopes is in progress. They will be used in these telescopes for the suppression of leptonic and pionic background.

The first wire chamber (close to D2 in Fig. 1) is under construction at the Forschungszentrum Rossendorf [11,12].

- The imaging surface of the magnet D2 has been determined in simulation calculations for the target located at 35 cm in front of the magnet iron. The transverse sizes of the stop telescopes were estimated taking into account the magnet geometrical aberrations in D2 and the particle deflection caused by multiple Coulomb scattering [13]. The transverse sizes of the modules do not depend on the particle momentum and would be recommended for all modules equal to 10 - 12 cm. This result agrees well with module sizes adopted in [10].

For the deuteron break-up experiment the expected missing-mass resolution has been calculated with protons to be detected in the forward and backward detectors (see Fig. 1) [14]. A full GEANT simulation with a proper event generator and track reconstruction procedure was carried out. An acceptable separation of the break-up process and the fragmentation accompanied by pion production has been found in the missing-mass spectrum.

The acceptance and missing-mass resolution has also been evaluated for detecting the process $p+p \rightarrow p+p+\phi \rightarrow p+p+K^+ + K^-$ with both the initial state protons polarized. The study of this process is very attractive due to the expected manifestation of the strange quark-antiquark pair contribution to the proton structure [15]. The calculation shows a reasonable performance of the facility for the proton pair in the forward detector and the positive meson registered in the side detector.

- The work on the preparation of the detectors at the JINR

Dubna for the deuteron break-up experiment comprises the design and beginning of production of the forward and backward detectors [16], of the development of a corresponding fast trigger system [17] and the readout electronics.

Cherenkov counters of total internal reflection have been proposed for π/K and p/d separation and background suppression at the trigger level in the d break-up studies. Cherenkov counter characteristics were studied by Monte-Carlo simulation using GEANT3 package for particle tracing [18]. The idea is to make use the total internal reflection for the Cherenkov light collection for particles with higher β . The advantages of the proposed Cherenkov counters are the possibility of operation at high Cherenkov radiation intensity (because of high refractive index) and the threshold tuning with rotation angle variation.

For confirmation of realization feasibility a relatively simple prototype of these Cherenkov counters was created and tested on the beam. Measurements were performed at the external beam of the ITEP accelerator with colleagues from ITEP and JINR. Results of measurements for protons and pions at beam momentum 1.5 GeV/c are presented in [19]. A dependence of the photomultiplier anode pulse height on the angle relative to the beam direction is clearly observed. During the counter beam testing it has been obvious that the counter characteristics can be improved by further developments.

• The work for the fast trigger electronic system, which allows to distinguish K^+ -mesons from protons and pions, was continued [17]. For the read-out of multi-wire chambers the development of a very compact electronic board, mounted directly on the chambers, has started. Most of the electronic functions, including data reduction, are executed by the large scale integrated chips RAL110/111, or possibly by new elements, which are actually studied at Rutherford Lab., and which will be specially adapted to the output signals of the wire chambers foreseen for experiments at the 0° Facility. Major components of the Fastbus-VME-DECstation data acquisition system [20] for the 0° Facility experiments were completed at the FZ-Rossendorf. The system was transferred to Jülich and hard- and software tests were performed. The digitizing Fastbus modules (QDCs and TDCs) fulfill the requirements of resolution and speed. The software developed so far works well for the standard hardware and allows to perform the standard operations, including elementary data analysis. For small experiments, in which detectors or experimental techniques are examined, we have developed a CAMAC-PC (connected via VIC-bus) data acquisition system. Up to $\approx 2 \cdot 10^5$ detector channels can be read per second. The software can be adapted to include any CAMAC-module. Programming of any desired measurement sequence or experiment control operation is relatively easy.

• The studies on a polarized gas target (storage cell) for the 0° Facility have been continued as well. It is hoped that the Filtext target [21] can be used later at the 0° Facility. But it is also checked whether equipment which is available from the INP Novosibirsk can be adapted to the use at the 0° Facility.

• First concept studies for a detection system for K^- mesons have begun [22].

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One of the main goals of the experimental program at the 0° Facility is the investigation of the subthreshold K^+ production in proton-nucleus interaction [1]. In order to get information about the reaction mechanism it is planned to measure the kaons in coincidence with other particles like protons and light fragments. In particular, from the measurement of correlated K^+ -deuteron pairs one hopes to unambiguously identify the two-step kaon production with the formation of an intermediate pion in the first step ($pN \rightarrow d\pi$, $\pi N \rightarrow K^+\Lambda$) [2,3]. Since the expected cross sections for such processes are extremely small it is planned, however, to first measure uncorrelated deuterons in order to calibrate the detection system.

Figure 1 shows the calculated momentum spectra of deuterons produced in $p+^{12}\text{C}$ collisions at $T_p=1$ GeV. The spectra are integrated over emission angles $\vartheta < 10^\circ$ which are accessible at the 0° Facility. The shaded area indicates those deuterons which are produced in coincidence with K^+ mesons via the two-step mechanism in the quasi-free process $p+N \rightarrow d+\pi$. It can be seen that they are peaked around $P_d \approx 800$ MeV/c and, thus, can be measured at the 0° Facility [1,5].

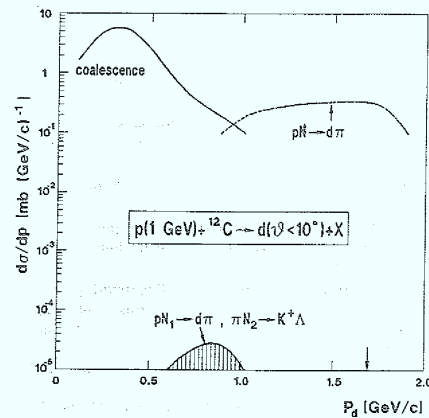


Fig. 1: Calculated momentum spectra of deuterons for $p+^{12}\text{C}$ collisions at $T_p=1$ GeV. The shaded area shows the spectrum of deuterons produced in coincidence with K^+ mesons [2,3]. The upper curves show the deuterons produced in the quasi-free mechanism $p+N \rightarrow d+\pi$ and via the coalescence mechanism [4].

Since at $T_p=1$ GeV one is far below the nucleon-nucleon threshold for the production of K^+ mesons ($T_p=1.58$ GeV) the involved cross section for the production of K^+d pairs is much smaller ($\sigma_{\text{tot}} \approx 10$ nb) than the total production cross section for uncorrelated deuterons. These deuterons can be produced via the coalescence mechanism [4] or via the quasi-free process $p+N \rightarrow d+\pi$ [4,6]. It can be seen from Fig. 1 that the quasi-free deuterons are peaked around 1.7 GeV/c. This shift to higher momenta relative to the deuterons produced in coincidence with K^+ mesons can be explained as follows: the kaons are produced in the second step of the two-step mechanism ($\pi+N \rightarrow K^++\Lambda$). Since one is far below the NN threshold only those pions which have the highest possible momenta can produce a kaon. The pions are produced together with the deuterons ($pN \rightarrow d\pi$) and, thus, only deuterons with very low momenta can be created in the case of kaon production.

The measurement of fast deuterons can also serve as a test of the coalescence model which is rather well established for low deuteron momenta ($P_d < 500$ MeV/c at $T_p=1$ GeV). At higher momenta there are still considerable uncertainties in the calculations since here high-momentum components of the deuteron wave function as well as kinematical and dynamical correlations have to be taken into account [4].

Summarizing, the 0° Facility is ideally suited to study the production of fast deuterons since a wide momentum range can be covered and measurements under forward angles ($\vartheta < 10^\circ$) are possible where the deuteron momentum spectrum is peaked [4].

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Formation Cross Sections and Recoil Momentum of Heavy Λ -Hypernuclei in the (p, K^+) -Reaction

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The production of hypernuclei is an interesting nuclear process in itself and, moreover, offers a source of information for heavy ion physics, astrophysics and cosmology. The possibility to produce hypernuclei in the (p, K^+) reaction has been first mentioned by Yamazaki [1], but Chrien [2] presumed that the cross section for this reaction might be very small. The hypernucleus formation consists of two steps as pointed out by Zofka [3] and Nifenecker [4]. The Λ hyperon production is the first and its attachment to the nucleus is the second step. The probability of both processes depends on many factors: The Λ production is determined by the beam energy, the momentum distribution of the nucleons in the nuclei, possible correlations of nucleons in clusters, etc.. The probability of the Λ attachment depends predominantly on the momentum distribution of the produced hyperons [5] and its modification due to Λ -N rescattering [6] but may be also influenced by the features of the nucleus to which the hyperon is attached [4].

The aim of this contribution is to present results of the calculation of the cross sections and of the recoil momentum for an experiment which is in preparation at the COSY accelerator in Jülich. The goal of the experiment is the investigation of the mechanism of the production of heavy hypernuclei in proton interactions with atomic nuclei.

It is planned to record fission fragments. This should allow a strong suppression of the very large background due to proton, neutron and pion emission. The recoil shadow method [7] will be used [8] for elimination of fragments originating in the target and arising from prompt fission of non-strange nuclei. This method assures that fission fragments can be observed in the shadow region in two cases. First, the Λ particle may attach to one of the fragments of the promptly fissioning hypernucleus. These fragments leave the target into all directions, but their detection in the shadow region is excluded by suitable diaphragms, unless the decay of the Λ particle kicks the fragment into the shadow region. Secondly, the heavy hypernucleus surviving prompt fission recoils out of the target, and the delayed Λ -decay induced fission occurring at a distance from the target downstream of the beam permits fission fragments to reach directly the shadow counter. Polikanov and his coworkers have successfully applied this method for obtaining information on the lifetime of very heavy hypernuclei in antiproton annihilation [9].

The phase-space folding model [10] was used to calculate the total cross sections of K^+ production in the sub-threshold region taking into account one step ($N+N \rightarrow K^++N+\Lambda$) and two-step processes ($N+N \rightarrow \pi+N+N$; $\pi+N \rightarrow K^++\Lambda$). The experimental results from Koptev et al. [11] and the model predictions are compared in Figure 1. As in both processes the kaons (K^+) are accompanied by Λ hyperons, it is natural to

assume that the folding model is also capable for a realistic estimation of the total cross section and momentum spectra of the Λ hyperon production. Therefore, the results of the folding model for the differential Λ -spectra were used as an input for the Boltzmann-Uehling-Uhlenbeck (BUU) calculations which enabled us to simulate the time development of the Λ -nucleon dynamics in the target nucleus. First, a hyperon may be attached to the nucleus without rescattering, when it has been produced with kinetic energy smaller than the attractive hyperon - nucleus potential (of order of 27 MeV [12]). In this case the resulting total cross sections for hypernucleus production in proton - uranium collisions, depicted in Figure 1, are rather small. The main contribution here comes from the two-step process due to two factors [13]: the cross section is larger for hyperon production (approximately one order of magnitude in the proton beam energy range from 1 to 1.5 GeV) and the hyperon momentum distribution is shifted to smaller momenta as compared to the one-step process (e.g. at a proton beam energy of 1 GeV the maximum of the Λ momentum distribution is close to 300 MeV/c for the two-step, and 700 MeV/c for the one-step process, respectively).

More involved calculations take into account the rescattering of the Λ hyperon in the nucleus. This leads to a slowing down of the hyperon and thus allows to capture even a Λ hyperon with high initial momentum by the mean field of the nucleus. The corresponding calculations performed within BUU demonstrate a very significant increase of the cross section for hypernucleus production (thick line in Figure 1) even by a factor of 30 in comparison to the model which ignores rescattering. The BUU calculations are capable also for an estimation of the distribution of e.g. momentum,

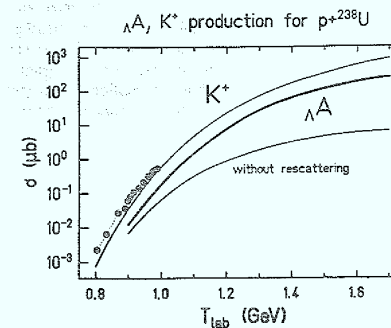


Fig. 1: Total production cross sections predicted for K^+ mesons in the $p+^{208}\text{Pb}$ reaction (upper curve) and for heavy Λ hypernuclei Λ_A in the $p+^{238}\text{U}$ reaction with and without rescattering. The dots are the experimental K^+ production cross section for $p+^{208}\text{Pb}$ [11].

excitation energy, angular momentum and mass of the produced hypernuclei. In Figure 2 these distributions are shown in histogram-like representation.

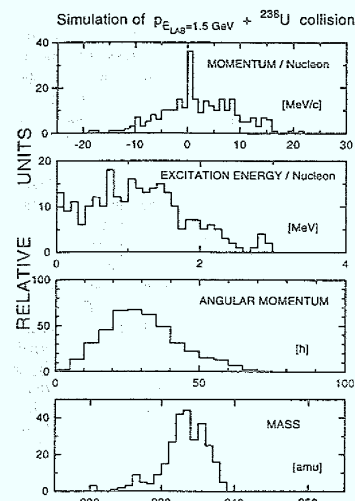


Fig. 2: Distributions of momentum, excitation energy, angular momentum and mass of Λ hypernuclei produced in the $^{238}\text{U}(p,K)$ reaction at 1.5 GeV.

It has been found that in the $p+^{238}\text{U}$ reaction at 1.5 GeV, the average mass A of the produced hyper-nuclei is 233 ± 3 . The angular momentum distribution first grows linearly and then drops to 0 at $\sim 70 \hbar$. The linear growth means that the sticking probability as a function of growing impact parameter varies rather slowly, until the impact parameter reaches a value comparable with the radius of the target nucleus. Then hypernuclei are no longer produced and the angular momentum distribution approaches zero. The momentum and excitation-energy distributions are rather broad. We expect a significant ratio of hypernuclei produced with low excitation energy (which will prevent them from undergoing prompt fission), and with high momentum (so that they will travel further away from the target before they undergo weak nonmesonic lambda decay).

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The production of particles in proton-nucleus reactions at energies below the corresponding threshold for free nucleon-nucleon interactions is a typical cooperative nuclear phenomenon, which is at present far from being understood. Contradictory physical pictures are used to interpret subthreshold production as for example the consideration of high-momentum components in the nuclear wave function, or the inclusion of clusters consisting of several target nucleons, or the discussion of multi-step processes. There is a close relation between the structure of the involved nuclei and the interaction mechanism. It is of principal importance for a correct description of nuclear reactions at relativistic energies to disentangle these two sides of the collision process. Subthreshold production is expected to be a sensitive tool for this purpose.

There are further reasons, peculiar to the type of the considered particle, which make the investigation of K^- production especially interesting:

- K^- mesons may be produced via resonances. Well established resonances are the $\phi(1020)$ meson, which decays with nearly 50% probability into a K^+K^- pair, and the $\Lambda(1520)D_{03}$ resonance, which is, in a first step produced together with a K^+ meson. The $\Lambda(1520)$ decays then with about 23% probability into $p+K^-$.
- Other candidates are the $f_0(980)$ and the $a_0(980)$ mesons the nature of which is still under discussion. It is not clear if they can be considered as usual $q\bar{q}$ states of the quark model or if they should be interpreted as K^+K^- molecules.
- There are recent discussions about a possible admixture of $s\bar{s}$ components in the ground-state wave function of nucleons. As a logical conclusion one may speculate about consequences for strangeness production also at subthreshold energies. The ϕ meson with its predominant $s\bar{s}$ structure may be expected to be a sensitive tool for this purpose.

So far, K^- subthreshold production has been investigated in the reactions $C+C$, $Si+Si$, and $Ca+Ca$ at incident energies between 1 and 2.1 GeV/nucleon [1]. There are, however, no data on K^- subthreshold production in hadron-induced reactions. The interpretation of this type of reaction is a necessary condition for the understanding of heavy-ion interactions.

We propose to measure the energy-dependence of K^- production in proton-nucleus reactions below the nucleon-nucleon threshold at 2.5 GeV incident energy in order to contribute to the solution of the above mentioned problems. The COSY energy range is well suited to this kind of investigation, and the planned 0° Facility [2] together with the expected beam quality of COSY offers good conditions for carrying out also difficult coincidence measurements.

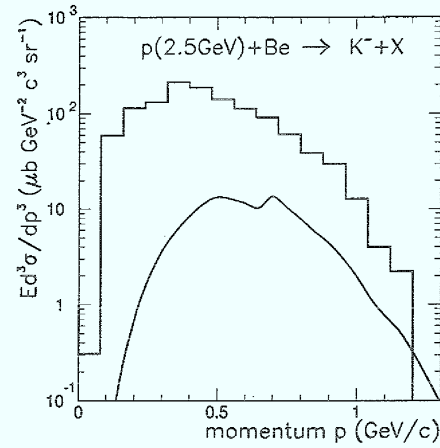


Fig. 1: Folding-model (curve) and ROC-model (histogram) predictions for the K^- -momentum spectra produced in $p+Be$ collisions at $T_p = 2.5$ GeV, integrated over kaon-emission angles $\theta \leq 10^\circ$.

Positively and negatively charged particles are separated in the 0° Facility and will be measured at the opposite sides of the magnet [3].

Since there are no data on K^- production in proton-nucleus ($p+A$) reactions at energies ≤ 2.5 GeV available, we are compelled to rely completely on model calculations in preparing the planned experiments. In order to minimize the uncertainties necessarily connected with such an approach two quite different models, the folding and the Rossendorf Collision (ROC) model, were applied. The inclusive K^- spectra as calculated by the two models are compared in Fig. 1 (for details see [4]). There is a discrepancy of about one order of magnitude in the cross section. This can be considered as a realistic estimate of the uncertainties we have to cope with in preparing the experiment.

The measurement of such an inclusive spectrum is usually considered as a first step to differentiate between models. However, there are problems because of the high background arising from other reaction channels. For both charge states of the kaon the production cross section is about two orders of magnitude smaller than that of pions (and protons). The problem is well-known from the planned investigation of K^+ subthreshold production [5]. In order to reduce the amount of data a fast trigger will be used in the K^+ experiment for finding on-line the true K^+ events out of the background [6]. This system will allow a π^+ suppression of more than 4 orders of magnitude and offer a K^+ detection efficiency of $\sim 15\%$. We intend to use

the same fast trigger also for the measurement of correlated K^+ and K^- mesons [3]. Then at the „negative“ side of the magnet only a relatively simple fast trigger for K^- mesons will be necessary, because the difference in the cross sections for K^- and π^- mesons becomes much less if they are measured in coincidence with identified K^+ mesons as can be seen from Fig. 2.

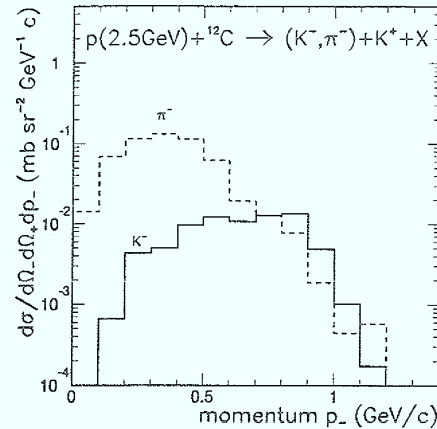


Fig. 2: ROC model predictions for the momentum spectra of K^- and π^- mesons correlated with K^+ mesons at angles $\theta \leq 10^\circ$ and in the momentum interval $0.15 \leq p(K^+) \leq 0.6$ GeV/c.

In the invariant mass spectrum (see Fig. 3) a rather clear signal from the production of the K^+K^- pair via the $\phi(1020)$ resonance can be seen. The region below 1020 MeV is mainly populated via the $a_0(980)$ and $f_0(980)$ resonances. If it turns out that these two resonances are not usual $q\bar{q}$ quark model states then a much smaller cross section in this mass region must be expected, at least from the point of view of the ROC model. Thus, the measurement of K^+K^- correlations can be expected to yield valuable information about the mechanism of subthreshold production as well as about the nature of the involved resonances.

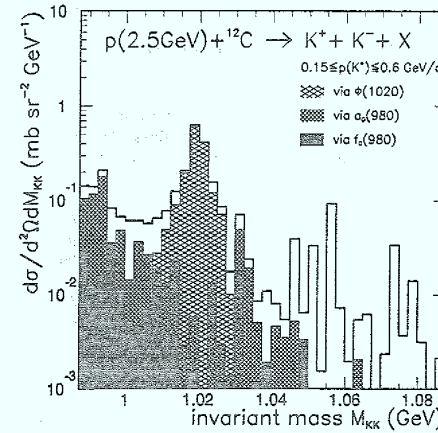


Fig. 3: Predicted invariant K^+K^- mass spectrum integrated over emission angles $\theta \leq 10^\circ$ and the contributions from $\phi(1020)$, $a_0(980)$ and $f_0(980)$ indicated.

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At present the main features of the $p\bar{p}$ interaction in the region 0.1-1.0 GeV/c of laboratory momentum p_Λ are understood. The measured cross sections [1,2,3,4] are described in the frame of the one-boson-exchange (OBE) approach [5]. However, it is interesting also to verify the low-energy $p\bar{p}$ amplitude, that is the S- and P-wave effective-range parameters. One of the ways to get experimental information in the range $p_\Lambda \leq 0.1$ GeV/c is to study the cross section and asymmetry of the reaction

$$p\bar{p} \rightarrow K^+\Lambda p \quad (1)$$

under the kinematical condition that the final Λ and p are produced near the threshold and at small angles. The 0° facility would offer such a possibility. Of course, the cross sections and asymmetries are affected by rescattering of $p\bar{p}$ which complicates the direct investigation of the $p\bar{p}$ low-energy amplitude since one needs to deal with triangle graphs in the final state interaction (FSI). However, in spite of this complication the observable effects should be sufficiently sensitive to the effective-range parameters and also to the off-shell behaviour of the $p\bar{p}$ amplitude.

We have considered the contributions to the cross section and asymmetry in the reaction (1) from the triangle graphs of the $p\bar{p}$ FSI as well as from the pole graphs of the one-kaon-exchange (OKE) and the one-pion-exchange (OPE) which are the main background contributions by the FSI. The results of the calculations for 1.65 GeV proton-beam energy are shown in the Figs. 1 and 2. The exclusive cross

section $\frac{d^5\sigma}{dp_p d\Omega_p d\Omega_K}$ for the case of $\theta_p = \theta_K = 0^\circ$ is shown

in Fig. 1 as function of the momentum of the ejected proton. Use of the $p\bar{p}$ amplitudes of Nagels et al. [5] yields the clear bump around $p_p = 0.8$ GeV/c with a maximum value of about 0.1 mb. Use of the effective range parameters obtained by Takahashi et al. [6] in the region $0.73 \leq p_p \leq 0.85$ GeV/c ($p_\Lambda \leq 0.1$ GeV/c) increases this maximum value to about 0.3 mb which demonstrates the sensitivity of the cross section to the parameterization. In Fig. 2 the analysing power A_x with respect to the longitudinally polarized target is shown. It is zero around $p_p = 0.8$ GeV/c which is due to the fact that near the threshold the $p\bar{p}$ pair is in the S state. This peculiarity in the behaviour of the analysing power could be an additional criterion of the $p\bar{p}$ FSI dominance at this point. The behaviour of the A_x at momenta > 0.8 GeV/c has been calculated separately for the diagrams with virtual pions only and with virtual kaons only. The rather distinct difference could be evidence to discriminate between both mechanisms.

We would like to acknowledge helpful discussion with A. Sibirtsev.

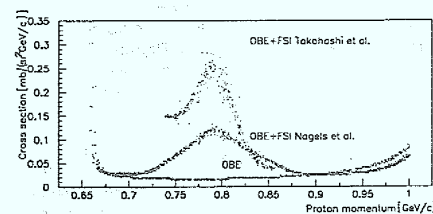


Fig. 1: Exclusive cross section for the reaction $p\bar{p} \rightarrow K^+\Lambda p$ at 1.65 GeV and ejectile angles $\theta_p = \theta_K = 0^\circ$ calculated without FSI and with two different FSI parameterizations.

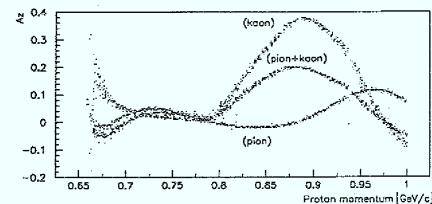


Fig. 2: Analysing power for the reaction $p\bar{p} \rightarrow K^+\Lambda p$ with longitudinally polarized target calculated with exchange of virtual pions and/or kaons in the FSI.

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In COSY-experiments, in which rare processes are studied, e.g. the sub-NN-threshold production of kaons in nuclei (COSY proposal no. 18, at the Zero-Degree-Facility), high luminosities are aimed at. They are obtained with thin, ribbon shaped solid targets at an internal target position. For 10^{16} protons per COSY-cycle luminosities in the order of $10^{33}/\text{s}\cdot\text{cm}^2$ can be achieved, giving rise to reaction rates, which may be difficult to be treated by large acceptance detectors and by the electronic and data acquisition system. The mean rate amounts to $\approx 10^8/\text{s}$ [1]. The maximum rates may be considerably greater, if the target is moved rapidly from its "parking position", where it must be placed during the beam preparation phase, towards the center of the beam. In this operation mode the rates are initially high and diminish, as the number of particles circulating in the ring decreases. The particle losses are mainly due to multiple small angle Coulomb scattering in the target. The bunch structure of the beam causes a further increase of the maximum rates by a factor ≈ 5 to ≈ 10 -compared with the average values-, when the maximum of the bucket passes the target. The experiments actually planned to be performed at the Zero-Degree-Facility do not require changes of the beam energy during the target irradiation phase. Thus also a debunched beam can be used. However in this case the beam momentum decreases continuously because of energy loss in the target. Implications caused by this effect depend on the momentum acceptance of the storage ring and on the dispersion at the target position.

The actual reaction rate follows from the target hit probabilities, which depend on the single particle emittances $\epsilon_{xy}(i)$, the β -functions at the target β_{xy} , and on the size and the position of the target. Also dispersion and momentum deviation interfere. A certain circulating particle passes the target plane within a rectangle of the width $x_0(i) = \pm\sqrt{\epsilon_x(i)\beta_x}$ and the height $y_0(i) = \pm\sqrt{\epsilon_y(i)\beta_y}$. The target hit probability is determined by the degree of overlap of the x_0, y_0 -rectangle with the target. As long as there is no overlap (x- and y-direction) the corresponding particle will not hit the target. Thus the mean target hit probability, which is proportional to the reaction rate, can be controlled through displacement of the target. The goal is to obtain a rate, which is nearly constant and which assumes a level acceptable for the detector read out electronics.

Any simple estimation of the needed target displacements is complicated because of the following facts:

- for a thin target, the complete Coulomb-scattering distribution must be considered.
- the particle density distribution in the transversal phase space strongly depends on the emittances ϵ_{xy} . It has a maximum for small ϵ -values and becomes very small for values close to the ring acceptance.
- the number of circulating particles decreases, the emittances increase strongly with time.

Therefore calculations were performed using, the Monte Carlo program STRAL [2]. In this code Coulomb scattering and the ion-optical transformation of the particle coordinates are simulated for each of the particles and for every revolution in the storage ring. Particles, whose $\epsilon_{xy}(i)$ -values exceed the ring acceptance, are removed. In an extended version of the program the target can be displaced according to a predefined function of the number of bunch revolutions in COSY, or the target position can be controlled by a certain number of desired target hits per revolution. Furthermore the x,y-distributions of the target hits and the variation of these distributions with the number of revolutions are determined.

The aim of the investigation, results of which are given below, was the determination of the type of source displacement, which leads to nearly constant rates. The following input parameters were used:

projectiles: protons of 1 GeV kinetic energy

target: vertically hanging carbon ribbon,
100 $\mu\text{g}/\text{cm}^2$, 2 mm wide, position of
lower border variable (controlled by hit
rate), higher border outside COSY
acceptance

ring acceptance: $\epsilon_{xy} = 130 / 35 \pi\text{-mm-mrad}$

β -functions at target: $\beta_{xy} = 6.6 / 2.7 \text{ m}$

dispersion at target: 0

The assumed ring parameters are close to values which are realistic for COSY, target place TP1 or TP2. The actual calculations were performed for a 1 mg/cm^2 target in order to save computing time. In this region the speed of beam degradation scales with the inverse of the target thickness [1]. The results, quoted below, are derived from the original ones by multiplying the number of revolutions by 10. Thus the results correspond to a target thickness of 100 $\mu\text{g}/\text{cm}^2$. In the simulation calculation 1000 protons were injected and accelerated. The position of the lower border of the target was controlled by the actual frequency of target hits. The aim at was set to a constant hit number per revolution of 5% of the initial particle number.

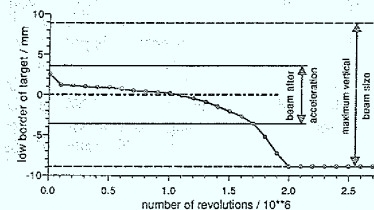


Fig. 1: Position of lower border of the target as a function of the number of particle revolutions in COSY, leading to nearly constant reaction rates.

The resulting target positions are shown in fig.1. The development of the beam during the target irradiation and the target hit rate are given in fig.2. The emittances are the mean single particle emittances ($\Sigma \epsilon(i)/(\text{number of particles})$). We see, that the target hit rate can be kept nearly constant up to $1.7 \cdot 10^6$ revolutions, corresponding to 1.2 s (at 1 GeV we have 0.7 μ s per revolution). As an average a circulating particle hits the target $\approx 70,000$ times before it is lost. For a nuclear cross section of 100 mbarn the reaction probability amounts to $5 \cdot 10^{-7}$ per target hit. Thus $\approx 3\%$ of the particles undergoe a nuclear reaction in the target. The luminosity becomes $3 \cdot 10^{33}$ /s-cm² during the irradiation phase.

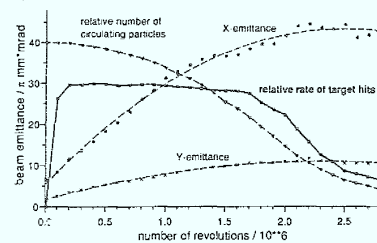


Fig. 2: Time dependence of the target hit rate.

Targets of the type described above will be used in the Zero Degree Facility. Reaction products are analyzed by a dipole magnet, requiring limitation of the reaction zone. Target widths ≤ 2 mm, as assumed above, correspond with the spatial resolution of 2 to 5 mm at the side detectors of the system. The vertical distribution of the target hits is less critical as it does not influence the momentum resolution.

However its knowledge is important for proper ray tracing of the observed reaction products. We see in fig.3, that the hits are concentrated in the region 0 to 5 mm above the beam axis, inspite of a maximum total height of nearly 18 mm. This behaviour is due to the concentration of the single particle emittances around 0.

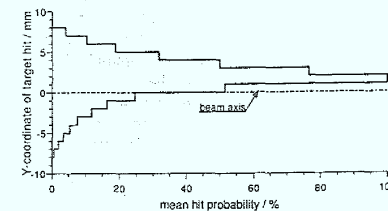


Fig. 3: Distribution of hits on the ribbon shaped target. The Y-coordinate refers to the beam axis, not to the variable position of the lower target border.

The simulation calculations indicate, that the reaction rates obtained with solid carbon targets can be well controlled by positioning the target. At a proton beam energy of 1 GeV the beam life time can easily be extended to ≈ 1 s. Longer life times should be achievable through thinner or smaller targets, or through slower motion of the target towards the beam axis.

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Test at the ITEP synchrotron of a prototype telescope for the K⁺ measurements with the 0° Facility:
Detector setup and performance of the components

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At the 0° Facility [1] a TOF system as well as Cerenkov, ΔE and veto counters will be used to identify unambiguously kaons among protons and pions which are produced with cross sections up to 6 orders of magnitude higher. A prototype of such a telescope has been tested at the synchrotron of the Institute of Theoretical and Experimental Physics (ITEP) in Moscow. The primary $T = 2.5$ GeV proton beam was used to produce secondary particles which were separated according to their momentum. Furthermore, these particles passed through two electrostatic separators where according to their masses clean proton, kaon and pion beams could be selected, see Fig. 1. All our measurements were done at $p = 545$ MeV/c, the lowest momentum where still a reasonable kaon rate was available and which is near the high-momentum end of the spectra which will be studied at the 0°-Facility side detector.

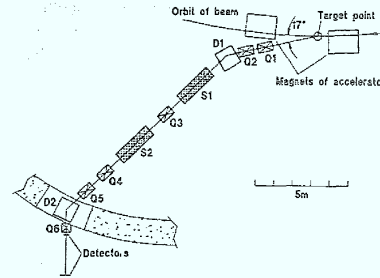


Fig. 1: The beam line for secondary particles with dipole magnets D1 and D2, quadrupole magnets Q1 to Q6 and electrostatic separators S1 and S2.

In this report we will concentrate on the performance of individual counters used. Detailed information about kaon identification and background rejection is given in [2].

Tab. 1: Components of the telescope

Counter	thickness	material
Start counter #1, horizontal	0.4 cm	Pilot U
Start counter #2, vertical	0.2 cm	Pilot U
Stop counter	1.0 cm	BC-408
Cerenkov counter	5.0 cm	GS-218
Degrader #1	7.3 cm	Fe
ΔE counter #1	1.5 cm	BC-408
ΔE counter #2	2.5 cm	BC-408
Degrader #2	1.2 cm	Fe
Veto counter	1.0 cm	BC-408

The telescope employed in this experiment (see Tab. 1) was designed for a particle momentum around 500 MeV/c. We adapted it to the actual beam momentum by an optimization of the degrader thickness. A START counter as proposed in [3] made of PILOT U, 2 mm thick, 5 cm

wide and 30 cm long was placed vertically close to the exit window of the beam line. For the approximately 30 cm long bent light guides Plexiglass GS 233 was used. The counter was read out by 1" photomultiplier tubes Philips XP2972 on both ends. A second scintillator strip was positioned horizontally in the center of the vertical START counter in order to act as an active collimator. This strip was directly read out by two 2" photomultiplier tubes Philips XP2020. This counter was used as reference for all time measurements since its simple geometry led to a better time resolution than was feasible with the START counter described above. The STOP counter, 1 cm thick, was placed 214 cm downstream of the START counter in a frame together with the remaining parts of the telescope (see Tab. 1). These components of a width of 10 cm and a height of 1 m were mounted vertically. Each end of the scintillators (STOP, ΔE_1 , ΔE_2 and Veto) was glued to a fishtail light guide which was then coupled to a 2" photomultiplier tube XP2020.

The data acquisition was based on a CAMAC system controlled by a PC 486. For each event all the amplitudes and the complete timing information were measured with LeCroy FERA ADCs and TFCs. The system gate signal was generated by a coincidence between the two START counters and the STOP counter, using analog meantimers to correct the timing signals for the light propagation time.

The counting rate during all measurements was typically below $2 \cdot 10^4$ cps so that chance coincidences were negligible.

Fig. 2 shows the time-of-flight spectrum for pions and protons measured with the unseparated beam. Due to the low production cross section kaon events are not visible in

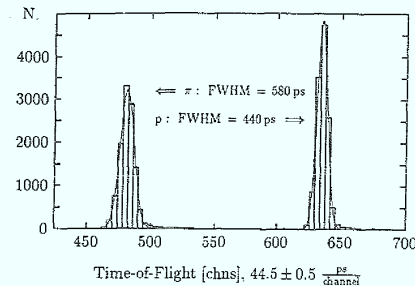


Fig. 2: Time-of-flight spectrum measured with an unseparated beam. The spectrum represents raw data with no offline suppression.

this plot. The spectrum indicates that the beam quality and the performance of the detector system are good. A ΔE spectrum is shown in [2].

Even with our preliminary Cerenkov counter using XP2020 for readout the spectra allow a good discrimination against pions, see Fig. 3. By the use of 3" photomultipliers the light collection efficiency will increase from 30% to 48%. The pion and proton spectra have been generated from the raw data by setting gates on the corresponding peaks in the TOF spectrum. The detection efficiency for protons is 18%, for pions it is 98%, implying a very efficient pion suppression in experiments aimed at kaon detection.

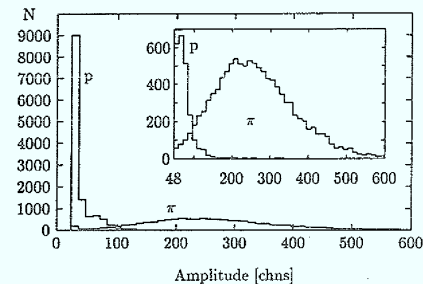


Fig. 3: Charge-amplitude spectra of protons and pions measured with the Cerenkov counter. The inset is the portion of the spectrum for higher channel numbers where pions dominate.

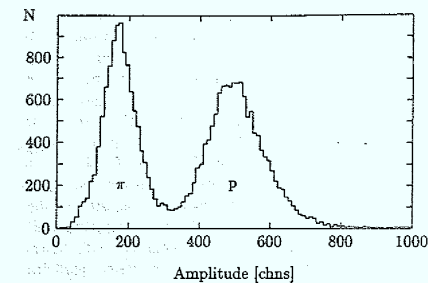


Fig. 4: Charge-amplitude spectrum measured with the START counter with an unseparated beam.

The charge-amplitude spectrum of the vertical START counter is shown in Fig. 4. Despite the expected low light output still a reasonable pulse-height resolution is achieved

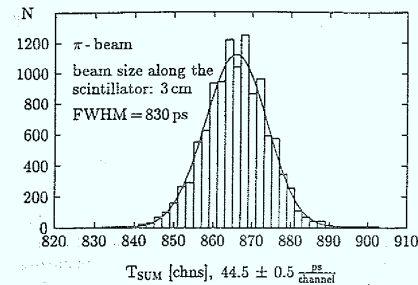


Fig. 5: Timing distribution of the vertical START counter for pions.

as is indicated by the shape of the well separated proton and pion peaks. The distribution of the pion events as a function of the time sum of both PMTs is shown in Fig. 5. The resolution of 830 ps FWHM still needs some improvement. Therefore we consider the following measures:

- (1) The amplitudes of the vertical START counter were limited to 200 mV for 545 MeV/c protons. Therefore the use of a constant fraction discriminator instead of a leading edge discriminator is mandatory. Furthermore the gain of the PMT XP2972 should be increased by modifying the voltage divider.
- (2) Because of the shorter wavelength of the light emitted by PILOT U the use of ultraviolet transmitting Plexiglass, e.g. GS218, instead of GS233 is expected to increase the light transmission by about a factor of 2. The time resolution will be improved accordingly. Alternatively, (standard) Plexiglass GS233 used together with scintillator materials with larger wavelengths (e.g.: BICRON BC408 or NUCLEAR ENTERPRISES NE 122) might lead to a better timing resolution despite their longer time constants. They are 'slower' but according to [4] photoelectron statistics is the dominating factor of the timing resolution.

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Test at the ITEP synchrotron of a prototype telescope for the K^+ measurements with the 0° Facility:
Particle identification and verification of GEANT simulations

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The side-detector system of the 0° Facility [1] will consist of a start-counter wall and of 13 telescopes which will allow to measure K^+ mesons in the momentum range $p \approx 150$ -470 MeV/c [2]. A telescope similar to the ones planned for the 0° Facility has been tested at ITEP, Moscow, where a proton synchrotron delivers protons with energies up to 10 GeV. Secondary particles are generated at internal targets (Be). Using an electrostatic and magnetic separator one can get pure pion, kaon or proton beams with very low background of other particles ($\leq 1\%$). The detection efficiencies for these particles have been extracted from the experimental data and compared with the results of simulation calculations performed with GEANT3.15, modified with MicRes, a code for the hadronic interactions.

At ITEP a telescope consisting of the components given in Tab. 1 and shown in Fig. 1 was used [3]. The thickness of the first degrader, which is necessary to get a good K/π separation in the ΔE counters, was optimized experimentally by maximizing the K^+ energy loss in the ΔE counters. For the ΔE counters a total thickness of 4 cm was chosen. This is equal to the thickness of the ΔE counters which will be used at the 0° Facility for momenta of $p_K \sim 450$ MeV/c. The flight path between the start and stop counters for the TOF measurements was 2.14 m.

Counter	thickness	material
Start counter #1, horizontal	0.4 cm	Pilot U
Start counter #2, vertical	0.2 cm	Pilot U
Stop counter	1.0 cm	BC-408
Cerenkov counter	5.0 cm	GS-218
Degrader #1	7.3 cm	Fe
ΔE counter #1	1.5 cm	BC-408
ΔE counter #2	2.5 cm	BC-408
Degrader #2	1.2 cm	Fe
Veto counter	1.0 cm	BC-408

Tab. 1: Components of the detectors used in the telescope. The start counters were 30 cm high and 3 and 5 cm wide, respectively. The height of all other counters was 100 cm, the width 10 cm. The thickness of the degraders was optimized experimentally. For all counters Philips XP2020 photomultipliers except the vertical start counter (XP2972) were used.

The particle momenta were extracted from the data by comparing the measured TOF difference of protons and pions with simulated TOF spectra for various momenta. In the simulations all energy losses (start counter, air between start and stop counter) were taken into account. This yields a momentum of 545 MeV/c. This momentum is higher than the momenta of the ejectiles at the 0° Facility, but at ITEP only a reasonable kaon rate was available at momenta higher than 500 MeV/c. The momentum acceptance $\Delta p/p$ of the separator was about 3%.

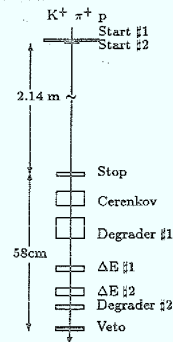


Fig. 1: Sketch of the telescope. The start counters were mounted perpendicular to each other in order to define the beam position. The second degrader in front of the veto counter was necessary to stop kaons which pass through the ΔE counters.

Runs with three different beams - separated kaon or pion beams and a mixed beam with kaons, pions and protons - were carried out during 27 hours and $5 \cdot 10^5$ pions, $5 \cdot 10^5$ protons and $7 \cdot 10^2$ kaons were collected. Fig. 2 shows ΔE spectra for pions and kaons. From these data the gates which have to be set on the various spectra to identify the three different particles were defined. For kaons and pions the separated kaon and pion beams, for the protons the mixed beams were used. A particle is identified as a proton, pion or kaon, respectively, if it lies inside the TOF gate of the corresponding run. The detection efficiencies for the cuts in various spectra were extracted from the data and are summarized in Tab. 2. For instance one can achieve a kaon detection efficiency of 34% with less than 0.01% for pions and protons if the informations ΔE_{Stop} , Cerenkov and $\Delta E_{\Delta E}$ are used. The values for the detection efficiencies using also the signals from the veto counter, which was installed to suppress the pions, are not shown in the table. It turned out that the pion suppression is not improved by this counter. These counters will, however, be useful to reject background of scattered particles at the 0° Facility. At the 0° Facility one expects an even better detection efficiency and background suppression [2] than the values shown in Tab. 2 since here - due to the lower K^+ momenta - the degrader thicknesses can be chosen smaller, a more compact setup is possible and, thus, edge effects (particles which are scattered out of the telescope) and background from hadronic processes contribute less. Furthermore, the total ΔE -counter thickness of 4 cm in the ITEP experiment was smaller than the optimum thickness for particle momenta $p \sim 545$ MeV/c and, therefore, a higher pion

			ΔE_{Stop}	Cerenkov	$\Delta E_{\Delta E}$	$\Delta E_{\text{Stop}} + \text{Cerenkov}$	$\Delta E_{\text{Stop}} + \Delta E_{\Delta E}$	$\Delta E_{\text{Stop}} + \text{Cerenkov} + \Delta E_{\Delta E}$
K	K	Exp.	87.2	93.8	37.0	84.8	34.1	33.6
		Sim.	94.3		47.5		46.3	
	π	Exp.	12.8	17.1	3.5	1.9	0.5	0.009
		Sim.	12.9		3.4		0.3	
	p	Exp.	1.4	99.7	0.033	1.4	0.0	0.0
		Sim.	0.05		0.05		0.0	
π	π	Exp.	98.7	97.3	55.4	96.3	55.2	54.5
		Sim.	94.8		58.9		56.5	
	K	Exp.	96.2	15.2	20.9	13.7	20.4	4.7
		Sim.	98.1		15.6		15.5	
	p	Exp.	4.1	0.6	0.4	0.040	0.033	0.020
		Sim.	0.09		0.16		0.0	
p	p	Exp.	95.7	88.9	99.5	84.9	95.2	84.7
		Sim.	99.9		99.8		99.7	
	K	Exp.	2.4	43.6	39.3	0.5	0.9	0.5
		Sim.	1.9		33.1		1.5	
	π	Exp.	1.1	1.4	39.0	0.19	0.8	0.18
		Sim.	5.2		34.7		2.5	

Table 2:
Measured detection efficiencies in %, relative to the number of particles in the TOF gates for different particles with gates set in the spectra of the corresponding detectors or combinations of these in comparison to the numbers obtained from GEANT simulations. The first column shows the particle for which the cuts in the various spectra are set, the second column the particle for which the detection efficiency is given.

suppression with the ΔE signals at the 0° Facility ($p \leq 475$ MeV/c) will be possible.

The main goal of the measurements was a test of our simulation calculations. In the simulations the geometry described above and a particle momentum of 545 MeV/c were used and a momentum spread of $\Delta p/p = 3\%$ was assumed. The beam profile was approximated by a rectangular distribution. The time of flight between the start and stop counters and the energy loss in the stop and the ΔE counters were calculated. Since the simulation of Cerenkov counters is not possible with GEANT3.15 this

spectra in comparison to the measured ones. It can be seen that there is good agreement between the simulations and the data. The simulated detection efficiencies are also given in Tab. 2.

It should be noted that in addition to the background reduction factors from Tab. 2 using the ΔE_{Stop} , Cerenkov and $\Delta E_{\Delta E}$ informations there is a substantial further reduction of pions and protons from the TOF spectra. In the case of a pure pion beam - that means that one assumes that there are no background particles in the pion beam - only 0.37% of the pions are lying in the TOF gate for the kaons. Thus, a total pion suppression of better than $9 \cdot 10^{-5} \cdot 3.7 \cdot 10^{-3} = 3.3 \cdot 10^{-7}$ can be achieved.

Summarizing, there is good agreement between the data measured at ITEP and our simulated spectra. We, therefore, expect that the K^+ -detection efficiencies and background-suppression factors from [2] which are necessary for the subthreshold K^+ studies at the 0° Facility will be achieved.

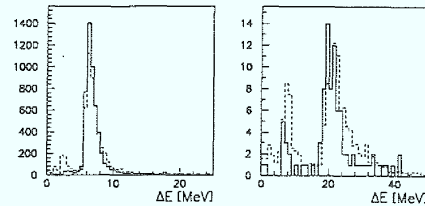


Fig. 2: Comparison of the sum of the energy losses in the two ΔE counters for pions and kaons. The solid line represents the experimental data, the dashed line the simulated spectra. The gate in the pion spectrum was set to $0 \text{ MeV} < \Delta E \leq 11 \text{ MeV}$ and in the kaon spectrum to $\Delta E \geq 18 \text{ MeV}$. The low energy peak in the pion spectrum is caused by pions which hit only the first ΔE counter. The low energy peak in the kaon spectrum is caused by μ^+ 's from K^+ decay.

counter is treated only as a passive degrader. To compare the simulated and measured ΔE spectra the latter were calibrated linearly: In the stop counter the maximum of the pion and proton peaks, in the ΔE counters the pion and the $\Delta E=0$ MeV peaks were used. Fig. 2 shows the simulated

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- [1] K. Sistemich et al.: contribution to this report, p. 41
- [2] A. Franzen, M. Büscher: contribution to this report, p. 57
- [3] R. Eßer et al.: contribution to this report, p. 52

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An array of telescopes consisting of scintillation and Cerenkov counters as well as degraders will be used for particle identification in the side detector of the 0° Facility [1,2]. The telescopes will be mounted on a movable platform coupled with the spectrometer dipole D2 so that all telescopes have a well defined position with respect to the spectrometer. Each telescope will be an independent module giving maximum flexibility in its use. Thus it can be installed in other experimental areas for tests and physics experiments. It is foreseen that the telescopes can be adjusted to the mode of operation of the 0° Facility by shifting them on the platform so that for different magnetic fields in D2 each telescope is hit by particles having momenta for which it has been optimized.

The telescopes have to be very compact because any gap between successive counters would reduce the angular acceptance. A solution with standard straight connections between PMTs, light guides and scintillators is rendered impracticable because the size of the PMTs and their bases would cause a large telescope thickness. On the other hand, long and bent light guides are not affordable since they would deteriorate in particular the very good timing resolution which is crucial for the required high suppression of background events. Therefore we have considered various coupling geometries including rectangular joints between scintillators and photomultiplier tubes in order to find an optimal geometry.

Each telescope will have a support structure which is open towards D2 thus avoiding that any structural material is hit by particles. It must not exceed the 10 cm width of the individual scintillation counters and have no protruding components which would interfere with neighbouring telescopes. Stress on the delicate scintillator material due to thermal expansion must be avoided as far as possible. Therefore a self-compensating support is foreseen using either springs or counterbalances. Optical coupling between PMs, light guides and scintillators will be realised by gluing so as to guarantee a reliable joint with optimal light transmission. It is planned to achieve the light tightness with tubes of black plastic foils which enclose each paddle including both photomultiplier tubes.

References:

1. K. Sistemich et al., contribution to this report, p. 41
2. O. Schult et al., Ann. Rep. KFA-IKP 1993, p. 39

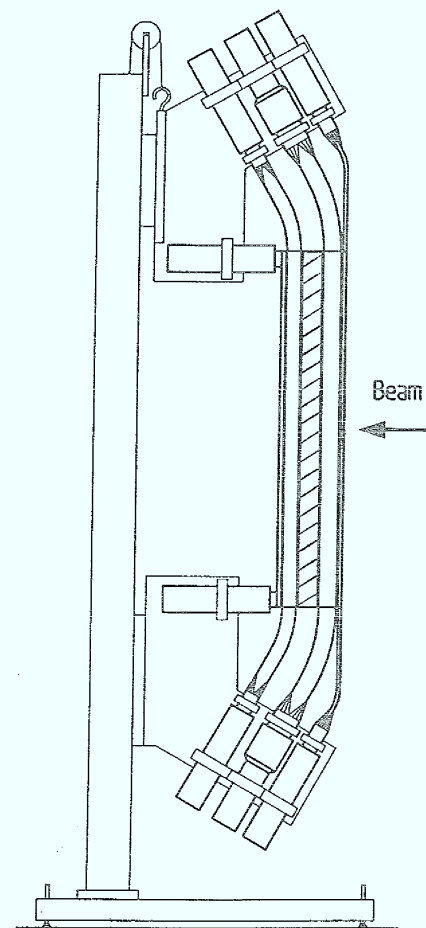


Fig. 1: Design variant of a telescope for the 0° Facility. It is indicated that the structure holding the upper photomultiplier tubes is movable and counterbalanced by a weight which compensates for thermal expansion. Along the beam the first counter is a stop counter which is followed by a Cerenkov counter, a degrader (hatched), a ΔE counter, a second thin degrader and a Veto counter.

The side-detector (SD) system of the 0° Facility [1] will mainly consist of multiwire proportional chambers for track reconstruction and of detectors for TOF and ΔE measurements to identify kaons and to suppress the large background (protons and pions coming directly from the target and particles scattered at the magnet iron). The passage of the ejectiles through the spectrometer and the time-of-flight and ΔE spectra were simulated. From these spectra the detection efficiencies and counting rates for kaons and background particles were determined. The simulations were made with GEANT3.15, modified with MicRes, a code for the hadronic interactions. All effects as decay in flight and hadronic interactions in parts of the dipole have been included.

The geometry of the 0° Facility and the detection system are shown in [1,2]. The SD system consists of a TOF-start counter wall and 13 telescopes which cover a momentum range $p=140-480$ MeV/c. The telescopes are placed perpendicular to the 0° trajectories. Each telescope is 10 cm wide corresponding to a momentum spread $\Delta p/p$ of 13% (1st tel.) - 7% (13th tel.) and will consist of TOF-stop, ΔE and veto counters, and above 300 MeV/c (8th telescope) in addition Cerenkov counters will be used. Furthermore, the telescopes contain degraders in front of the ΔE counters in the imaging surface of the spectrometer dipole to tune the energy of the ejectiles in such a way that the K^+/π separation is optimized. Two apertures at the entrance of the spectrometer magnet restrict the ejectile angles to $\theta \leq 15^\circ$. The two wire chambers also serve for further background suppression.

In the simulations the response of the telescopes and wire chambers have been simulated by tracking the particles from the target through the magnetic field into the telescopes. For the simulation of the kaons a 'realistic' momentum and angular distribution from [3] was started at the target. For the simulation of the background, protons

with $T_p=1$ GeV were started in front of a ^{12}C target and the reaction products were traced. Here particles coming directly from the target as well as the particles which are scattered at the magnet iron were taken into consideration. It turned out, that the number of particles which are scattered in the magnet iron is roughly equal to the number of protons and pions coming directly from the target, see Fig. 1.

With a realistic luminosity of $4 \cdot 10^{32} \text{cm}^{-2} \text{s}^{-1}$ (on an average $4 \cdot 10^{14}$ protons per second on a $20 \mu\text{g}/\text{cm}^2$ ^{12}C target with 10^{10} protons in the ring) one expects a total counting rate of about 20000 events per telescope per second. Only those particles are accepted which hit the right combination of start and stop counters corresponding to emission angles at the target of $9 \leq 10^\circ$. With this coincidence matrix one can suppress 35% of the scattered background particles. Gates were set on the TOF and ΔE spectra in each telescope. Fig. 2 and 3 show the spectra in the 7th telescope (275-300 MeV/c) for TOF and $\Delta E_{\Delta E}$. It can be seen from the figures

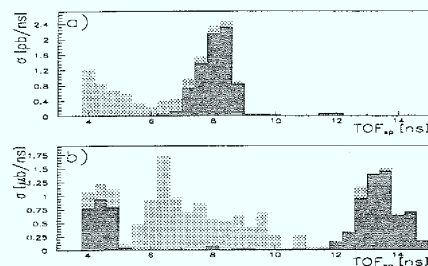


Fig. 2: a) TOF spectra for kaons and myons, b) TOF spectra for the background events. The gate for K^+ detection is set to $7.1 \text{ ns} \leq \text{TOF}_{\text{sp}} \leq 9.0 \text{ ns}$.

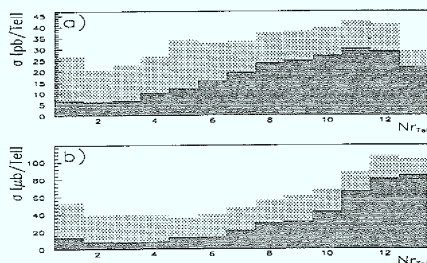


Fig. 1: Number of events in the stop telescopes 1 through 13. The upper spectra (a) show the kaons (black) and the myons (dotted) from kaon decay, the lower (b) the background events (direct particles black, scattered particles dotted).

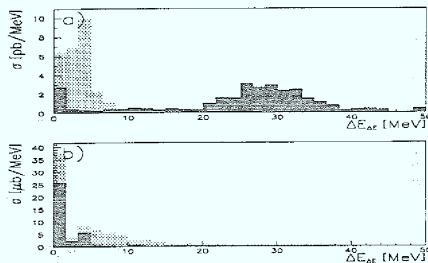


Fig. 3: a) $\Delta E_{\Delta E}$ spectra for kaons and myons, b) $\Delta E_{\Delta E}$ spectra for the background events. The gate for K^+ detection is set to $14 \text{ MeV} \leq \Delta E_{\Delta E} \leq 40 \text{ MeV}$.

Telescope Nr.		1	2	3	4	5	6	7	8	9	10	11	12	13
Decay in flight		31	31	32	32	33	33	34	35	35	36	36	36	36
TOF	K	100	92	92	92	88	99	96	88	97	99	95	89	100
	p, π scat.	16	20	17	25	15	34	28	32	51	51	43	30	37
$\Delta E_{\Delta E}$	K	88	81	78	71	72	77	69	67	68	59	67	52	51
	p, π scat.	22	24	7	27	10	11	17	8	17	21	12	12	15
Veto	K	95	94	88	97	97	97	95	88	97	97	93	96	91
	p, π scat.	15	20	19	30	27	27	36	37	46	48	56	74	77
Wire chambers	K	95	89	92	94	93	97	98	98	96	97	93	94	100
	p, π scat.	8	13	7	5	6	7	5	9	6	11	4	7	8
All	K	78	68	64	70	65	74	65	54	65	58	60	48	49
	p, π scat.	0	0	0	0	0	0	0	0	0	0	0	0	0

Tab. 1: Detection efficiencies for kaons and background particles in %. Since all protons and pions coming directly from the target (~ 11000 particles in the simulation) can be rejected by using TOF and $\Delta E_{\Delta E}$, only the numbers for protons and pions scattered at the magnet iron are given. The numbers for the survival probability can be taken from line 'Decay in flight'. 'TOF' is the time-of-flight difference between the start and stop counters and ' $\Delta E_{\Delta E}$ ' the energy loss in the ΔE counters of the telescopes. 'Veto' includes the time-of-flight difference between the stop and the veto counters in one telescope. 'Wire chambers' uses only the height information from the wire chambers. 20-50% of the kaons are lost in the counters due to edge effects. The efficiencies for kaons are relative to the number of kaons reaching the TOF-stop counter. They have to be corrected for decay in flight between target and stop detectors. Note, that the number in the line 'all' is not the product of the other numbers.

that the suppression of the particles coming directly from the target (protons and pions) is very good. One can reject all direct background protons and pions by only using the TOF and the $\Delta E_{\Delta E}$ information, in accordance with [4]. However, some of the scattered particles lie inside the gate for kaon identification. But, since the scattered particles come from other directions than the direct ones, the information from the two wire chambers will be used for further suppression. Thus, with the information on the trajectories a detection efficiency of nearly 100% for the kaons and a suppression factor of more than 90% for the scattered background can be achieved. It must be noted, however, that this information can only be used in the off-line analysis; it cannot be included in the fast trigger, whereas the height information from the start and stop counters can be used on-line (see below). Because of the good suppression of the direct background particles the results for the detection efficiencies for each telescope are given in Tab. 1 only for the scattered background. One can see, that the $\Delta E_{\Delta E}$ signal is most efficient to reject scattered background particles - apart from the off-line information one gets from the wire chambers. Using TOF, ΔE , veto (Rejection of particles giving a signal in the veto counters shortly after the TOF signal. As the kaons will be stopped in the ΔE counter, the μ 's from kaon decay will give a delayed signal in these counters ($\tau_K = 12.4$ ns) and the wire chamber information one can gain detection efficiencies between 20 and 25% for kaons - including decay in flight - while suppressing all background particles (~ 23000 particles in the simulation). Due to decay in flight only 30 to 35% of the kaons will reach the stop telescopes. The simulation of Cerenkov counters is only possible with GEANT3.16, which was not available during these studies, but test measurements at ITEP/Moscow showed, that one can suppress the background by a factor of ~ 10 using these counters.

The total rate summed over all telescopes will be $\sim 3 \cdot 10^5 s^{-1}$. Using the information from the time-of-flight and energy loss, and height information from the start and the stop counters with a spatial resolution of $\Delta x = 4$ cm for the stop and $\Delta x = 5$ cm for the start counters, one can achieve an on-line suppression by a factor of $\sim 4 \cdot 10^2$ leading to 750 events per second on tape, which is possible.

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1. K. Sistemich et al., contribution to this report, p. 41
2. O. Schult et al., IKP Annual Report 1993, p. 39
3. A.A. Sibirtsev, M. Büscher, Z. Phys. A347, 191-202 (1994)
4. S. Kopyto, M. Büscher, IKP Annual Report 1993, p. 49
5. A. Franzen et al., contribution to this report, p. 54

Design of a Multiwire Proportional Chamber for the 0°-Facility

L. v. Horn, D. Gotta, H. Ohm, H. Wyrwich^a, Th. Kirchner^b, H. Müller^b, Ch. Schneider^b

At the 0°-Facility the tracks of the ejectiles will be measured with the use of two Multiwire Proportional Chambers MWC1 and MWC2 [1], providing basic information for the reconstruction of the ejectile momenta and emission angles at the target.

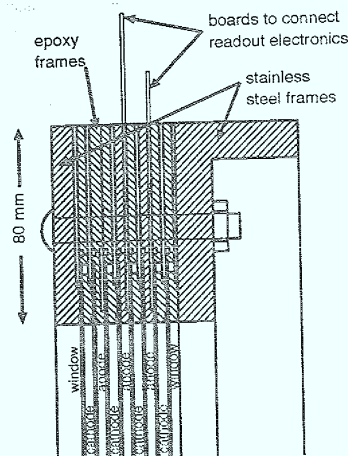


Fig. 1: Drawing of the design of the multiwire chamber MWC1. A cut through the upper part of the frame is shown.

A stable operation of a multiwire chamber under high voltage and a uniform response over the whole active area can only be guaranteed if the electrodes stay precisely at their positions. For this reason, wires and foils are mounted on their frames with a certain tension which for the wires is close to their tensile strength. The resulting forces on the frame will cause a bending of it which means a reduction of the wire tension. Usually, for multiwire chambers the frames are chosen so massive that the reduced tension will stay in the regime where a stable operating is guaranteed. However, for the chamber close to the dipole magnet D2 of the 0°-Facility, there will not be sufficient space, so that this simple approach is not possible (see also [2]). Therefore, for this chamber, the wires and the foils are mounted on prebent frames which are made of epoxy material and stainless steel to gain stability, see Fig. 1. The special assembling method requires a certain knowledge of the shape of the bent frames which was calculated by the „Finite Element Methode“ (FEM) at ZAT, KFA-Jülich. To reduce computer calculation time the FEM code was run for a simplified model of the chamber where smaller grooves were omitted. Because of symmetry, it is sufficient to use a model which corresponds to a quarter of the whole frame. Tensions in the material and the torsion due to the load of anodes and cathodes were investigated in the calculations [3] but no critical values were detected. They would have

required a redesign. The main question for the calculations was the shape of the bent frame. In the calculations a slight decrease of the foil tension, which may happen due to relaxation effects, was taken into account. For this reason, the calculation was repeated for another load assuming only half of the foil tension, which would mean a strong relaxation, see Fig. 2. A change of the shape of the frame caused by a decrease of the foil tensions will lead to a higher wire tension by increasing the length of the wires. But the two calculations show that even for an assumed extreme reduction of the foil tension the wire tensions will stay in the range where a stable operation is guaranteed.

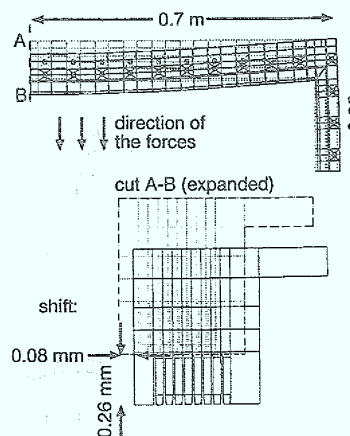


Fig. 2: The model of the multiwire chamber MWC1 on which the „Finite Element“ calculations were performed (dashed lines). The bent frame is drawn with solid lines. The shift has been increased by a factor of 100 to make it visible.

In the assembling process the wires and foils will be mounted on frames which are prebent to the calculated shape. So the FEM calculations prove that the stability of the chamber frame will be sufficient for the proposed assembling and mounting method providing informations on the bending of the frame which are essential for the assembling process.

^a ZAT, KFA Jülich, Germany

^b FZ Rossendorf, Dresden, Germany, supported by BMFT

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1. K. Sistemich et al., contribution to this report, p. 41
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The Multiwire Chamber MWC 1 for the 0°-Facility

Ch. Schneider^a, Th. Kirchner^a, H. Müller^a, L. v. Horn, H. Ohm, V. Kruglov^b

At the 0°-Facility two multiwire chambers will be installed for the reconstruction of the particle trajectories. The main task of these chambers is therefore the position measurement of the particle tracks.

In this contribution we concentrate on the chamber MWC 1 which is, together with the start scintillators, located near the magnet. Since the achievable position resolution is restricted to about 5 mm mainly due to small angle scattering [1], we decided to build a chamber in which only anode readout is employed. So we can use conductive foils as cathodes.

With the chosen wire distance of 2.54 mm we are below the necessary resolution and have an easy connection to the pin raster of the readout electronics, RAL 110, 111 [2]. For the cathodes we have looked for a material with high resistance in order to limit the current inside the chamber in the case of spark breakdowns. Furthermore the conductive material should have a high electron work function because the photo electrons produced in the cathode material increase the noise and worsen the operating conditions inside the chamber. From our point of view a foil with carbon as conductive material fulfills these conditions better than the usually used aluminium.

To test the counting behaviour of such a type of chamber we built a small test chamber with a sensitive area of 200 mm x 145 mm and with wire distance 2.5 mm. The wires with diameters 20 μ m are made from tungsten covered by a thin gold layer. The mylar foil we used has a thickness of 26 μ m with a 6 μ m thick carbon layer on it. Test measurements with a 2.5 GeV beam consisting of pions, protons and kaons were performed at the ITEP Moscow.

In this experiment our test chamber and a beam monitor were located between arrays of scintillation counters used as trigger for the efficiency measurement of the chamber. If we compare the sensitive areas and spatial positions of the counters, the angular acceptances α_1 , α_2 , and α_3 (see Fig. 1), along the vertical direction we get for the trigger counters $\alpha_3=26^\circ$, for the test chamber $\alpha_1=6.6^\circ$ and for the beam monitor $\alpha_2=6.5^\circ$. From the beam monitor we know

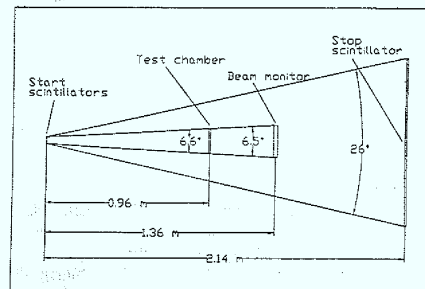


Fig. 1: Schema of the detector arrangement from the side view.

that most parts of the gaussian like profile of the beam are inside the sensitive area of the beam monitor. The beam monitor can detect 88% of the full beam due to the size and position of its sensitive area. That is why we can detect with the given angular acceptance of the test chamber only about 90% of the full beam. From Fig 2 we see that we measure 91% of all beam particles. This is about 100% efficiency in our estimation about the angular acceptances of the counters. A second indication that the detection efficiency is close to 100% is the plateau in Fig. 1 for high voltage > 2900 V. This result gives confidence in the reliability of the concept chosen for the chamber.

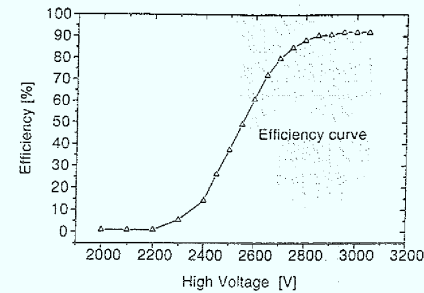


Fig. 2: The Efficiency curve shows the ratio of counted to the triggered events in the chamber.

The complete outline of the start chamber with a sensitive area of 1303 mm x 350 mm is shown in Fig. 3. The outer dimensions are 1403 mm x 510 mm. For the determination of the particle position in horizontal (x) and vertical (y) direction we have 3 anode planes with 512 wires each, denoted as Anode I-III in Fig. 3. The plane Anode I with the wires in vertical direction measures the x co-ordinate and the planes Anode II+III with inclined wires, ± 30 degree, the y co-ordinate. Corresponding to the 3 anode planes there are 4 cathode planes denoted as Cathode I to IV. The frames of all anode and cathode planes are made from a fibre glass epoxy compound called Stesalit.

Cathodes II and III have mylar foils with carbon layers on both sides, Cathodes I and IV have only one layer of carbon on the side that is oriented to the anode planes. Because the gap between anode and cathode equals the width of a Stesalit frame the cathodes and anodes are mounted always on the same side of the Stesalit frames. In Fig. 3 these are the sides in direction to the frame Support I.

The anodes are on ground and the high voltage is on the cathodes, so that the window foils should be again on ground potential to have electrostatic equilibrium conditions in the chamber. Therefore we use an additional cathode frame (Window I) as fixing frame for the front window to have the distance of 5 mm between the cathode and the window foil. The window foil consist of 25 μ m Kapton.

On the top sides two support frames (Support I, II) stabilise the chamber. Since the chamber MWC 1 is placed in the fringing field of the magnet D2 the support frames are made from stainless steel. The second gas isolating window is fixed on the frame Support II.

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- a Forschungszentrum Rossendorf, Germany, supported by BMFT and Forschungszentrum Jülich
b JINR Dubna, Russia

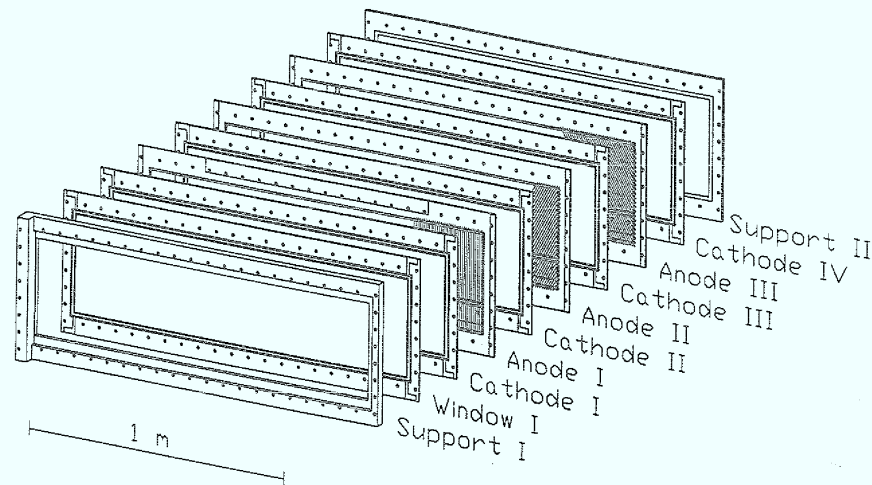


Fig. 3: Components of the start chamber, the second window is fixed on the frame Support II

The goal of the present article is the estimation of the widths of the stop modules of the side detector system at the COSY ZDF setup [1]. Estimations were made taking into account the imaging properties and aberrations of the spectrometer dipole D2 as well as particle deflection due to multiple Coulomb scattering in the various components of the setup.

We consider the K^+ mode [2] with a homogeneous magnetic field of 12.8 kG. The placement of the stop telescopes [3] corresponds to the setup defined earlier. We deal with ejectiles with emission angles of $-10^\circ < \vartheta < 10^\circ$ relative to the axis of the proton beam impinging on the target. The vertical emission angle is $\phi=0^\circ$.

In order to determine the imaging surface of the magnet D2, two particles with the same momenta and horizontal angles $\vartheta = -5^\circ$ and $+5^\circ$ were generated as emitted from the target located at 35 cm in front of the magnet iron. The intersection point (X_L, Z_L) of these trajectories defines the focal surface.

In real experimental conditions the trajectories will be distorted by multiple scattering. This distortion is higher for low momenta. On the other hand the corresponding spread for low and high momenta will be almost equal, because of the long flight path for higher-momenta particles. The crossing points of particles with equal momenta with uniformly distributed emission angles define a 'spot' in the XOZ plane. The sizes of this 'spot' are determined by the continuous character of multiple-scattering fluctuations. In fact, the envelope of the intersection points of all trajectories is a volume which is oriented along the track of particle with $\vartheta = 0^\circ$, i.e. along the corresponding module axes (Z_L coordinates). The transverse size (X_L coordinates) sets a lower limit for the width of the designed module since the module width must be at least not less than the transverse spot size.

The projected spot sizes defined purely by geometrical aberrations are presented in Fig. 1 at minimum and maximum momenta and in Fig. 2 for K mesons including multiple scattering. To obtain these distributions two tracks with the same momenta and uniformly distributed in the range of emission angles ($-10^\circ, +10^\circ$) were generated. In these distributions the point $Z_L = 0, X_L = 0$ corresponds to the 'middle' point defined as the crossing point of tracks with $\vartheta = -1^\circ, +1^\circ$. It can be seen that the spot shape is increasingly nonsymmetric with momentum. When multiple scattering is included this asymmetry is somewhat smoothed.

From these figures we conclude that the spot width σ_X (standard deviation) does not depend on particle momentum and approximately equals $\sigma_X = 3.0$ cm for K mesons. It means that in the momentum range 150 - 600 MeV/c the module transverse sizes would be recommended the same for all modules and equal to be in the range 10 - 12 cm ($\pm 2\sigma_X$).

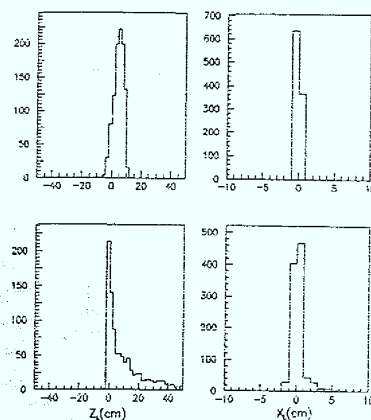


Fig. 1: The spot sizes defined by aberration effects only in the module local frame at minimum (upper) and maximum (lower) momenta (~ 150 and ~ 600 MeV/c).

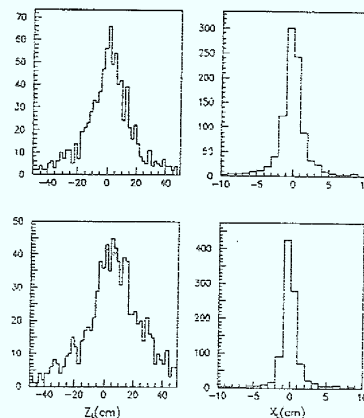


Fig. 2: Spot sizes for K^+ mesons in the module local frame (accounting for aberration and deflection) at minimum (upper) and maximum (lower) momenta.

^a HEPITU Tbilisi, Georgia; supported by the International Science Foundation

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Expected missing-mass resolution in the deuteron-fragmentation experiment at the ZDF
and conditions for studies of the $\bar{p}p \rightarrow p\bar{p}p$ process

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One of the crucial points for the deuteron break-up study at the 0° Facility (ZDF) is the possibility to distinguish between simple break-up with only three nucleons in the final state and fragmentation with pion production. Separation of both processes can be obtained from the missing-mass spectrum deduced from the measured momenta of the forward and backward emitted protons. The quality of the separation depends as well on the missing-mass resolution as on the relative intensity of the processes at the special kinematic choice.

All the simulations were performed using the GEANT program package. The description of the set-up included field maps of the dipole magnets D1 and D2, the material in the particle paths and the spatial resolution of the multiwire chambers ($\sigma_t = 330 \mu\text{m}$ for the forward MWPC and $\sigma_b = 200 \mu\text{m}$ for the backward MWDC). Energy losses, multiple scattering, and nuclear interactions of the particles in the materials were taken into account. A deuteron gas-cell target [1] was considered with a cylindrical aluminum cell of 2 cm diameter, 20 cm length (along the beam), and 0.5 mm wall thickness. Momentum and angular dispersion of the beam of 2 mm diameter in the interaction region were neglected. The set of pd -interaction events for 2.5 GeV beam energy was produced using the code of the Rossendorf collision model [2]. The proton momenta at the target were determined from the track coordinates in the detectors by the reconstruction procedure described elsewhere [3]. The resulting missing-mass spectra for an ideal pointlike target and the extended one are shown in Fig. 1. The broadening of the mass-line for the missing neutron, caused by the large dimension of the interaction region, does not abolish the possibility to separate the break-up channel and the rather low background from the pion-production channel.

The same kind of calculation has been performed for a study of ϕ -meson production in the interaction of the polarized proton beam with a polarized proton target: $\bar{p} + \bar{p} \rightarrow p\bar{p}\phi$. This process attracts attention as a possible way to search for manifestation of strange quark-antiquark pairs in the structure of the proton [4]. As a first stage we have looked for the acceptance of the spectrometer while detecting the proton pair in the forward detector and the K^+ meson from ϕ decay in the side detector. The calculation has been performed for a beam energy 6 MeV above threshold. The field strength in the dipole magnet D2 was optimized for the momentum acceptance yielding a detection probability of 23% determined by the acceptance and the K^+ decay. The simulated ϕ missing-mass line has a FWHM of 1.3 MeV which demonstrates the promising conditions for this study.

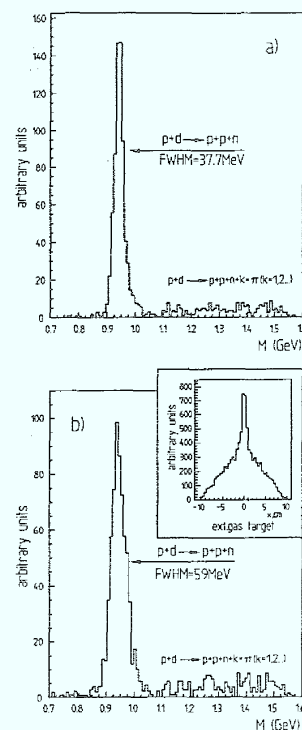


Fig. 1: Calculated missing-mass spectra with a: pointlike target, b: cell target with deuterium-density distribution as shown in the insert.

^a Joint Institute for Nuclear Research, Dubna, Russia; supported by grant 93-02-3745 of the Russian Foundation for Fundamental Research

^b FZ Rossendorf, Dresden, Germany, supported by BMFT

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The preparation of the Backward Detector System for the 0° Facility

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The backward detector assembly is an essential part of the 0°-Facility (ZDF) detector system permitting the study of particle ejection in a kinematic region forbidden for free N-N collisions [1,2]. It consists of three drift chamber packages, scintillation counter hodoscopes, and anticoincidence veto counters. The detector is dedicated to angular and momentum measurements on particles with emission angles close to 180°. Mass identification is achieved by $\Delta E1$, $\Delta E2$, and Veto counters. The hodoscope can also provide event arrival times when it is operated as a trigger in coincidence measurements with other ZDF detectors.

The most time consuming part of the drift chamber production - manufacturing of the carrying frames and assembling of the chambers - is completed at present. Wire

winding and mounting of the chambers are scheduled for the next year. The work on the lay-out of the hodoscope now is in the final stage. The scintillators are under production and the counter-prototype modules have been prepared. The modules consist of 36 cm long plastic scintillators with cross sections of 6×2 cm² and 6×0.5 cm². The scintillators from both ends are viewed by PMT XP4222B. The first test of the modules has been carried out at the JINR phasotron. The beam consisted of positrons, pions and muons of 250 MeV/c momentum at relative intensities of about 1 : 0.5 : 0.5. Reasonable time and amplitude properties of the modules were found (Fig. 1 and Fig. 2, respectively). The time resolution of the thick scintillator is sufficient for time-of-flight measurements between the backward and the forward detectors. The measured amplitude distributions have been compared with

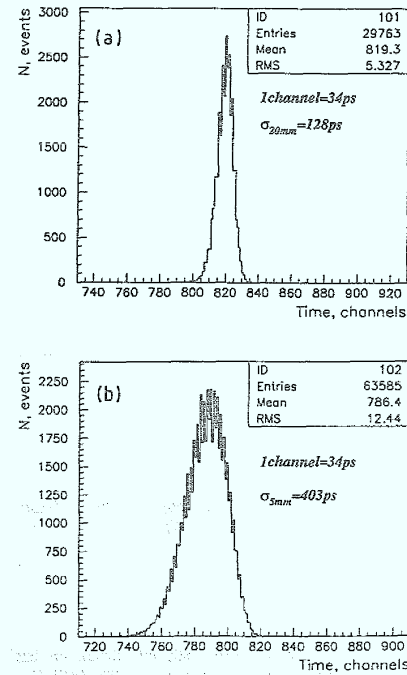


Fig. 1: Time spectra measured for coincidences between the hodoscope modules: a) Spectrum for 20 mm thick start module and 20 mm thick stop module; b) Spectrum for 5 mm thick start and 20 mm thick stop module. Single counter RMS values are given.

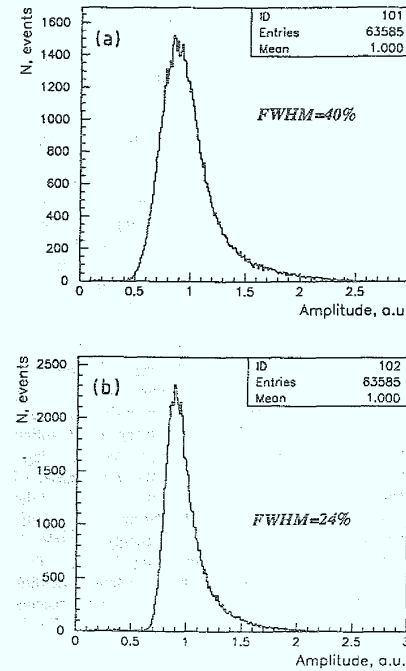


Fig. 2: Amplitude spectra measured with the hodoscope modules of 5 mm (a) and 20 mm thickness (b).

energy-loss distributions calculated with the GEANT program package. Relative FWHM of about 30% and 10% were obtained for the 5 and 20 mm thick counter, respectively, in reasonable agreement with the measured numbers of Fig. 2.

Two-dimensional distributions of the energy losses $\Delta E1$ and $\Delta E2$ in the scintillators of the two subsequent hodoscope layers have been calculated for the ejectile pions, kaons and protons. It is obvious from Fig. 3 that the pion and proton regions are well separated in spite of the angular and momentum dispersion of the ejectiles within the backward spectrometer acceptance and the rather crude segmentation of the hodoscopes (eight modules of equal width in a layer). The measured detector dispersion does not significantly broaden the regions and thus the good proton-pion separation will be maintained in the experiments.

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* Supported by grant 93-02-3745 of the Russia Foundation for Fundamental Research

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1. V.I. Komarov et al.: Ann. Rep. 1992 of the IKP, KFA Jülich, Report Jül-2726 (1993), p. 26
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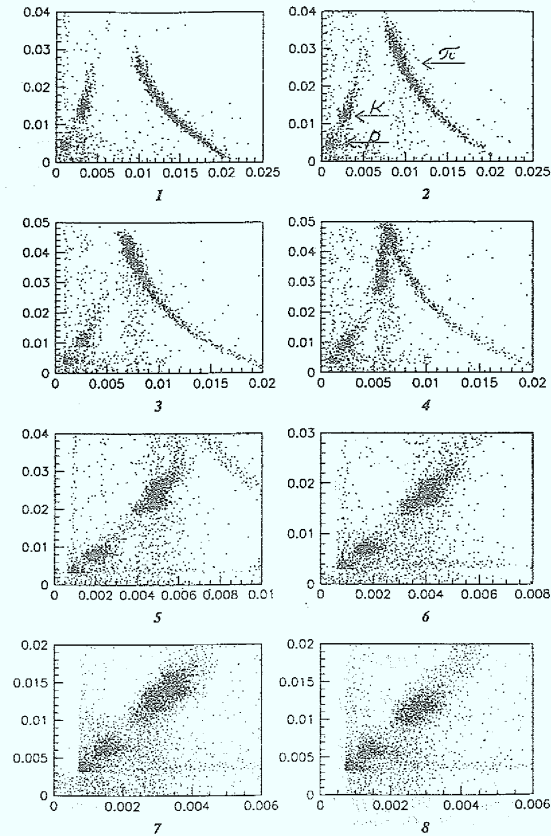


Fig. 3: GEANT simulated energy losses in the counters of the two hodoscope layers (abscissa: 5 mm thick counters of the first layer; ordinate: 20 mm thick counters of the second layer). The plots 1 to 8 correspond to the 8 counters hit by particles with momenta increasing from 250 MeV/c in the first to 600 MeV/c in the eighth counter.

the integrated ΔE signals. This guarantees a minimum energy loss in the ΔE counter, whereas Cerenkov and Veto counters are used to reject high-energy particles. The OR-ed signals of all stop telescopes K_j define the 1st level trigger for the side detector. This trigger is the start signal for all TDCs, the gate signal for all ADCs, and the start signal for the readout of the MWPCs.

For a precise timing of the trigger, it is essential to put a lot of care into the lay-out of the whole electronic system of discriminators, meantimers, the TOF trigger, and the interconnection of all components.

- Constant fraction discriminators will be used to compensate for the amplitude walk effect. Meantimings derive a position independent signal from the 2 photomultipliers of each scintillator.
- Special integrators give a position independent energy signal of every 2 photomultipliers of each AE counter. The subsequent leading edge discriminators allow to set a lower threshold on the energy loss in the AE counters. This results in a rejection against particles with a small energy loss, e.g. scattered high-energy pions coming from the iron of the magnet D2.
- Individually shielded twisted pair cable reduce the cross talk between neighbouring digital signals and therefore reduce the time shift due to this cross talk.
- A sharp veto timing is necessary not to reject products from the 12ns kaon decay.

The time walk measured with a prototype of a CAMAC Constant Fraction Discriminators developed in collaboration with JINR Dubna [3] is shown in fig. 2. This unit is now under redesign in order to reduce cross talk. A new design will be available in early 1995. It will be equipped with on board meantimers described in [4].

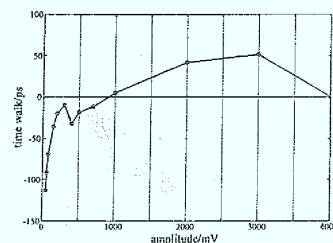


Fig. 2 The time walk versus input-pulse amplitude measured with the CAMAC Constant Fraction Discriminator developed by JINR Dubna and the IKP. For a wide amplitude range the maximum time walk is below 50ps.

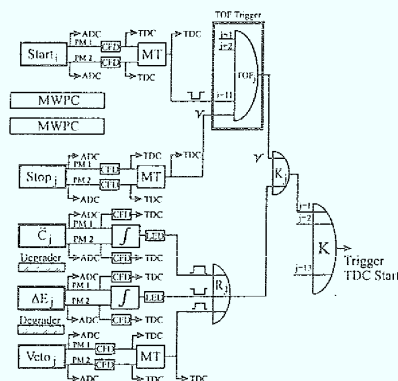


Fig. 1: Principal trigger scheme for one stop telescope. Start- and stop counter define the input signals for the TOF trigger. ΔE counters guarantee a proper energy loss for any triggered particle, whereas Cerenkov and Veto counters are used to reject high-energy particles. Upper and lower end of all scintillators are equipped with photomultipliers.

The fast trigger is planned as a one-particle trigger but can easily be extended to a multi-particle trigger. It is based upon events in the stop scintillation counters. According to the emission angle of the particle at the target, a hit in one of the stop counters corresponds to a hit in one of up to 11 out of 23 start counters. Thus for each event in one of the stop counters, the fast TOF trigger has to check whether the time difference between the signal in the stop and the corresponding start counter falls into the TOF window of the particle searched for, e.g. a kaon.

Equipped with Constant Fraction Discriminators (CFD) and Meantimers (MT) start and stop counters define the input signals of the TOF trigger. Leading Edge Discriminators (LED) are used to set a lower threshold on

Measurements on standard (unshielded) twisted-pair cables show appreciable time shifts for digital signals (rise time 1ns, width 20ns) which must not be neglected, if a timing accuracy of a few 100ps is aimed at. First an overall increase of the propagation time through a 30m long 100 Ω cable by about 4ns is observed, if the cable is rolled up (Fig. 3). An additional time shift of one pulse is observed which is caused by a second pulse propagating in an adjacent cable pair. As is seen from Fig. 3, it depends on the cable geometry and the time difference of the adjacent pulses. The time shift amounts up to 1ns in an overlap region of about ± 10 ns around the falling or rising edge of the adjacent pulse. In order to minimize this cross-talk effect, it is planned to use individually shielded twisted-pair cables for the fast digital signals between the front-end electronic units.

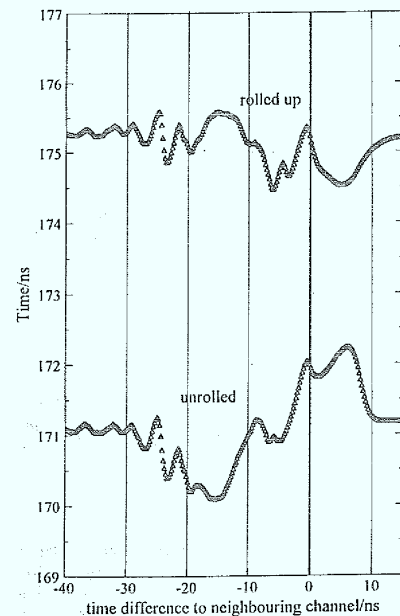


Fig. 3: The pulse propagation delay in a rolled compared to an unrolled, unshielded 100 Ω twisted-pair cable of 30m length. The structure in the delay is caused by a signal on the neighbouring channel. (negative time difference means that this neighbouring signal is earlier).

A special VME module provides the up to 11 TOF windows per stop telescope. The stop signal is directly fed into the assigned module, whereas the start signals are distributed to all units by FAN OUTs. The stop signal opens a gate which is adjustable for a time interval from 2 to 20ns in steps of 5ps according to the time window required in the TOF spectrum. By individually setting the delay in a range from 0 to 20ns for up to 16 start signals, one module provides a maximum number of 16 coincidences/TOF windows per stop counter. All gates and delays are both manually and remotely adjustable. Thus every stop counter looks individually at its corresponding start counters for a particle having the desired time of flight. The signal from the stop counter defines the timing of the trigger. Fig. 4 shows a block diagram.

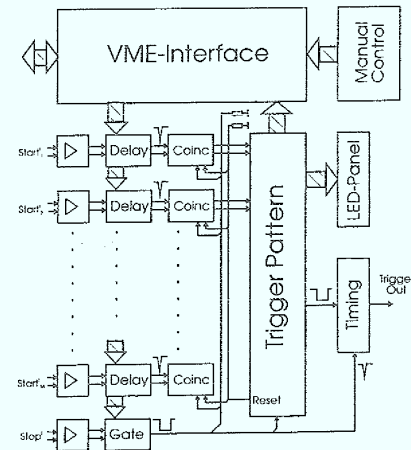


Fig. 4: Block diagram of the VME unit for the time of flight trigger. The main part consists of the 16 coincidences. All delays and the common gate width are remotely adjustable in steps of 5ps.

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- [2] A. Franzen, priv. comm.
- [3] with A. Schikov from Moscow State University and JINR Dubna.
- [4] R. Schleichert et al., IKP, KFA Jülich, Report Jül-2879 (1994), p. 48

In this paper we consider the Cerenkov counter of total internal reflection (CTIR) proposed for π/K and p/d separation and background suppression in the trigger for the detector system of the ZDF. This type of Cerenkov counters we plan to use in the experiments for the investigation of K^+ subthreshold production [1] and to study the exclusive deuteron break-up [2].

The CTIR characteristics were studied by Monte-Carlo simulation. The GEANT3 package has been used for particle tracing.

For the number of Cerenkov photons which reach the photocathode we have the well known expression:

$$\frac{dN_{ph}}{dL} = 370 \cdot \sin^2 \vartheta_C \cdot \varepsilon(\delta, E) \cdot \Delta E$$

where $\vartheta_C = \arccos(1/\beta n)$ is the opening angle of the Cerenkov cone. We assume that $\varepsilon(\delta, E)$ is the light collection (propagation) efficiency which is factorizable and uniform over E and is equal to the fraction of the light reaching the photodetector $\varepsilon(\delta, E) = \Delta\phi(\delta, \vartheta_C)/2\pi$. ΔE is the energy range of the irradiated Cerenkov photons. Using

$$E = h \frac{c}{\lambda} = \frac{1240 \text{ nm}}{\lambda} \text{ eV}$$

ΔE is equal to 1.2 eV for lucite with $n = 1.49$. δ is the particle entry angle. The acceptable angular range is equal to $(\delta_0, \delta_0 + \Delta\delta)$ with the radiator inclination angle δ_0 and the particle angular spread $\Delta\delta$.

The number of generated photoelectrons in the photomultiplier (PM) is equal to

$$N_{pe} = \frac{dN_{ph}}{dL} \cdot \frac{\Delta z}{\cos \delta} \cdot \varepsilon_{PM},$$

where ε_{PM} is the quantum efficiency of the PM (in our simulation $\varepsilon_{PM} = 0.2$) and $\Delta z/\cos \delta$ is the particle path length in the Cerenkov radiator with the thickness Δz .

The idea is to make use the total internal reflection (the angle of total internal reflection: $\vartheta_R = \arcsin(1/n)$) for the Cerenkov light collection for particles with higher speed than $\beta \cdot c$.

In the simulation program we have used quartz and lucite as a CTIR radiator. CTIR consists of a rectangular radiator $50 \times 10 \times 2 \text{ cm}^3$. The CTIR is located just behind the STOP counters. Its width of 10 cm has been defined by the TOF-system STOP-module transverse size [3]. The counter width defines the acceptable momentum range which has not affected significantly the CTIR efficiency.

To avoid wave length shifters the PM should have a quartz window to be sensitive for the ultraviolet Cerenkov light. The PM has to have good timing characteristics.

In our simulation program we have used the most unfavourable geometry: the distance of CTIR from the pointlike target was equal to 150 cm, half opening angle of

particle trajectories was taken equal to 5° . The distortion of that angle due to the inhomogenous magnetic stray field and the light absorption in the radiator were not taken into account.

The counter geometry and arrangement was described in the GEANT framework. The inclination angle for one series of simulation was fixed. The particles were generated at the pointlike target and traced to the counter straightly. The variations in particle trajectory and momentum were included in acceptable ranges. For each event we have calculated the amount of the Cerenkov light which reaches the PM. For the light fraction detected by the PM we used a simplified algorithm for the arc $\Delta\phi$ estimation in each generated event (instead of circular iterations). The 'active fraction' (or detectable arc) has been defined as

$$\Delta\phi = 2 \arccos \left(\frac{\tan(\vartheta_R - \delta)}{\tan \vartheta_C} \right);$$

if $\vartheta_R - \delta > \vartheta_C$, total internal reflections don't occur and $\Delta\phi = 0$. At last the number of calculated photoelectrons has undergone a gaussian fluctuation to take into account the collection efficiency of the first dynode.

The following factors influencing the particle detection efficiency have been taken into account:

- the particle momentum (velocity) variation for one counter in the acceptable range, because of kinematical and geometrical fluctuations and due to energy losses. This factor is not significant for one module: $\Delta p/p = 0.02 \Delta \vartheta_C / (\beta \gamma)^2 = 3\%$ ($\Delta \vartheta_C = 1^\circ$).
- variations of particle entrance angle and position due to the acceptable angular range and multiple Coulomb scattering. The factors such as the radiator optical dispersion, the light attenuation, energy losses and multiple scattering inside the radiator were not considered. It is evident that these factors are not significant (maybe except the light attenuation).

In the presented results, pulse heights from PM are expressed in the photoelectron number, but, as it is evident from the assumptions listed above, this number was estimated approximately. On the other side, they reflect the relative pulse heights correctly.

In Fig. 1 the particle detection efficiency (or the arc 'active' fraction) for various inclination angles (δ_0) is shown. Corresponding pulse heights for lucite radiator are presented in Fig. 2. From these figures it is evident that the amplitude discrimination can be used to enhance the trigger efficiency.

To summarize, we recommend to use CTIR for particle separation and background suppression. The advantages are the possibility to work at high radiation intensity (because of high refractive index) and to tune the threshold by variation of the inclination.

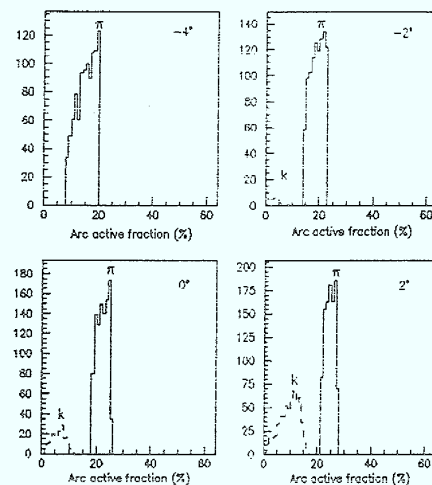


Fig. 1: Arc active fraction of 0.8 GeV/c for various inclination angles δ_0 , $N_0(\pi)=N_0(\kappa)=10^3$.

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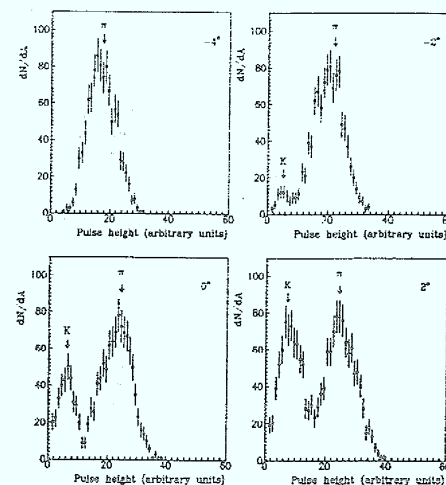


Fig. 2: Pulse heights at 0.8 GeV/c for various inclination angles δ_0 , $N_0(\pi)=N_0(\kappa)=10^3$.

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^b JINR Dubna, Moscow, Russia

The basic idea of the Cerenkov counter suggested by us earlier [1] is the following: Using the mechanism of light propagation in a radiator by means of total internal reflection one can change the detection efficiency of particles emitting the Cerenkov radiation by turning the radiator relative to the particle direction. The promising results of our simulation calculations were presented in paper [1]. In order to confirm the feasibility of the suggested mechanism, a relatively simple prototype of the Cerenkov counter was built and tested at a secondary beam of the ITEP accelerator.

The Cerenkov counter prototype was a plane radiator made of lucite with a photomultiplier (FEU-118) glued on one end of it. The voltage divider of the photomultiplier has been optimized for output signals in the linear range. The anode-signal duration was approximately 100ns. The radiator was polished on all sides except the opposite side of the photomultiplier and wrapped up in black paper to suppress diffuse reflection. The dimensions of the radiator were $36 \times 10 \times 4$ cm³. The radiator was narrowed down to 6 cm on the photomultiplier side. The diameter of latter was 4.5 cm.

The counter was placed on a horizontal frame that allowed to place it under angles in the interval from -90° to $+90^\circ$ relative to the beam direction. The angle -90° corresponds to an entrance of the particles from the free unpolished end of the radiator.

Measurements were performed on the secondary beam of the ITEP accelerator. Tuning the operating regime of the beam channel elements allowed to choose particles of the desired momentum. The intensity of the positive-particle beam with a momentum of 1.5 GeV/c containing protons and pions was $(3+5) \cdot 10^4$ for 1 s spill duration and an acceleration cycle of 5 s. The dimension of the beam at the counter was approximately 4 cm in diameter.

The photomultiplier anode pulses were recorded with an ADC (10 bit LeCroy 2249A) placed in a CAMAC crate and analyzed on-line by an IBM-PC/AT. The ADC and the data acquisition were triggered by the beam signal that was obtained from the coincidence of three scintillation counters. These counters were located on a base at distances of about 20 m and the delay selection allowed to discriminate protons and pions in the beam.

The analysis of the photomultiplier anode pulse-height spectrum proceeded mainly for the beam on a distance of about 25 cm from the photomultiplier. To check the dependence of the pulse height on the particle-entrance point a series of measurements was carried out at a distance of about 10 cm from the photocathode and no differences were found.

The investigation of the scintillation properties was made at a beam momentum 0.85 GeV/c (that is the threshold for

protons). The counter was exposed under angles -30° corresponding to π registration. The pulse-height spectra corresponding to the response from protons and π mesons are presented on the Fig. 1. It can be seen that for protons the signal is absent, so the scintillation in the radiator is not observed.

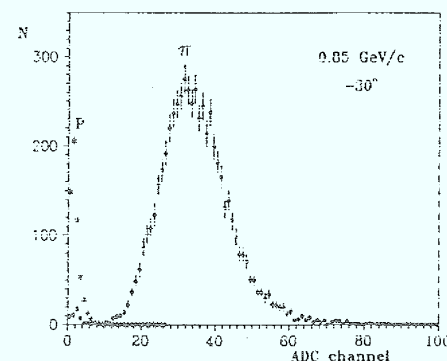


Fig. 1: Pulse-height distributions for protons and pions at 0.85 GeV/c for the positioning under -30° .

The dependence of the Cerenkov signal on the orientation angle was investigated at the beam momentum 1.5 GeV/c contained both protons ($\beta = 0.85$) and pions. (Protons with such a momentum produce Cerenkov light equivalent to that of deuterons at 3.0 GeV/c the detection of which will be of interest at the 0° Facility). At this momentum the angle of the Cerenkov radiation in the radiator with a refractive index of $n = 1.5$ is 36° for protons while for pions it is equal to 48° , and the total internal reflection angle equals to 42° .

Results of measurements for protons and pions are presented on Fig. 2. A dependence of the photomultiplier anode pulse height on the angle relative to the beam direction is clearly observed. At rotation towards negative angles the pulse height from π mesons rises earlier than that for protons. As mentioned above the cone of Cerenkov radiation for protons is narrower, therefore the increase of the amplitude for protons is sharper. At the perpendicular positioning of the radiator a decrease of amplitude is caused in the same way - the signal is well extinguished by multiple reflection. At positive angles a decrease of the number of reflections occurs including also the reflections from the free side of the radiator. That causes some increase of the signal. For pions the total internal reflection takes place even at the perpendicular positioning of the radiator and an increase of the amplitude at negative angles is more slightly. When both protons and pions induce Cerenkov light, the possibility of discrimination is well demonstrated on Fig. 3 corresponding to angle -10° .

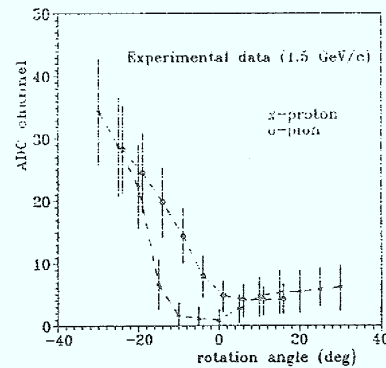


Fig. 2: Measured pulse heights for protons and pions vs. the rotation angle at 1.5 GeV/c.

During the counter testing it has become obvious that the counter characteristics can be improved by some further developments. For instance an utilisation of the wedge-shaped radiator (adapted to the photomultiplier) allows to decrease the number of reflections from the radiator's free end improving the light-reflection conditions near the photomultiplier simultaneously. Selection of the photomultiplier and optimisation of operating conditions demand further investigations (the transparency of the radiator has not been studied). It will be useful to study the different connection conditions of the photomultiplier to the radiator, for instance, the inclination of PM relatively to the axis of the radiator. The utilisation of quartz as radiator can also improve the discriminating properties of the suggested Cerenkov counter because of the sensitivity of quartz in the ultra-violet region.

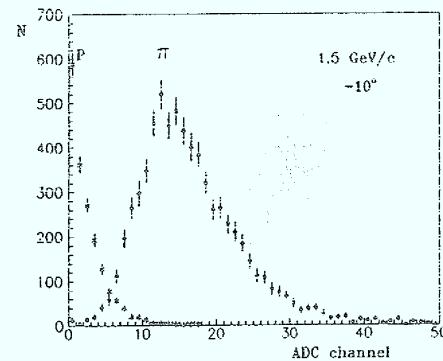


Fig. 3: Pulse-height distributions for protons and pions at 1.5 GeV/c, corresponding to the rotation angle -10° on fig. 2.

The performed studies of the Cerenkov-counter prototype allow us to believe that such detectors can successfully be used for the separation of protons and deuterons in the FORWARD arm [2] of the ZDF setup (momentum region 1.4 ± 2.0 GeV/c) by means of a careful selection and the optimisation of its parameters. Because of the limited space for the location of the apparatus in the FORWARD arm, the suggested Cerenkov counter would be preferable over Silica Aerogel counters. The necessary location of counters at very short distance from the beam-pipe requires a careful estimation of the necessary volumes to decrease the background load.

The authors would like to thank Dr. I Minashvili for helpful discussions. The research group from Tbilisi would like to note that the research described in this publication was supported in part by Grant No. Ph1-8678-0925 from the ISF.

- a HEPITU Tbilisi, Georgia; supported by the International Science Foundation
- b ITEP Moskau, GUS
- c JINT Moskau, GUS

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The 0° Facility which is currently being built [1] will allow to simultaneously measure positively and negatively charged ejectiles which leave an internal target under forward angles $|\theta| \leq 15^\circ$. It has been proposed [2] to measure the K^- production in $p+A$ reactions at energies close to and below the nucleon-nucleon threshold at $T_p = 2.5$ GeV. In this article we describe a concept of a detection system for the proposed studies.

For the detection of the K^+ (K^-) mesons and for the reduction of the background consisting mainly of π^+ and p (π^-) and of particles scattered in the magnet iron the following concept is suggested: The K^+ mesons will be detected in the side-detector system [3] which will offer a K^+ -detection efficiency of $\sim 15\%$ (including decay-in-flight between target and detectors) and a background suppression by more than 4 orders of magnitude [3,4]. It covers a momentum range of $p_{K^+} \sim 150 \dots 600$ MeV/c and allows to select events where a K^+ meson is produced with almost no background from other reaction channels.

As shown in Fig. 2 of [2] roughly the same number of K^- and π^- are produced in events where a K^+ meson is formed. Thus, only a moderate background suppression has to be achieved by the detection system for the negatively charged particles.

The principal layout of the $K^-(\pi^-)$ detection system can be seen in Fig. 1. The TOF-start counter wall can be similar to the one used for positively charged ejectiles, i.e. it consists of 18 scintillators with a height of ~ 25 cm, a width of 5 cm and a thickness of $1 \dots 2$ mm. For the TOF-stop counter wall 20 'standard' scintillators of 10 cm width and 1 cm thickness can be chosen. This wall allows to detect ejectiles in the momentum range $p = 150 \dots 1000$ MeV/c which is equal to the full momentum range accessible for negatively charged particles at the 0° Facility. According to raytracing calculations the detector heights must be between 40 cm ($p = 150$ MeV/c) and 80 cm ($p \geq 500$ MeV/c) in order to cover the full vertical angular acceptance given by the spectrometer magnet.

For further background suppression it is planned to also read out the ΔE signal from the TOF-stop counters, see below. Two wire chambers will allow to measure the particle flight paths for track reconstruction and for suppression of background particles scattered in the pole pieces of the spectrometer dipole, see also [3].

The performance of the full detection system has been simulated using 'realistic' K^+/K^- and K^0/π^0 events calculated with the Rossendorf Collision Model [2]. Only events where a K^+ meson is detected and a negatively charged ejectile reaches a TOF-stop counter were analyzed. Decay-in-flight and hadronic interaction in the magnet and detectors have not been simulated yet. Figure 2 shows a 2-dimensional plot of the TOF of $K^-(\pi^-)$ versus the time difference $\Delta t_{+/-}$ between the detection of the K^+ and the K^-

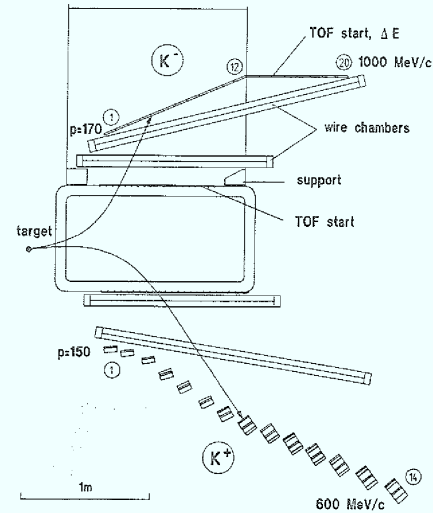


Fig. 1: Principle layout of the proposed detection system for the K^+/K^- studies. For the detection of the K^+ mesons the system described in [3] will be used. It consists of a TOF-start counter wall and 14 telescopes with TOF-stop, Cerenkov, ΔE and veto counters. Negatively charged particles will be detected in a similar TOF system. For further background suppression the ΔE signal from the TOF-stop counters can be used.

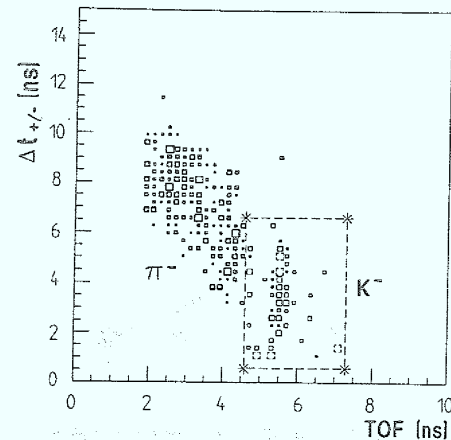


Fig. 2: Simulated TOF spectra of correlated K^+/K^- and K^0/π^0 events summed over all counters.

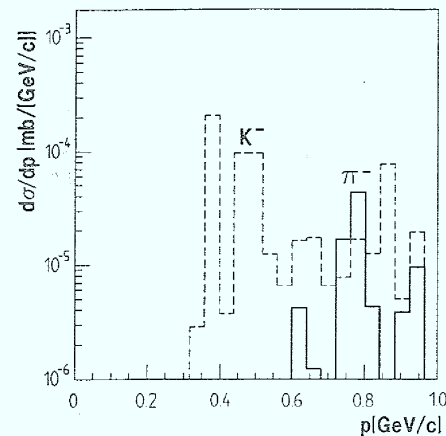


Fig. 3: Momentum spectra of particles lying inside the gates shown in 2.

(π^-) in the corresponding TOF-stop counters. It can be seen that - even if the sum over all counters (and, thus, ejectile momenta) is used - there is a good separation between the K^- and π^- . This allows to select K^- events by setting a 2-dimensional gate in the spectrum, see Fig. 2. The resulting momentum spectra of particles lying inside this gate is shown in Fig. 3. A comparison with Fig. 2 from [2] shows a clear suppression of the π^- events. Only for high-momentum ejectiles ($p \sim 800$ MeV/c) there is a remaining π^- background. However, this background can also be totally rejected (on the level of statistics in our simulations) if one analyzes the TOF informations from Fig. 2 for individual stop counters. As an example, Fig. 4 shows the same spectrum as Fig. 2 when only scintillator #18 of the negative stop-counter wall ($p \sim 800$ MeV/c) is read out. Clearly, the separation between K^- and π^- is much better and the π^- events can be rejected completely.

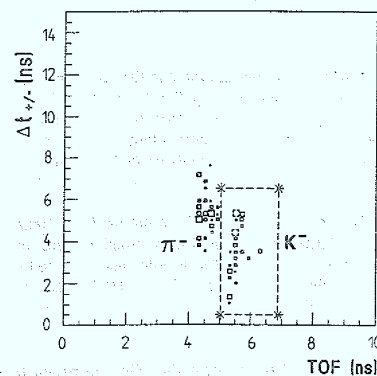


Fig. 4: Same as Fig. 2 but only for counter #18 of the stop-counter wall for negatively charged particles.

From Fig. 3 one can estimate the counting rates for correlated K^+/K^- events. Taking into account a total detection probability of 15% [3] (30%) for K^+ (K^-) we expect 25 events per minute integrated over all momenta ($p_{K^+} = 150 \dots 600$ MeV/c, $p_{K^-} = 170 \dots 1000$ MeV/c).

In addition to background suppression via TOF measurements it will be possible - especially for momenta $p \lesssim 600$ MeV/c - to use the ΔE signal from the TOF-stop wall to distinguish between K^- and π^- . As an example Fig. 5 shows the ΔE spectrum in counter #10 ($p \sim 400$ MeV/c) for K^- and π^- . From this spectrum it can be concluded that an additional π^- suppression of at least 1 - 2 orders of magnitude for lower momenta will be possible.

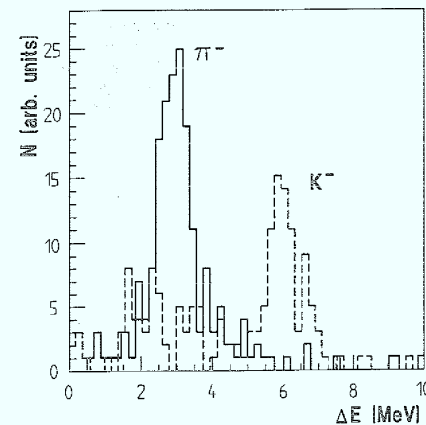


Fig. 5: Energy-loss spectrum for K^- and π^- in TOF-stop counter #10 ($p \sim 400$ MeV/c).

Summarizing, the detection system presented here will allow to measure correlated K^+/K^- events from p+A interactions at $T_p = 2.5$ GeV and even at lower beam energies where the cross sections decrease rapidly (see [2] and reference therein). It is currently under investigation if the system can be installed 'step by step' using first only the TOF-stop wall for K^- identification. In this case one would use the information on the time difference between emission at the target (calculated off line using the information on the K^+) and detection in the wall. This simpler setup should allow to measure at $T_p \approx 2.5$ GeV and moderate cross sections. For further measurements at lower energies, the full setup would be required.

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The 0° Facility is designed as a multi-purpose device for experiments at internal targets in COSY-Jülich and will consist of three dipole magnets [1]. For the experiments it will be necessary to ramp the dipoles with various speeds up to ~ 1 T/s, depending on the maximum field in the 0° Facility and the COSY dipoles as well as the ramping speed in COSY. The basic parameters of the power supplies for the three magnets (i.e. maximum output voltage and current, maximum ramping speed dI/dt) are, to a large extent, determined by these experimental requirements which we have investigated for the various envisaged measurements.

It should be noted that for the smaller dipoles D1 and D3 an existing power supply will be used which offers a maximum output voltage of $U_{\max} \approx 75$ V and a maximum ramping speed of $dI/dt_{\max} \approx 2500$ A/s. Obviously, the power supply for the large spectrometer dipole D2 should be designed such that these given parameters can be utilized in the best possible manner.

Figure 1 shows the time dependence of the magnetic field in the dipoles D1/3 and D2 of the 0° Facility for the proposed subthreshold K^+ studies with a flat-top field of 1.3 T. Note that the field strength in all three dipoles must be equally high in order to guarantee a stable closed orbit bump of the circulating COSY beam. The flat-top beam energy at which the measurements will be made is in the range 0.9 ... 1.2 GeV and the radius of curvature of the circulating beam in the 0°-Facility magnets will be $\rho = 4.05 \dots 4.92$ m. The assumed ramping speed in COSY is $dB/dt = 0.9$ T/s which is close to the maximum possible ramping speed in COSY.

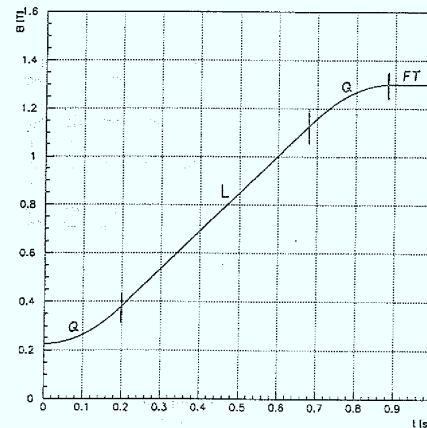


Fig. 1: Magnetic field in the dipoles of the 0° Facility during ramping and flat-top for a mode of operation with $B_{\text{flat-top}} = 1.3$ T. As in the case of the COSY dipoles the field ramp consists of two quadratic roundings (Q) during 0.2 s, a linear increase of the field (L) and a constant flat-top value (FT).

For the calculation of the maximum ramping speed $dI/dt_{\max}$ and voltages $U_{\max}$ the values given in Table 1 for D1/3 and D2 were used as well as the $B(I)$ and $L(I)$ dependences from Fig. 2a,b. The latter values were obtained from 2d and 3d field calculations using the POISSON and PROF1 program packages [2]. For the calculation of the inductivity the formula $L = I^{-1} \cdot dW/dI$, where the current I and the energy W stored in the magnets were taken from the calculations.

Table 1: Basic parameters of D1/D3 and D2 used in our calculation.

	D1/3	D2
Number of turns per coil	24	120
Resistance of one coil [mΩ]	2.6	56.75
Current at $B=1.58$ T [A]	2910	1375

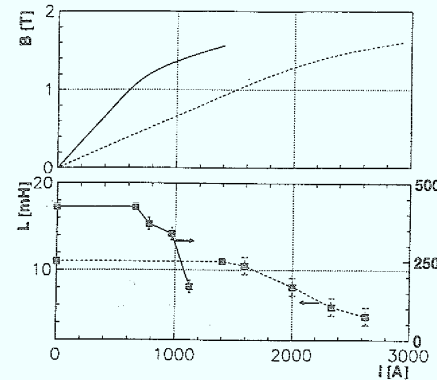


Fig. 2: $B(I)$ and $L(I)$ dependence of the dipoles D1/3 (dashed lines) and D2 (solid line).

Figure 3 shows the values of dI/dt and the necessary output voltage U for the two power supplies during the field ramp depicted in Fig. 1. It can be seen that the maximum values achievable with the existing power supply for D1/3 are not exceeded. For the D2-power supply we deduce a maximum ramping speed of $dI/dt_{\max} \approx 1300$ A/s and $U_{\max} \approx 550$ V.

For the K^+ mode of operation with a flat-top field strength of 1.58 T it turned out that the COSY ramping speed must be reduced to $dB/dt \approx 0.4$ T/s in order not to exceed the maximum values for dI/dt and U for D1/3. The main parameters for the various modes of operations are listed in Table 2.

Summarizing, we have deduced the following parameters for the D2 power supply: a maximum current (flat-top) of $I_{\max} \approx 1375$ A, a maximum ramping speed of $dI/dt_{\max} \approx 1300$ A/s and a maximum output voltage of $U_{\max} \approx 550$ V.

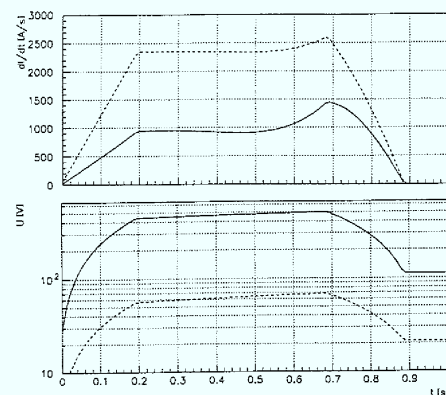


Fig. 3: Ramping speed dI/dt and required output voltage U of the power supplies for D1/3 (dashed lines) and D2 (solid) for the field ramp from Fig. 1.

Table 2: Main parameters for the power supplies for D1/3 (first number) and D2 (second) for the various modes of operation. The values for dB/dt are given for the linear part of the field ramp.

	K_1^+	K_2^+	d breakup	η production
$B_{flat-top}$ [T]	1.3	1.58	1.58	0.62
dB/dt_{COSY} [T/s]	0.9	0.4	0.7	1.0
$dB/dt_{D1/23}$ [T/s]	1.56	0.69	0.7	0.55
dI/dt_{max} [A/s]	2500/1300	2500/1300	2500/1300	850/400
U_{max} [V]	75/550	50/350	50/350	30/200

A power supply with these parameters will allow to make optimum use of the existing power supply for D1/3 and offer a high flexibility for the various modes of operation at the 0° Facility.

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Status of the Experiment
Proton Therapy at COSY, Pre-Therapeutic Physical Experiments on Phantoms with Proton Beams
at Energies between 50 MeV and 250 MeV

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Orienting measurements with the COSY beam in the KFA water phantom

First test measurements were performed at the low-energy-experimental area (Niederenergie-Meß-Platz, NEMP). The energy of the extracted proton beam was 295 MeV. A vessel with dimensions of 52 cm x 52 cm x 45 cm filled with water was used as a phantom. The proton Bragg peak occurred at about 51 cm. The beam monitoring was done by 300 μ m thick plastic scintillators.

For the measurements of dose profiles and microdosimetric spectra, a 1/2" Rossi counter was used, which was mounted on a position scanner. The detector is a proportional counter with a spherical sensitive detection volume made of tissue equivalent plastic A150, which has an inner diameter of 1.27 cm. The detector is filled with methane based tissue equivalent gas. The pressure of the gas was set to 75 mbar, so that the detector simulates a 1 μ m sphere of unit density (tissue) material. The spectrum of secondary charged particles created in wall as well as the energy deposition of the primary protons and of the secondary charged particles in the sensitive counter volume is therefore very similar to what would be measured in a spherical tissue cell of 1 μ m diameter.

Tissue equivalent proportional counters (TEPC) measure the energy deposition of single charged particles crossing the sensitive volume. The resulting frequency distributions are scaled in terms of lineal energy, y , which is defined as the ratio of energy imparted by a single charged particle and mean chord length of the detection volume. The product of the frequency mean lineal energy (first moment of the measured frequency distribution) and the total number of events detected is directly proportional to the dose deposited in tissue at the measurement location.

Figure 1 shows dose distributions (energy weighted frequency distributions) in terms of the logarithm of lineal energy. The maxima of these distributions shift to higher lineal energies towards the Bragg peak.

Pile up effects caused by the irregular distribution of protons over a COSY pulse impaired the measurements. The following experiments will have to account for this situation. Additionally a better spatial resolution will be achieved, if a smaller proportional counter is used. An improvement of the signal to noise ratio can provide measurements at lower values of lineal energy. Finally a better statistic is necessary to interpret the measurements at the edges of the proton beam. The precision of this first test measurement were clearly effected by pile-up effects, misalignment and width of the proton beam. In addition the beam energy exceeded the range of application.

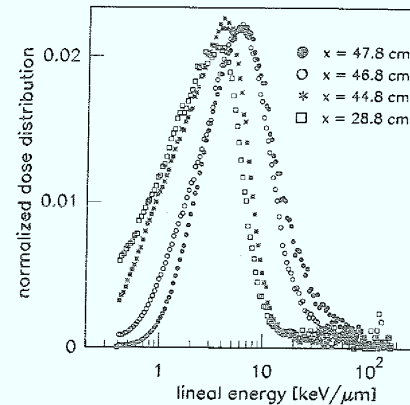


Fig. 1: Microdosimetric dose distributions as a function of depth in water

We planned to gate the TEPC data acquisition system by the scintillator veto. However, timing problems and pile up effects haven't allowed us to use this feature yet.

Comparison of different Monte Carlo programs for use in proton dosimetry

Before a therapy treatment planning system based on the Monte Carlo method can be developed, different codes have to be studied to find out about their different features and their advantages in the energy region of 70 - 250 MeV. We used four different codes to calculate the depth dose profile of protons in water and compared them with recent data from PSI, Villingen, Switzerland.

- Characterization of the used Monte Carlo programs:

1. HETC

HETC (High Energy Transport Code) is a module of HERMES /1/ (High Energy Radiation Monte Carlo Elaborate System) which was extended at KFA. Its main applications are radiation transport and detector simulation in medium and high energy physics. Apart from charged particle ionization energy loss based on the Bethe-Bloch formula it generates and transports secondary particles like Neutrons, protons and recoil nuclei. For primary particles small angle Coulomb scattering is performed. Using the Combinatorial Geometry, one can look at energy deposition, yield and flux of primary and secondary particles at any given point of the geometrical set-up.

2. MC4 /2/

MC4 is a Monte Carlo detector simulation code, developed at KFA-IKP for particle detectors in medium and high energy physics. It is also integrated in the HERMES system and uses the same nuclear collision model as HETC. The main difference concerning the use in proton dosimetry, is the different handling of range fluctuations. HETC performs range straggling at the end of the proton flight path when the particle reaches the cut-off energy (1 MeV). In MC4, energy loss along the track is optionally straggled with a Gaussian or a Landau distribution.

3. TRIM /3/

TRIM (The Transport of Ions in Matter) calculates the penetration of ions in solids for initial energies from 10 eV to 2 GeV. It is mainly used to calculate the final distribution of the ions, target damage, ionization and phonon production. Secondary particle production is not included.

4. PTRAN /4/

PTRAN (Proton Monte Carlo Transport Program) is especially designed for use in treatment planning and dosimetry in proton therapy. It deals with the penetration of 50- to 250-MeV proton beams through water phantoms. Energy loss straggling in Coulomb interactions is sampled from the Vavilov distribution. Multiple-scattering is treated according to Moliere's theory. Protons in nonelastic nuclear collisions are absorbed, their energy is deposited locally. Secondary particles are not produced. The program simulates monoenergetic pencil beams. It calculates depth dose profiles, lateral dose profiles and proton spectra as functions of depth.

- Fit of Experimental Data from PSI

With these four Monte Carlo programs mentioned above, we tried to fit experimental data from PSI, Switzerland /5/. Energy and energy spread of this bragg peak were not determined experimentally but estimated by PSI. To get in good agreement with the position of the bragg peak of the PSI-data at 13.7 cm, we had to run the MC programs with different proton energies ranging from 137.0 to 139.5 MeV (see Fig. 2). The peak width could be approximated only with HET and MC4 by differing the energy spread. The other two codes handle only with monoenergetic pencil beams and therefore have narrower peaks.

The estimated energy and energy spread of the proton beam depends strongly on the Monte Carlo program, because they are using different energy-range-tables and secondary particle production and transport is not always included. For proton dosimetry purposes, the development of a new code is therefore recommended. It should be able to deal with protons up to 300 MeV and to produce secondary particles like neutrons, light ions and also delta electrons, to allow to study their influence on the dose distribution.

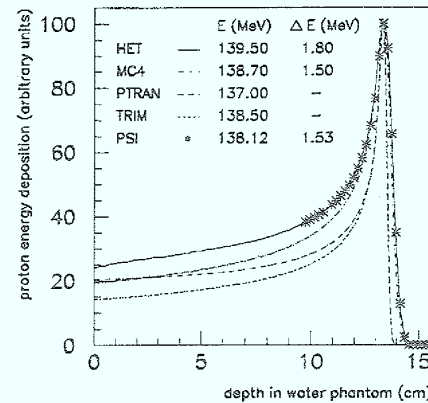


Fig. 2: Fit of PSI-data with different Monte Carlo codes

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Status and first Tests of highly granulated Detector Systems for Pre-Therapeutic Physical Experiments at COSY

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For the pre-therapeutic measurements at COSY [1,2] modern detector systems used in fundamental research experiments will be applied.

Experimental Setup

For different beam momenta the spatial flux distribution of primary and secondary particles after penetrating certain absorbers will be measured. That is done by highly granulated detector arrays which enable the observation of a few thousand pixels in the beam-phantom-system at the same time. For that silicon microstrip detectors of different geometric shapes and scintillating fibre hodoscopes are built. Silicon microstrip detectors offers the possibility to determine the differential energy deposition simultaneously with high spatial and energy resolution. The application of scintillating fibres guarantees the quasi tissue equivalence of the detector system.

Fig.1 shows the basic setup using a solid phantom wedge. Every incident beam particle is localized by a startdetector (A). To determine the flux distribution the main detectors (C) are placed between the phantom wedge (B) and a following tissue equivalent absorber (D) which allows the observation of backscattering. Replacing the solid phantom wedge by a water basin one can study the recent effects of protons on a liquid phantom. Beyond that the design of the scintillating fibre hodoscope allows also the application inside the liquid.

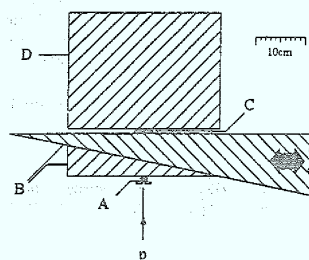


Fig.1: Basic experimental setup with a solid phantom wedge.

Detector design and status

The startdetector system consists of two scintillator counters and a pair of silicon microstrip detectors. The scintillators serve the timing signal and a clear beam definition. With the microstrip detectors (active area $3.2 \times 3.2 \text{ mm}^2$, spatial resolution $200 \mu\text{m}$) the position of

each incident particle is measured to define a "micro-beam".

The "main detector" for the measurement of the proton propagation in the phantom can be equipped with three different detector arrays:

- A pair of two crossed linear silicon microstrip detectors ($2 \times 2 \text{ cm}^2$ active area, $200 \mu\text{m}$ resolution) for a high precision position measurement
- A silicon ring detector (4 inch $\varnothing$, $300 \mu\text{m}$ spatial resolution) for investigating the outer region of the diffracted beam
- A fibre hodoscope ($3.2 \times 3.2 \text{ cm}^2$ active area, 1 or 2 mm resolution) for applications in water or behind the phantom wedge

The startdetector system is built and tested in different environments. Prototypes of each main detector are built and tested. Optimized main detectors are designed and under construction.

Test of the experimental principle

During a test measurement at Big Karl with a simplified solid phantom (consisting of several plexiglass blocks) and an proton energy of 200 MeV the perfect behaviour of detectors and data acquisition system could be shown. As expected one sees the broadening of the "beam profile" (fig.2a) measured in the microstrip array and also an increasing energy loss (fig.2b) with increasing phantom thickness. [3]

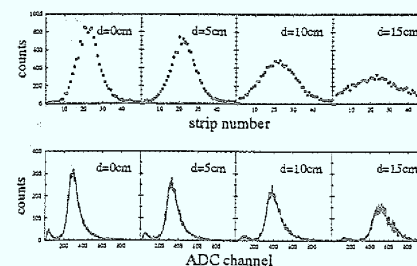


Fig.2: a) "beam profile" and b) energy loss spectra measured with microstrips behind a phantom of thickness d .

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Thermoluminescence materials are widely used in conventional radiotherapy because of their high sensitivity, small size, self integrating behavior, fast analysis, linear response in the therapeutic dose range and the possibility to use them as in vivo dosimeters.

In heavy particle fields the dose response is strongly LET dependent for most materials; it was shown however /1,2/, that the sensitivity of the high temperature peak of $\text{CaF}_2:\text{Tm}$ (TLD 300) is LET independent up to 50 keV/ μ , whereas the sensitivity of the low temperature peak decreases rapidly with increasing LET. In some high LET fields like neg. pion or helium fields it was demonstrated, that it is possible to determine dose and average lineal energy simultaneously from the analysis of a single glow curve /3/. With a proper RBE vs. LET model this allows to determine biological equivalent dose. In therapeutic neutron fields it is possible to separate neutron and gamma-dose /4/.

It is therefore planned to determine the potential of this dosimetry method at the pretherapeutic 200 MeV proton beam of the COSY accelerator at KFA Jülich. The glow peaks of different materials will be compared to ion chamber dose measurements and to the results of the microdosimetric measurements taken with proportional counters. After this calibration the use of TLD's allows to measure a complete 3-dimensional dose distribution within one run of the treatment plan program. Because of their high sensitivity they can be produced in thin foiles, which do not change the range of the protons measurably and therefore will not disturb the dose distribution to be measured.

In 1994 two TLD readout apparatus (Harshaw 200 A+B) were installed with proper annealing and handling equipments. One is located at KFA for fast readout after irradiation and one at Wuppertal University for further analysis. The TL and temperature signals are interfaced to a PC, that reads out the values every 80 msec. The readout of a single TLD takes 60 sec and all glow curves are stored on disk for later reference. The first version of an analysing software has been developed. The handling procedure still has to be optimized in order to achieve a reproducibility of better than $\pm 3\%$.

For the first irradiations in the proton beam 45 $\text{CaF}_2:\text{Tm}$ (TLD 300) chips of size $3 \times 3 \times 0.9$ mm have been selected. They were individually calibrated at the Co^{60} -source of the Institute for Medicine (IME) at KFA. The calibration dose was chosen to be 0.25 Gy at a dose rate of 0.5 Gy/min. Figure 1 shows a typical glowcurve readout after Co^{60} irradiation. The sensitivity of P_1 is about 400 nA/Gy and the P_1/P_2 -ratio is 3. The peak P_0 at 100 °C is unstable at room temperature and fades away in about one day. It can be used to determine the time between irradiation and readout.

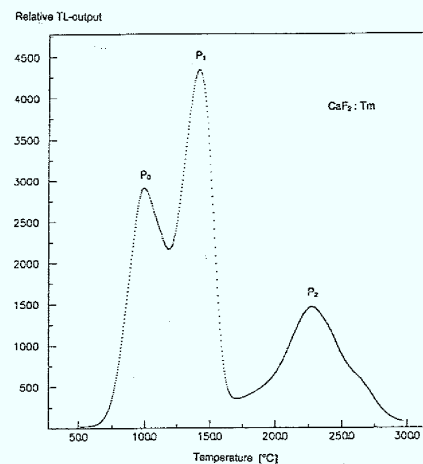


Fig. 1: Glow curve of $\text{CaF}_2:\text{Tm}$ thermoluminescent material after Co^{60} -irradiation

At the proton beam line 2 TLD holders were constructed, that fit into the $50 \times 50 \times 40$ cm³ water phantom built by KFA. An array of 7×7 TLD chips can be positioned perpendicular to the proton beam axis at a grid size of 1.5 cm. The positioning accuracy is 0.1 mm and the extra plastic material in the beam is 2 mm water equivalent. The center position is located at the same water depth as the 0.1 cc ionization chamber (IC 18), such that with the same phantom setting the results can immediately be compared.

In early 1995, as soon as proton beam is available, the depth dose distribution in water will be measured with the 0.1 cc chamber and with TLD chips. The glow curves will be analysed to find any indication for high LET radiation components in particular in the bragg peak falloff region. The results then will be compared with the corresponding microdosimetric measurements.

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Calculation of the Biological Effectiveness for Proton Beams

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The use of protons in cancer radiotherapy offers significant advantages over radiations conventionally employed. Tumor volumes in the body can be irradiated more precisely than it is possible with photon, electron, pion or neutron beams. The biological effectiveness of therapeutic pencil proton beams was studied with regard to the initial energy distribution of an extracted beam and to the depth of the beam penetration in water. Such data will be useful to assess the biological effectiveness of protons in future therapy installations with electronically modulated scanned beams.

It is possible to characterize the biological effectiveness of a given radiation by a determination of its microdosimetric parameters. A main parameter used in experimental microdosimetry, lineal energy, y , describes an actual energy deposition to the target normalized to its mean chord length. Even if the whole system is irradiated with just a single radiation modality the target can experience a whole spectrum of energy deposition events. These spectra of energy deposition events (microdosimetric distributions), denoted e.g. $d(y)$, and their moments are used to describe the ionization density of the radiation in the target. The assessment of RBE is based on an assumption that there exists a biological weighting function, $r(y)$, which relates a given, observed biological effect to measured or calculated $d(y)$ distributions. This distribution can be obtained, when the proton energy spectrum at the point of interest is known.

Proton transport calculations in water for a proton energy range from 70 to 250 MeV were performed with the Monte Carlo codes PTRAN /1/ and MC4 /2/ in order to obtain proton energy spectra at a given depth in water. If one replaces tissue by water and neglects differences in the electron collision cross sections, the interaction of protons can be simulated by using cross sections for water vapor. The momentum spread of an input semi-monoenergetic

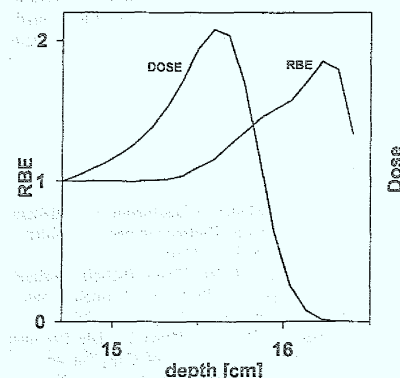


Fig. 1: RBE and dose as a function of depth for a 150 MeV pencil beam. The dose is given in arbitrary units.

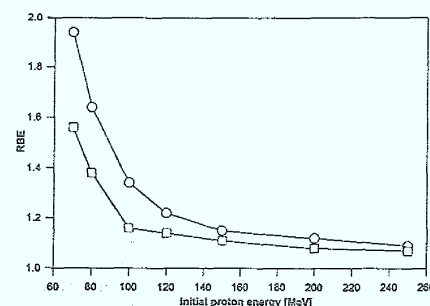


Fig. 2: RBE in the center of the Bragg peak (maximum dose) (O) and averaged over FWHM (□) as a function of initial proton energy for a monoenergetic proton beam.

proton beam, $\Delta p/p$, was varying from 0 to 0.003. Next, the microdosimetric response function approach was applied to calculate the biological effectiveness of a proton beam at a depth of interest with the use of a biological weighting function for early intestinal tolerance in mice published by Loncol et al. /3/. The biological response function describes the probability of a biological effect after an event of energy deposition with a given lineal energy, i.e. it describes for a biological endpoint the cumulative probability that a subcellular target will respond to a specific target-averaged ionization density. In the calculation of the relative biological effectiveness for protons, RBE, the whole distribution of possible energy deposition was taken into account and normalized to an effect for Co-60 gamma-rays. A computer code PMIC has been prepared to analyze the results of the Monte Carlo calculations in order to obtain microdosimetric characteristics of the proton beam such as e.g. the $y_d(y)$ distributions and their moments. Homogeneous spherical targets of a uniform sensitivity, 1 μm in diameter, were considered in our calculations.

The results show that RBE values depend on the initial energy spread of the input proton beam and on the depth of beam penetration. RBE for intestinal tolerance in mice approaches 1 in the plateau region and gradually increases with the proton penetration depth. In the center of the Bragg peak, at maximum dose delivery, RBE values range from about 1.1 (250 MeV beam) up to 1.9 (70 MeV beam). The rise of the RBE decreases with initial proton energy. Averaged over FWHM of the Bragg peak we found 1.1 and 1.6, respectively. At the end of the Bragg peak the RBE was found even higher (1.3 for 250 MeV beam to about 2.3 for 70 MeV beam) but only about 1% of dose is delivered with such a high RBE values. Figure 1 illustrates the depth dependence of the RBE in the region of the Bragg peak for a 150 MeV pencil beam. The RBE in the center of the Bragg peak and peak averaged is illustrated in figure 2 as a

function of incident proton energy. The calculated RBE values are only slightly depending on the energy spread of the proton beam for the small spreads used. When increasing the energy spread, we obtained a decrease of the RBE behind the Bragg peak and a slightly RBE increase in the Bragg peak. The increase of the RBE with depth results in a shift of the biological dose distribution (RBE times physical dose) to larger depths in comparison to the physical dose distribution, e.g. about 1mm for the 150 MeV beam. By this, using an average RBE results in an underestimation of the effectiveness of the beam at the end of its range. In addition an radial increase of the RBE has been found. The above results indicate that for therapy planning with electronically modulated pencil proton beams it would be necessary to account for a depth and radial dependence of proton RBE, especially for low penetrating beams.

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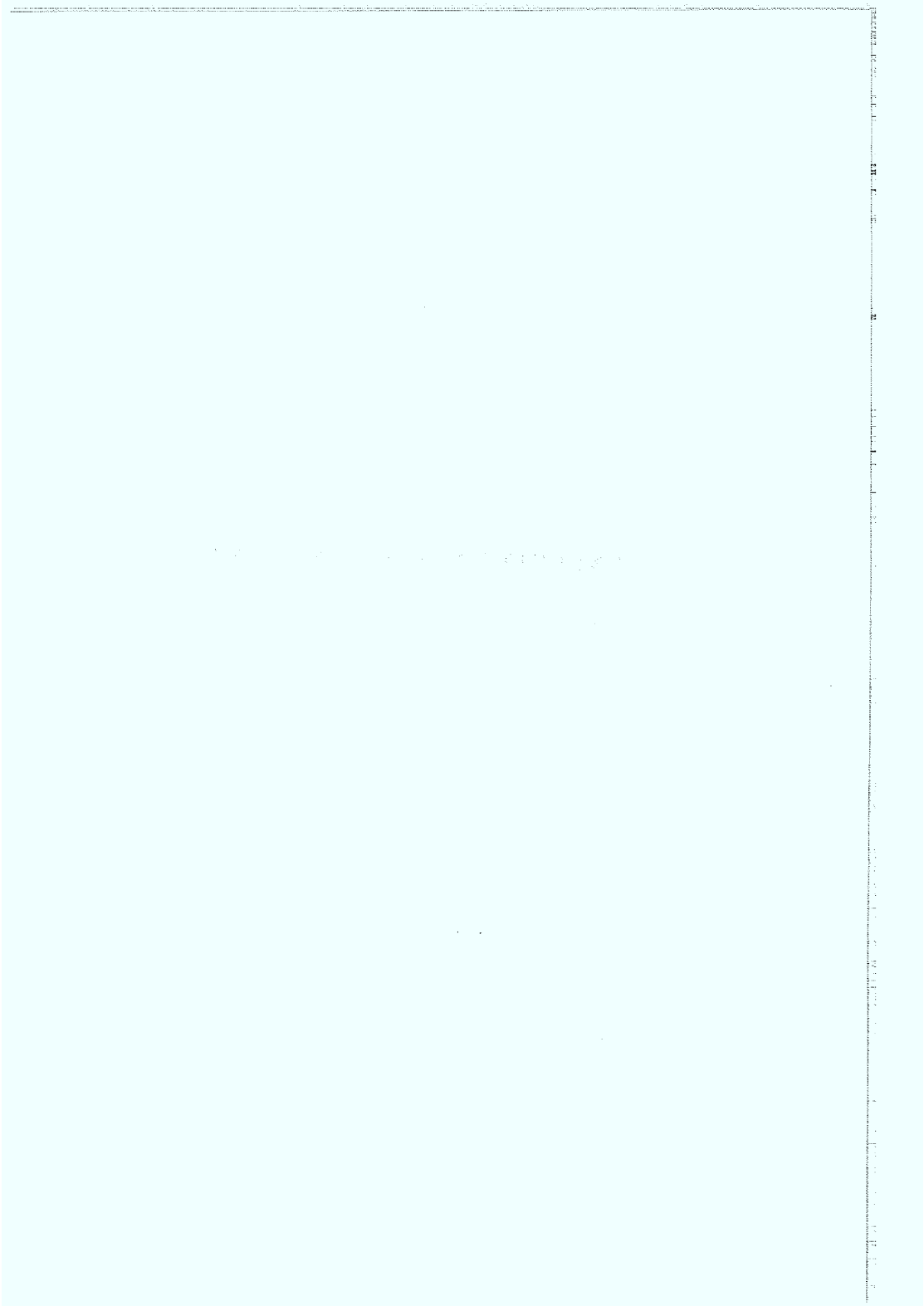
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1.2 EXPERIMENTS AT EXTERNAL FACILITIES



Development of an Aerogel Cerenkov Detector with Wavelengthshifter Lightguide Read-Out for the Observation of Single Kaon and Double Kaon Production at SAPHIR

Status Report

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For elementary particle physics experiments at intermediate energies -e.g. SAPHIR [1]- π -K-separation can be performed by aerogel threshold Cerenkov detectors [2]. In that special case the possibility of an aerogel Cerenkov detector with wavelengthshifter lightguide (WSL) read-out is under investigation [3].

Test detector

The test detector is assembled and will be tested at the Bonn 2.5 GeV-Synchrotron test beam facility in early 1995. A related read-out electronic circuitry was developed consisting of a three channel ECL discriminator and a micro controller for remote control via a RS-232 serial interface. A graphic user interface allows one to set the discriminator thresholds interactively.

Radiator improvement

To increase the light output of a Cerenkov detector with read-out by WLS, the spectral matching between radiator and the WLS can be done by doping the radiator itself with a fluorescent dye as wavelengthshifter.

The possibility of doping the radiator with a fluorescent dye, in our case aerogel, was investigated during a stay of M. Bockhorst at the Université Claude Bernard, Lyon, France. The optical properties of the samples show that it is possible to dope aerogels with the dye POPOP if the hypercritical drying is done via solvent exchange. Both the concentration of the POPOP and the homogeneity are good enough for the use as radiator material [4] and an enlargement of the light output by a factor 2 seems possible.

Alternatively investigations for the improvement of the WLS geometry and spectral properties were made, which show the possibility of an enhancement of the light output by a factor 4 without doping of the radiator.

Detector Simulation with GEANT

The aerogel Cerenkov detector (ACD) was integrated into the SAPHIR-GEANT detector simulation software as instrument for the designing of the Cerenkov detector geometry. The first results of the simulations show that WLS-hits occur but a careful design reduces them below 10 % for Kaons produced in different reactions e.g. $\gamma p \rightarrow K^+ A$. Figure 1 shows the calculated ADC values of the ACD of all particles in a tagged photon energy range of 1.4-2.9 GeV. The refractive index of the aerogel was assumed to be 1.04.

Rates of Double Kaon Production

Using the following data the daily rate of double K^+ events for the reaction $\gamma D \rightarrow HK^+ K^+$ ($H = \text{Dibaryon}$, $S = -2$) can be estimated:

- Tagging rates with the TOPAS II tagging system
- Hard- and software acceptance for single K production
- The estimated cross section

The planned maximum tagging rate is 10×10^6 Hz, a rate of 10^6 Hz can be used now and is assumed. The acceptance for single K production is about 0.1 [5] using the kinematical fit without any TOF identification, for KK production it will be in the same range. The estimated cross section for the above described reaction is ca. 2.5 nb [6]. Using a D_2 target of 8 cm thickness a reaction rate of 3 identified double K^+ events per day is expected.

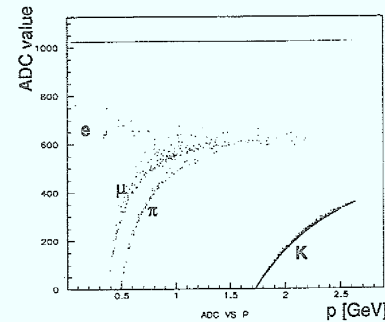


Figure 1: Simulated ADC values

Future tasks

- Simulation of different reactions with strangeness (kaon) production for the determination of the hardware- and software acceptance of SAPHIR. All these simulations will be done with the Cerenkov detector.
- Evaluation of physics data of the SAPHIR detector to improve the particle identification mechanisms.
- The Cerenkov detector will be build in modules. One module with dimensions appropriate for the SAPHIR will be built after successful measurements with the test detector.

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Further Progress in Offline-Analysis of PS185

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for PS185 collaboration

The experiment PS185 measures charged tracks to analyse the reaction $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$. At beam momenta of 1992 MeV/c the processes $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\gamma$, $\bar{p}p \rightarrow \bar{\Lambda}\Sigma + cc$ and $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\pi^0$ are also energetically possible. These reactions contain two vertices of p, π^- and $\bar{p}, \pi^+$ and in addition an unobserved neutral particle. The only way to distinguish between those four channels is via missing mass and Dalitz plot. For this analysis a clear event pattern is essential.

To improve the resolution of the detector-system in the offline analysis a first step was made by a reconstruction of driftchamber hits via isochrones [1]. In the further development a new algorithm for three-dimensional track finding, the so called matching, has been incorporated. This routine assigns a value to each combination of four two-dimensional tracks. This scalar criterion indicates the quality of a matching.

For this valuation we took advantage of the special setup of PS185. The experiment consists of two spatially separated multiwirechambers, a driftchamber (DC) and a proportional chamber (PC) with two orientations each. Therefore it is possible to define a set of three-dimensional tracks for each system independently.

These skew lines are compared pairwise and opening angle, distance of DC track to PC trackcenter and vice versa are calculated. The distribution of the quality criterion yield as a weighted sum over the squares of these three quantities is shown in figure 1. Every two-dimensional track may vote for one of all possible combinations according to the calculated criterion. If all four concerned two dimensional tracks agree in the combination it is declared as a matching.

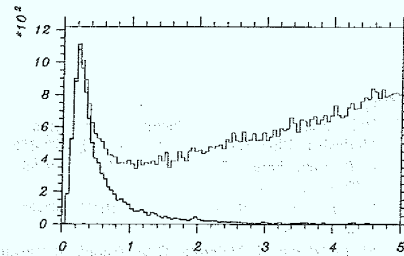


Figure 1: quality criterion
Dots: all combinations
Line: matchings only

This and all other distributions have small but long tails for matched combinations. A numerical cut in one or more of those quantities can not lead to such a result. With this routine the complete set of combinations is checked and the consistence criterion reduces this multitude of track candidates to a small class of matchings almost without numerical restrictions. For each single two dimensional track the best possibility is found and none is used twice.

This new algorithm increases the number of tracks used for a full (four-track) matching by a factor of 1.6. Whilst the number three-track matchings decreases by a factor of 3. This avoids the former buffer overflow and is an indication for the achieved full reconstruction of three dimensional tracks along with a reduction of combinatorial background.

This clear track pattern leads to the aspired reduction of geometrical solutions per event, which is reduced to 44% in average. It is a large advance towards a plain event reconstruction of maximal one physical significant vertex combination per trigger. This is already achieved for 80 % of all events with at least one vertex combination. Figure 2 shows a Dalitz plot for events with exactly one solution. A cut in $M^2(> -0.03 \text{ MeV}^2)$ has been applied. There is a clear signal of $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ and $\bar{p}p \rightarrow \bar{\Lambda}\Sigma$ events.

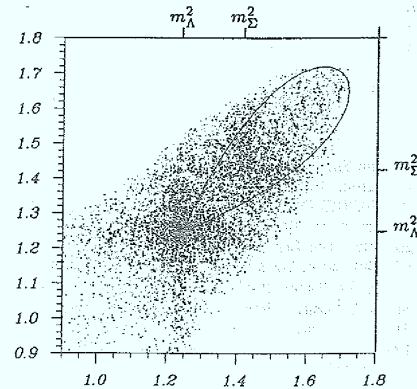


Figure 2: Dalitz plot for $\bar{p}p \rightarrow \bar{\Lambda}\Lambda X$
indicated is the contour of $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\gamma$

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Measurement of the reaction $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ near threshold

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The PS185/2 collaboration studies the $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ reaction to resolve a possible structure near the reaction threshold[1-3]. For the data taking in 1994 the complete PS185 experiment had to be reinstalled. The wire chambers were partially rewired and cleaned. The target was a newly build sandwich target in the design of former experiments with four CH₂ cells and one carbon cell. Two sets of μ -strip planes were put into the beam right upstream the target to measure the incoming antiproton's track.

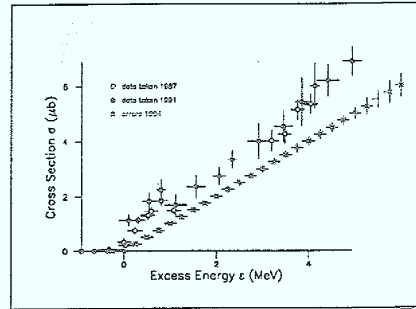


Figure 1 The excitation function of the reaction $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ from its threshold up to 5 MeV. The open and full circles are data taken by the collaboration [1,2]) in the past. The stars assume a linear cross section $\sigma = \epsilon \cdot \frac{\mu b}{MeV}$, while their error bars are calculated from the statistics obtained at each energy.

During three weeks in June and July 1994 data were taken at seven different momenta near the reaction threshold. Intermediate momenta were obtained using a degrader in the beam. The momenta were chosen to scan the irregularity observed in the previous data[1,2] with good statistics. Those data indicate a width of the structure of the order of 1 MeV. At the position of the irregularity the cross section drops from 1.84 ± 0.17 (2.24 ± 0.38) μb at 0.8 MeV to 1.48 ± 0.15 (1.68 ± 0.39) μb at 1.08 (1.12) MeV excess energy in reference [1] ([2]). To achieve statistical significance as well as the resolution required to resolve this structure the momentum was varied corresponding to steps of 0.5 MeV in the excess energy. The average antiproton beam intensity was around 700 KHz, giving a luminosity around $5 \cdot 10^{28} sec^{-1} cm^{-2}$. Data were accumulated until the statistical error on the cross sections could be expected to become 5% or less. Due to the segmentation of the sandwich target four energy bins are measured at each momentum setting. Having the vertex constrained by the measurement of the beam track with the silicon μ -strip detectors upstream the target the energy resolution is expected to improve by more than a factor three in the kinematical fitting procedure to be less a 100 KeV.

P (MeV/c)	degrader	ϵ (MeV)	$\bar{p}$ $\times 10^{10}$	$\bar{\Lambda}\Lambda$ (expected)
1439.7	in	0.5	1.735	31
1439.7	out	1	9.23	823
1442.6	in	1.5	6.06	1018
1442.6	out	2	5.58	937
1445.6	in	2.5	1.61	829
1445.6	out	3	2.495	1113
1448.5	in	3.5	2.40	1283
1448.5	out	4	2.16	1351
1451.6	in	4.5	1.16	828
1451.6	out	5	2.32	1853
1454.5	in	5.5	0.43	429
1454.5	out	6	0.75	882
1477			0.9	2000

Table 1 The momentum settings with and without degrader and the excess energy of the most upstream cell are listed in columns one to three. Columns four and five contain the number of antiprotons used at this setting and, using a linear dependence of the cross section on the excess energy, the expected number of $\bar{\Lambda}\Lambda$ pairs.

In total $36.8 \times 10^{10} \bar{p}$ were used for data taking. The readout rate was $0.87 \times 10^{-4} / \bar{p}$, resulting in about 33 million events on tape. Assuming a linear cross section $\sigma = \epsilon \cdot \frac{\mu b}{MeV}$ and using acceptance and efficiencies of [2] the number of expected $\bar{\Lambda}\Lambda$ events can be calculated. A total of around 12000 $\bar{\Lambda}\Lambda$ can be expected after analysis. The estimates for the individual momentum settings are listed in the table.

- 1) University of New Mexico, Albuquerque
- 2) Los Alamos National Lab.
- 3) CERN, Geneva
- 4) Carnegie Mellon University, Pittsburgh
- 5) Universität Erlangen
- 6) Albert Ludwigs Universität Freiburg
- 7) Rutherford Appleton Lab
- 8) University of Illinois at Urbana Champaign
- 9) The Svedberg Laboratory, Uppsala

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Estimate of the Quality of a Measurement of Spintransfer in $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ with PS185 K. Rührich for the PS185 collaboration

The error on the depolarization D_{nn} has been estimated for a measurement of the reaction $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ with the PS185 experiment[1] and a frozen spin target[2].

A Monte Carlo program was used, which created events in pure phase space as well as measured distributions. It did not include effects from interactions of the reaction products with detector material (multiple scattering, δ -ray production or hadronic interactions). The target was assumed to be a cylinder of 1 cm diameter and 1 cm length oriented along the beam direction. The event trigger is based on the charged-neutral-charged signature of the events. It therefore implies the installation of veto counters, which cannot be closer to the target than what its vacuum chamber allows. Those veto counters were simulated as being a rectangular beam veto counter of 1 cm² and a larger 12x12 cm² scintillator with a beam hole of the size of the beam veto counter. For tracking and baryon identification the standard PS185 detector was simulated. The distance of the veto counters to the first wire chamber plane was 1.6 cm and the distance between the last wire chamber plane and the hodoscope was 11 cm. To approximately estimate material effects, events having one or more of the charged particles with β less than 0.1 were discarded. An event was accepted if it had no veto counter hit and at least one hit in the hodoscope. Furthermore it was required that the tracks of all four charged final state products (only charged Λ decays were considered) are seen at least in both projections in one of the chamber stacks. A dead zone for tracks with laboratory angles of $90^\circ \pm 30^\circ$ was introduced. For a determination of D_{nn} it is essential, that the hyperon or antihyperon is identified. At least one track must have had hits in the magnet chambers. Individual counter efficiencies were assumed to be 100%.

The events were generated with a momentum of 1642 MeV/c. The beam profile was assumed to be of gaussian shape with a width in x and y (eg. perpendicular to the beam) of 1 mm. The $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ angular distributions of [3] for cross sections and polarizations were taken into account.

The results of this simulation can be summarized as follows: Many events are lost due to the cut on the decay point of the hyperons/antihyperons implied by the veto counters. In case of a pointlike target about 50% of the events are lost, if the veto counters are only 1 cm away. The loss increases to 27% if the distance is 3 cm. The loss due to backward going pions is typically between 2% and 3%. A second loss occurs when tracking is required. The finite size of the detectors as well as the inefficiency for tracks at 90° kill another third of the events. The overall efficiency is as small as 33% in the case where the veto is 1 cm away and becomes less than 20% for a minimum decay length of 3 cm. The dependence of the acceptance on the production angle as well as on the decay baryon's angle is smooth and varies only little. No regions with zero acceptance have been found.

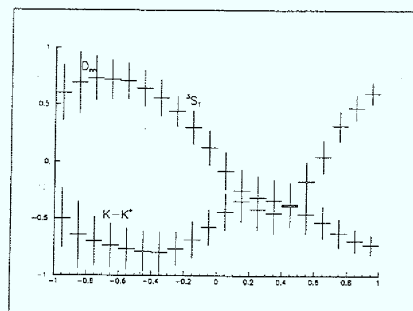


Figure 1 The depolarisations D_{nn} are taken from [4], one using a 3S_1 quark-gluon model, the other one being based on a $K-K^*$ exchange. The errors on D_{nn} are based on the data of [3] and were evaluated using the calculated acceptance and assuming a one week running at LEAR.

A measurement of D_{nn} must be decisive on different models. The statistical error for each angular bin was estimated using an simulation of the experiment with event generation based on measured angular distributions[3] and a proper detector description. About 20000 $\bar{\Lambda}\Lambda$ events after decay branching ratio, trigger and reconstruction conditions were accepted and used for the error calculation. The error is in the here considered cases only weakly dependent on the actual strength of the polarization. It is determined by the event statistic as well as by the quality of the measurement of $\bar{\Lambda}\Lambda$ polarization with an unpolarized target, which is a necessary input to determine D_{nn} . Assuming a realistic beam rate of 500 KHz an eight days run could result in the statistical errors shown in the figure.

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During the year 1994 data taking was completed and data analysis was continued on the $\bar{p}p$ annihilation into the two meson final state $\phi\phi$ observed via the decay into the four prong signals of $K^+K^-K^+K^-$.

The primary goal of the JETSET (PS202) experiment is to explore the meson mass spectrum above 2 GeV/c². The experiment searches for s-channel resonances which decay into four kaons. In formation mode, the experiment is sensitive from near threshold to the highest mass accessible with the LEAR machine, around 2.4 GeV/c². In the four-kaon final state, we have an enhanced sensitivity to resonances with high strangeness-content. There are many conventional mesons with a large strangeness admixture, however, these are generally expected to be suppressed in $\bar{p}p$ annihilations. The reason for this is that the nucleon (antinucleon) is known to be composed primarily of $u(\bar{u})$ and $d(\bar{d})$ quarks, and to couple most readily to final states composed of these quarks. This dynamical selection rule is related to the more general principle - OZI rule -, that the dominant hadronic processes involve the minimum of quark pair creation/annihilation. The cross section for $\bar{p}p \rightarrow 4K^+$ should be suppressed relative to the dominant multi-pion reactions. However, this suppression might be overridden by the formation of a flavour-neutral intermediate state, such as a gg exotic. It is a general feature expected for exotics that they should be less flavour-specific in their couplings than conventional $q\bar{q}$ states. Hence the OZI rule cannot apply to exotic-formation reactions. For this reason the reaction $\bar{p}p \rightarrow 4K^+$ has special interest in the search for exotic mesons above 2 GeV/c².

Another dynamical selection rule observed in $\bar{p}p$ annihilations is that the reaction proceeds primarily through two-meson channels. In the four-kaon final state near threshold this would correspond chiefly to $2\phi(1020)$ production, with probably some contribution from the tail of the $f_0(980)$, and from the $f_2(1270)$, $f_0(1300)$, and $\alpha_2(1320)$ at the highest values of K^+K^- invariant mass. With a deviation from ideal mixing $\theta-\theta_{ideal}$ of only 1° to 4° , the ϕ is nearly 100 % $s\bar{s}$. If the OZI rule is interpreted to forbid strangeness production in $\bar{p}p$ annihilations then the process $\bar{p}p \rightarrow \phi\phi$ can proceed only via the small $u\bar{u}, d\bar{d}$ component present in the ϕ wave function. From this model one can derive an upper limit $\tan^2(\theta-\theta_{ideal}) \approx 2.5 \cdot 10^{-5}$ for the ratio of cross sections $\sigma_{\phi\phi}/\sigma_{\omega\omega}$ for production from $\bar{p}p$ annihilation. Although the cross section of $\bar{p}p \rightarrow \omega\omega$ has not been measured directly, an estimate can be obtained from the total $2\pi^+2\pi^-2\pi^0$ cross section [1], which was measured to be about 5 nb in the energy range of our experiment. There are many reaction channels that contribute to this final state. If we estimate that it is 10 % $\omega\omega$ then the expected $\phi\phi$ cross section is 10 nb.

The observed cross section for the $\bar{p}p \rightarrow \phi\phi$ reaction is much larger. The reason for a large violation may not lie in the existence of strong resonances in the formation channel. The origin of OZI violation in $\bar{p}p$ annihilation reactions is one of the important topics and it is probable

that the physics behind the more general phenomenon is playing a role in the $\phi\phi$ channel. The experimental question remains, however, whether resonance formation may also be contributing to the large cross section for this reaction, and if so, what the resonance parameters might be.

The PS202 experiment utilizes a molecular hydrogen-cluster gas jet target located in a straight section of the LEAR antiproton storage ring at CERN. A cut view of the detector is shown in Fig. 1. The antiprotons circulating in the machine pass from left to right in the figure. The jet is normal to the beam and lies in the horizontal plane. Apart from the constraints imposed by the jet fixtures, the detector is azimuthally symmetric. It is divided at polar angle $\theta_{lab} = 45^\circ$ into barrel and forward regions. Beginning at the interaction vertex, the particles entering the acceptance of the detector traverse the elements in the following order:

1. thin-walled beam pipe, a corrugated section of 300 μ m stainless steel extending 30 cm downstream of the target, of elliptical cross section with horizontal and vertical half-axes of 7.8 and 3.8 cm, respectively;
2. inner trigger scintillators, finely segmented in azimuthal angle ϕ and divided between barrel and forward;
3. drift chamber, composed of packed layers of cylindrical straws operating in proportional mode;
4. silicon pads of thickness 300 μ m for measurement of dE/dx ;
5. threshold Cerenkov counters filled with either water or liquid freon, for fast pion rejection;
6. outer trigger scintillators, three-layer hodoscope forming pixels for fast track finding and luminosity measurement;
7. electromagnetic calorimeter composed of a lead-scintillating-fiber sandwich.

Several small dedicated silicon detectors for luminosity monitoring are also present. The RICH in the forward region and the MWPC (z-chamber) in the barrel were present only for the later running periods. The experiment has no magnetic field.

The online 4K trigger was based upon the charged-particle multiplicity from the trigger scintillators combined with low pion multiplicity from the Cerenkov counters, and low gamma multiplicity from the calorimeter. The azimuthal hit pattern in the inner trigger scintillators was required to be consistent with transverse momentum balance.

The size of the jet target was about 8 mm f.w.h.m., compared with the typical $\bar{p}$ beam dimensions of about 3 mm r.m.s.. The density of the jet target was typically $4 \cdot 10^{12}$ atoms/cm², averaged over the profile of the beam. During later runs this was raised to $6.5 \cdot 10^{12}$ atoms/cm². A typical stored beam of $4 \cdot 10^{10}$ antiprotons circulating at a frequency of 3.2 Mhz (1.4 GeV/c) gave a luminosity of $5 \cdot 10^{29}$ cm⁻²s⁻¹. Assuming full acceptance and modest experiment dead-time, a $\phi\phi$ cross section of 4 μ b at this luminosity would yield a $4K^+$ trigger rate of about 0.5 Hz. The actual trigger rate at this luminosity varied from 100 to 500 Hz,

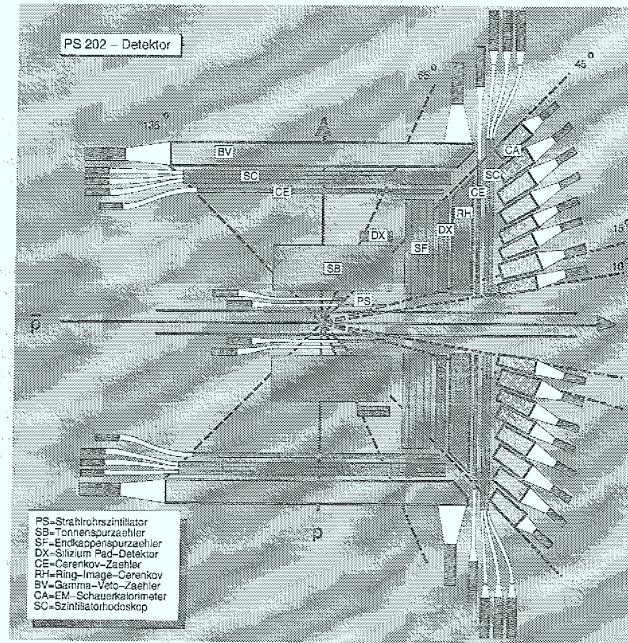


Figure 1: Layout of the JETSET (PS202) experiment, showing the major detector components.

depending on the beam energy and the Cerenkov radiator in use.

The acceptance is much less than unity, owing to the large loss of particles in the low polar angle region, where the tracking is limited by the LEAR beam pipe. The priority of keeping the tracking close to the interaction region, in order to avoid losses from kaon decays, conflicts with the need to extend the coverage as low as possible in polar angle. The compromise implicit in the final detector design gives an effective cutoff near $\theta_{lab} = 15^\circ$, averaged over ϕ . The resulting acceptance for $\phi\phi$ varies strongly with the production angle between 50 % and zero. In spite of its acceptance deficiency, this experiment provides the first detailed measurement of these $\bar{p}p \rightarrow 4K^\pm$ cross section, and in particular the excitation function of the $\phi\phi$ channel, with good statistical precision and excellent energy resolution. The analysis leading to a sample of four-kaon events proceeds via four steps; four-prong reconstruction, 4K candidate selection, kinematical reconstruction, and confirmation using particle identification (PID).

With $\Delta E = [E(p4) - E_o]_{min}$ we define ΔE as the minimum of the difference between the total energy in the laboratory system of the final four K state and the total energy E_o of the initial $\bar{p}p$ state.

An example of the ΔE distribution for all four-prong events is shown in Fig. 2a. The same distribution following 4K candidate selection is shown in Fig. 2b. Here one begins to see evidence of an edge at $\Delta E = 0$ which is the expected signature of the four-kaon reaction. However, there is still a large background present in the sample, as evidenced by

the preponderance of events with $\Delta E \gg 0$. To eliminate this background, the reconstructed four kaon momenta are required to be confirmed by the PID measurements: dE/dx in the silicon pads and number of photoelectrons in the Cerenkov counters. In addition to confirming good 4K events, the PID tests are also used to distinguish the correct kinematical solution from the incorrect one, in the case that there were two. Following PID confirmation, the ΔE distribution is as shown in Fig. 2c. The remaining non-4K background can be largely eliminated by a cut at $\Delta E < 0.02$ GeV. A crude estimate for the background that remains under the peak in the $\Delta E < 0$ region for the final sample, shown by the dashed line in the figure, is about 15 %.

A plot of the invariant masses of two-kaon pairs m_{12} vs m_{34} is shown in Fig. 2d. The most prominent feature of the plot is the large accumulation around the mass of the $\phi(1020)$. The peak is centered at 1019.4 ± 0.4 MeV and has a f.w.h.m. of 14.6 ± 1.1 MeV, which reflects the finite precision in the kinematical reconstruction and not the intrinsic width of the ϕ . Because the charge of a track is not measured by the detector, each event contributes three entries in the mass plot. The band along the upper diagonal in the plot is the combinatoric reflection of the $\phi\phi$ peak. About 70 % of the 4K events in this sample taken at a $\bar{p}$ momentum of 1.4 GeV/c are associated with the $\phi\phi$ peak, with the remaining 30 % spread out in a broad continuum over the surface of the mass plot. To obtain a quantitative measure of the $\phi\phi$ fraction in the sample, we have applied a maximum likelihood fit to the data, under the hypothesis that each kaon either comes from ϕ decay or

Efficiencies for the Forward Hodoscope in JETSET

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The circular three layer hodoscope[1] has been employed in various triggers of the experiment PS202 at LEAR at CERN (JETSET). The scintillators are expected to have an efficiency for charged particles very close to 100%. Nevertheless, due to fluctuations of the pulse heights as well as geometrical inefficiencies due to little gaps between the counters, their absolute efficiency is somewhat smaller.

During data taking the JETSET experiment has recorded events with multiplicity 2 in coplanar configuration. These events come mainly from elastic proton-antiproton scattering. They can be used to determine the efficiency of individual counters in the hodoscope. The accepted elastic events cover the full azimuthal angular range and the polar angular region from 15° up to about 40°.

The events were selected after criteria for elastic scattering. They have one track in the forward tracking detector and a second track beyond 45° due to the trigger condition. The forward track's impact point on the hodoscope was calculated and the expected hit scintillators determined. To make sure that the particle would traverse the scintillators, one cluster of hits was required in the calorimeter further downstream, with its center not further away from the track impact than half the cluster radius. Due to multiple scattering, especially for 1992 and later when the implementation of the RICH detector brought more material into the particle's track, the particle might not just hit the scintillator in the straight extrapolation but the left or right neighbour. As a prediction therefore the scintillator on the track and its both neighbours in each layer were taken. The efficiency was calculated as the ratio of hits found over hit predictions for that scintillator. In case of two or three scintillators hit due to secondary effects like backscat or 6-rays only the central one was taken into account.

The efficiencies have been determined for the data taking periods in 1991 and 1992. Integrated over longer run periods the statistical errors, being as large as 30% for individual runs, can be neglected. The resulting efficiencies[2] are always larger than 95%, but typically they vary between 97.5 and 99%. An inefficiency of about 1.1% is introduced due to the 400 µm gaps between the counters. Thus efficiencies larger than 99% cannot be expected. It was found that some counters systematically have lower efficiencies in the order of 95% only. This must be attributed to a slightly asymmetric mounting as well as gravitational pull, which makes the scintillators in the lower hemisphere bend downwards a bit and thus increase the gaps between them in certain areas.

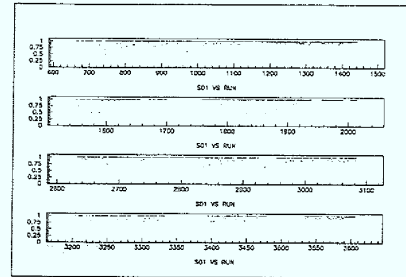


Figure 1 Run by run efficiencies for wedge scintillator FSS01 for 1991 and 1992. The average efficiency is around 99%. The fluctuations to low efficiencies mainly arise from runs with very few events only.

The investigation also showed that with the chosen voltages and discriminator threshold settings there is hardly any difference when using the TDC hit or the recorded ADC with an offline cut. The difference between ADC and TDC efficiency is less than 1%.

To illuminate the stability with time the figure shows the run by run efficiency of one of the straight counters of the forward hodoscope. The runs with lower efficiency have in most cases very little statistics. This looks more or less the same for all other counters. Variations of the high voltage, as they were adjusted in the beginning of the data taking hardly influence the efficiencies, but rather improve the ADC measurement.

To summarize the efficiency of the scintillation counters in the forward hodoscope of the JETSET experiment are close to 100%. The observed inefficiencies can to their larger part be attributed to the gaps between the counters.

- 1) University of Bari and INFN, Bari
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- 3) Universität Erlangen-Nürnberg
- 4) Albert Ludwigs Universität Freiburg
- 5) University of Genova and INFN, Genova
- 6) University of Illinois at Urbana Champaign
- 7) University of Oslo
- 8) The Svedberg Laboratory, Uppsala

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The aim of the JETSET[1] experiment at LEAR, CERN is to study the $\phi\phi$ final state in $\bar{p}p$ annihilations in flight. The ϕ -mesons decay to about 49% into two charged K-mesons. Based on the signature of four charged final state kaons an apparatus with good tracking and particle identification in the kinematically allowed angular region was built. A large fraction of possible background is avoided by selecting events with tracks in a limited angular range (15° to 65°). The lower limit is imposed mainly by LEAR's vacuum beam pipe. This causes a hole in the experiment for events with the ϕ mesons close to the beam axis ($-0.8 \leq \cos\Theta^{CMS} \leq +0.8$). Those events have a typical signature of one or two fast forward going kaons and two slower kaons at larger angles. Making use of the circular forward hodoscope [2] a trigger [3] was formed to catch some of those events.

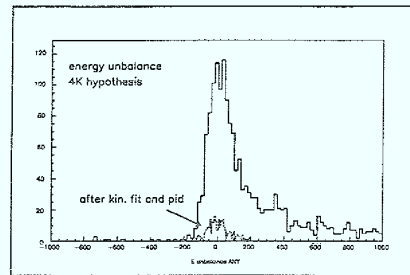


Figure1 Energy unbalance after kinematic calculation. The solid line is the distribution before particle identification. The dashed distribution remains after pid.

About 350K events were recorded at two beam momenta 1465 and 1505 MeV/c. The accumulated events correspond to a luminosity of $40 \mu b^{-1}$.

The analysis was based upon three major characteristics. There must be two clear tracks in the forward tracker. This is necessary to define the production vertex, which together with the small angle pixel in the forward hodoscope defines the third track. A slow, large angle kaon will only give a signal in the beam pipe scintillator, where it is slowed down or even stopped. It eventually decays and in 63.5% the decay product is a muon accompanied by a neutrino which is not detected. The muon track can be reconstructed. Its impact point on the surface of the beam pipe was assumed to be the decay point of the kaon. This is a crude approximation, but Monte Carlo studies have revealed a geometrical acceptance of about 6% and a reconstruction efficiency of still 5%. The kinematic was calculated for the reconstructed events for all possible mass hypotheses. In figure 1 the deviation from energy conservation (energy unbalance) is displayed as calculated from the track directions before and after particle identification for the 4K mass hypothesis. Background, which appears mainly at large energy unbalance, is al-

most completely eliminated. The distribution of events in these spectra are well reproduced by the Monte Carlo simulations.

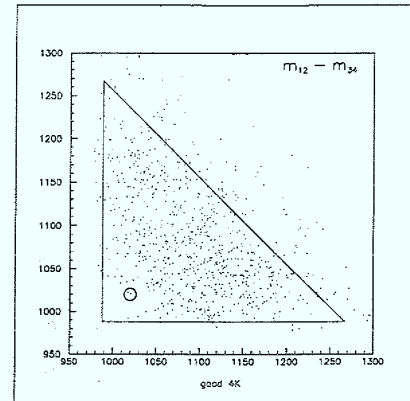


Figure2 The Goldhaber plot for the 4K hypothesis. The triangle indicates the kinematically allowed region. The circle is at the position of the $\phi\phi$ -peak and has a radius of $7 \text{ MeV}/c^2$.

The second figure shows the distribution of invariant masses for pairs of tracks assuming all being kaons. In the region where $\phi\phi$ events should appear, indicated by the circle at masses of $1020 \text{ MeV}/c^2$, only few events are left. This is in agreement with an estimated cross section in the order of two μb . The major reason for the small number of final events was the trigger prescaler (2^6) which was necessary to keep the trigger rate below 10% of the total rate.

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G.Baur, R.Bröders, M.Chanel[†], M.Dahmen, D.Grzonka, K.Kilian, C.Lacroix[†], P.Lefèvre[†], D.Möhl[†],
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Like in 1993 [1] a measurement parasitically to the PS202-experiment at LEAR/CERN was done in order to observe anti-hydrogen ($\bar{H}^0$) production. The underlying process is the creation of an electron-positron pair (e^+e^-) during the exchange of two photons emitted by the beam anti-proton ($\bar{p}$) and a target proton (p) [2] [3].

$$\bar{p}p \rightarrow \bar{p}pe^+e^- \rightarrow (\bar{p}e^+)pe^-$$

The e^+ – if it has the appropriate velocity – can be captured by the $\bar{p}$ thus forming an $\bar{H}^0$ -atom, which leaves the accelerator through a hole in the next bending magnet. Behind this dipole – still in vacuum – the $\bar{H}^0$ ionizes in an 30μ Al foil. The e^+ is stopped in a subsequent silicon-counter ($700\mu m$, $1200mm^2$) leaving all of its kinetic energy while the $\bar{p}$ due to its high energy crosses this and all following detectors, namely a second silicon-counter ($500\mu m$, $1200mm^2$) and the scintillation counters (BC-404, 5 mm, $100cm^2$) for time-of-flight measurement, inducing signals proportional to the individual energy losses. The e^+ annihilates in the silicon and the two 512 keV photons (γ) are detected in a surrounding six-fold segmented NaI-detector.

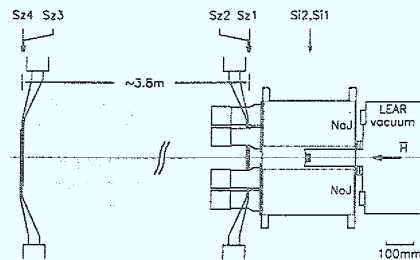


Figure 1: Experimental setup

The data analysis is based on following measured values: scintillator TDCs for time-of-flight and so velocity information from the p/ $\bar{p}$, scintillator ADCs for velocity check via $(dE/dx)_p$, silicon ADC's for $(dE/dx)_p + T_e$ and $(dE/dx)_p$ resp. to ensure the existence of pe^- or $\bar{p}e^+$, and NaI ADCs for detection of the two e^+ -annihilation γ 's which then indicate at least the existence of the $\bar{p}e^+$ -pair.

Data were taken during three weeks at different beam momenta: in steps of 20 from 1180 to 1280 MeV/c, in steps of 30 from 1315 to 1375 MeV/c, and at 1500, 1550 and 1800 MeV/c. The choice of these momenta was governed by the physics program of the PS202-experiment.

The time-of-flight spectrum in figure 2 consists of mainly two components: a peak at high velocity which is due to fast flying particles from $\bar{n}$ induced reactions producing e.g. multiple pions, and a peak representing particles with velocities slightly slower than the beam itself. These particles are supposed to be p or $\bar{p}$ mainly due to charge

exchange reactions where a $\bar{n}$ scatters at a n and a $\bar{p}p$ -pair is produced. The other source of this type of events is the $\bar{H}^0$ -production.

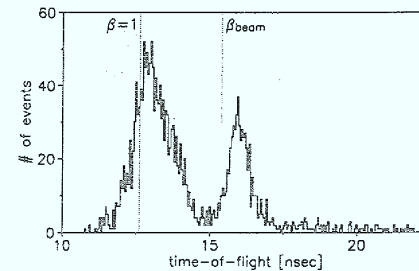


Figure 2: Time-of-flight spectrum at 1345 MeV/c

With the time-of-flight trigger 904,520 events were acquired. Applying several appropriate cuts with four standard deviations around the expected observables the following candidates remained: four events with a clean e^+ -annihilation but no acceptable signals in the silicon counters, five events with a clean e^+ -annihilation but time-of-flight and scintillator ADC values indicating multi pion events, and at least three events showing time-of-flight values in the expected range, consistent dE/dx measurements, silicon signals with the correct signal height for the $\bar{p}e^+$ -pair from $\bar{H}^0$, but no NaI-signature in opposite crystals which would indicate a e^+ -annihilation doubtlessly. All six crystals would have to be summed up to fulfil the required 2×512 keV signature. On the other hand, calibration measurements with a ^{22}Na -source have shown that the total efficiency, i.e. geometrical plus detection of the two γ 's, is around 35 %.

In 1995 an improved setup will be installed which allows for particle identification using a magnetic separator system. Heavier nuclei will be used as a target to improve statistics since $\bar{H}^0$ -production cross section is proportional to Z_{target}^2 .

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Study of the Roper Resonance Excitation in Light Nuclei

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Experiments on the excitation of the $P_{11}(1440 \text{ MeV})$ resonance in α -scattering at $E_\alpha = 4.2 \text{ GeV}$ have been performed at Saturne for the nucleon [1] and for nuclear targets [2] (^2H and ^{12}C). This type of excitation corresponds to a compression of the nucleon and is directly related to the radial degree of freedom of the system. It is of particular interest to study nuclear medium effects on this hadronic degree of freedom.

In addition to the medium effects discussed in ref. [3] a swelling of the nucleon in the nuclear medium has been discussed in the literature [4]. This will effect the excitation energy of the compression mode which can be estimated using energy weighted sum rules [5]

$$E_\alpha^2 = \frac{3\hbar^2}{m_N} K_N \frac{1}{\langle r_N^2 \rangle} \quad (1)$$

The nucleon compressibility K_N is also related to the radius. Using $K_N \sim 1/\rho$, where ρ represents the density $\rho = 4\pi/3 \langle r^2 \rangle^{3/2}$, this gives $K_N \approx 9/4\pi \langle r^2 \rangle^{3/2}$. From this we obtain

$$E_\alpha^2 \approx \frac{27\hbar^2}{4\pi m_N} \frac{1}{\langle r_N^2 \rangle^{5/2}} \quad (2)$$

From this expression we see that already a small change of the nucleon radius would result in a relatively large change of the excitation energy of the monopole excitation. Estimates of the swelling of the nucleon based on K -scattering total cross section measurements [6] yield $r \sim 1.1 r_0$. Such a change of the nucleon radius in nuclear systems would give rise to a lowering of the monopole excitation by about 80 MeV. Certainly, in our experiment we cannot expect to observe such large effects due to the fact that α -scattering is only sensitive to the surface of the nucleus which has a reduced density.

Experimental spectra for ^2H and ^{12}C are given in fig. 1. Similar to the scattering from the nucleon two components are observed in the spectra, target excitation and coherent π -production from projectile excitation. Calculations of the latter process in quasifree kinematics are given by the solid lines which give a good description of the spectral form above the π -threshold. In this calculations the Fermi-motion is taken into account using the Hulthen wave function [7] for ^{12}C .

The difference between the experimental data and the solid lines in fig. 1 shows the effect of target excitation (see fig. 2). In this figure a similar spectrum for the case of the nucleon [1] is included (upper part). For this, the solid line presents the spectral shape of the P_{11} -resonance using the π -N P_{11} phase-shifts [8]. The fact, that a reasonable description is obtained supports the interpretation of excitation of the Roper resonance. For the nuclear systems this excitation is also observed. The excitation energy is not much changed, whereas the width is quite a bit in-

creased. Calculations of the spectral shape of the Roper resonance in quasifree kinematics are given by the solid lines, the Fermi-motion of the nucleon is taken into consideration in the same way as for projectile excitation.

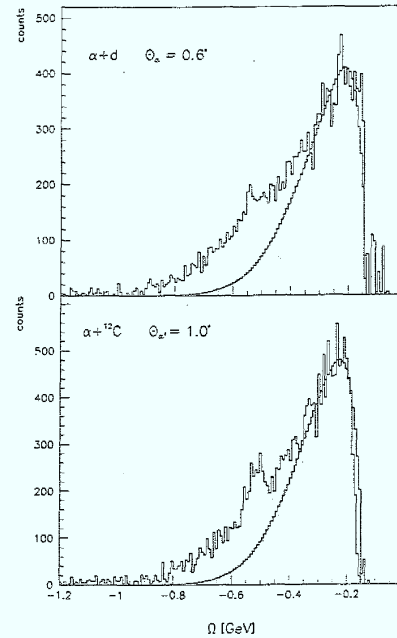


Figure 1: Missing energy spectra (Ω) for inelastic α -deuteron and α - ^{12}C scattering. The solid lines indicate the contributions from projectile excitation.

This gives a reasonable description of the spectral shapes with an increase in width consistent with the data. For the deuteron this indicates that the present experiment is in good agreement with the previous data [1]. For ^{12}C the spectra are also reasonably described, a shift of the resonance is not observed within the experimental uncertainties. This may exclude a swelling of the nucleon in the scalar field of α -scattering, at least for the surface nucleons in ^{12}C .

In the same experiment the elastic α -scattering from ^2H and ^{12}C is measured. The data are already published [9] together with an optical model analysis.

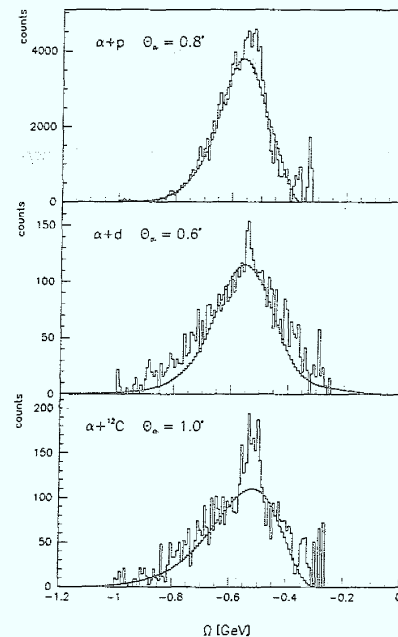


Figure 2: Difference spectra between the experimental Omega-spectra in fig. 1 and the contributions from projectile excitation (solid lines in fig. 1). In the upper part a spectrum for the nucleon is included which is well described using the π -N P_{11} phase shifts. The solid lines for the deuteron and ^{12}C are obtained in the same way taking into account the Fermi motion in these systems.

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Meson Production Experiment PROMICE at CELSIUS

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for the PROMICE collaboration

The research program of the PROMICE (PROduction Of Mesons In CELsius) experiment at The Svedberg Laboratory uses the CELSIUS storage ring in Uppsala (Sweden) with the aim to study fundamental reaction mechanisms for the π^0 and η meson production and to investigate the role of different resonances for the reaction cross section, especially close to the production threshold. In 1994 first results from proton proton collisions and the $pd \rightarrow {}^3\text{He}\eta$ reaction at 1250 MeV and 1276 MeV were obtained [1].

The experimental setup has been improved last year. In the forward direction a window-detector for fast particle separation in the trigger decision was added and two planes of a straw-chamber detection system were installed additionally to the unchanged thin plastic-scintillator hodoscope (Quirl) and the segmented calorimeter. The internal cluster jet target and the well working central detector for the measurement of neutral and charged tracks were unchanged. Behind the calorimeter is a new vetohodoscope.

In February '94 a 1 GeV proton beam was successfully used for testing both the setup changes and different triggers. Data for the $pd \rightarrow {}^3\text{He}\eta$ reaction were taken subsequently, with ~ 6000 $pd \rightarrow {}^3\text{He}\eta$ events expected.

In July '94 was a run for Bremsstrahlung measurement. In October/November '94 a second data taking for the $pd \rightarrow {}^3\text{He}\eta$ reaction took place at four energies close to threshold and at the magic momentum for the η -meson production (930 MeV, 965 MeV, 1037 MeV, 1100 MeV proton beam). From a very first check and the online control program the number of measured reactions can be estimated to 30,000 events per energy setting.

Very preliminary results for the improved particle separation are shown in figure 1, where the energyloss in the last plane of the Jülich-hodoscope is given versus the energy deposit in the ${}^3\text{He}$ stopping range hodoscope plane.

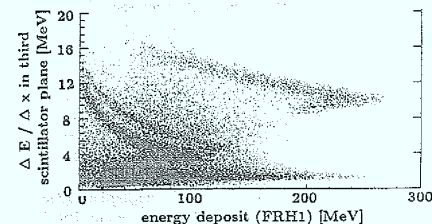


Figure 1: Energy loss in third Jülich hodoscope plane (dE) versus energy deposit in the first range hodoscope plane (E) from proton-deuteron collisions at $T_p = 1037$ MeV.

The possible Missing Mass resolution is two to four times better with the new tracker for the ${}^3\text{He}\eta$ production, as shown in figure 2.

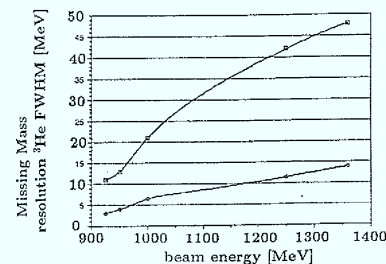


Figure 2: Missing Mass resolution for $pd \rightarrow {}^3\text{He}\eta$ versus beam energy. Upper curve old setup, lower curve including the new installed straw-chamber system.

The η is identified by the missing mass of the ${}^3\text{He}$ recoil nucleus. In addition, a full event reconstruction is possible by measuring the two η decay photons and the ${}^3\text{He}$ recoil. This nearly background free event sample has certainly a much lower counting rate. Figure 3 gives a preliminary missing mass distribution for full reconstructed events from the November run. It contains about 10 % of the data.

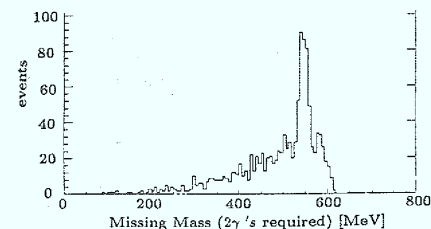


Figure 3: Missing mass distribution for $pd \rightarrow {}^3\text{He}\eta$ event at $T_p = 1037$ MeV.

In the final analysis angular distributions and total cross sections will be evaluated to get more information about the remarkable strong η final state interaction.

Reference:

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One of the challenges in medium energy physics is to find unambiguously objects which are not simple states of a quark-antiquark pairs or three quarks like mesons or baryons. Non-standard quark states are conceivable and have been predicted. If found, they will provide strong constraints on the QCD confinement potential and the extrapolation of the asymptotically free quark-gluon models into the non-perturbative region. One class of non-standard objects are six-quark particles or dibaryons. Of special interest are the strangeness containing dibaryons. Historically called "H-dibaryons" several experiments[1] have been performed to search for those particles. Typical upper limits for the production of H-dibaryons are in the order of a few hundred nanobarn. Here I discuss the possibility of an experiment to detect the double strange charged dibaryon H^{++} in the reaction chain

$$pp \rightarrow H^{++} \rightarrow \Sigma^+ \Sigma^+$$

The threshold momentum is 1.352 GeV/c corresponding to a mass of 2.3788 GeV/c².

Detecting the reaction $pp \rightarrow \Sigma^+ \Sigma^+$ would be the first measurement of a $\Delta S = 2$ transition outside the neutral meson systems. This reaction is principally not forbidden by any quantum number selection rules. But because it involves a transition between different isospin/strangeness states (I=1 and S=0 to I=2 and S=-2) it can only proceed via the electroweak interaction. In the standard model (SM) neutral current interactions do not change flavor. Therefore the exchange of at least two charged W-bosons is needed and thus the process is suppressed by at least m_W^2 . Due to the exchange of the heavy boson masses the interaction is extremely short ranged. In the meson exchange model (MEM) a K^0 from the one baryon would have to convert into a $\bar{K}^0$ to get absorbed by the other baryon. The transition $K^0 \leftrightarrow \bar{K}^0$ is well known and within the MEM the reaction $pp \rightarrow \Sigma^+ \Sigma^+$ measures the short range $\Delta S = 2$ in the neutral K system. As discussed by [2] a determination of the short range strength of the $\Delta S = 2$ transition could constrain the masses of higgs particles or the heavy partners of the gauge bosons in left-right symmetric models.

The production of $\Sigma^+ \Sigma^+$ in proton-proton reactions is a slow process and the reaction time is short. The interaction time and therefore the chance for two d-quarks to turn into s-quarks is enhanced, if intermediate resonances would exist [3]. At the position of the resonance the $\Sigma^+ \Sigma^+$ production should be enhanced. Because the nonresonant background in the reaction $pp \rightarrow \Sigma^+ \Sigma^+$ is expected to be extremely small, this process is perfectly suited to study dibaryonic resonances which might decay into $\Sigma^+ \Sigma^+$. They would be I=S=2 particles. Bag model calculations[4] put the mass of the H^{++} in the region of 2.7 GeV/c².

In the MEM the process $pp \rightarrow \Sigma^+ \Sigma^+$ is only possible if the exchanged meson is or becomes its antiparticle. In this picture a maximum estimate for the cross section

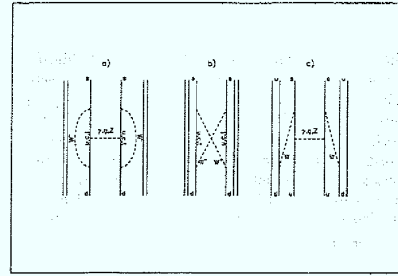


Figure 1 Diagrams for $\Delta S = 2$ transitions in proton collisions. Quark lines are full, bosons lines are broken. The exchange of at least two W bosons is compulsory in the standard model.

is provided through the reaction $\bar{p}p \rightarrow \bar{\Sigma}^+ \Sigma^+$. Here the particles of one arm of the process $pp \rightarrow \Sigma^+ \Sigma^+$ are replaced by their antiparticles. Cross sections for this reaction are of the order 1 to 10 μb . One could also imagine a dissociation of one proton into a Σ^+ and a virtual K^0 . To get absorbed by the other proton the K^0 has to oscillate into its antiparticle $\bar{K}^0$. The $K^0 - \bar{K}^0$ oscillation frequency is proportional to the mass difference $\Delta m = (3.522 \pm 0.016) \cdot 10^{-6} \text{ eV}$ between K_s and K_L . The oscillation period is about $\frac{1}{3}$ nsec corresponding to an oscillation length of a few cm. In hadronic interactions as considered here the interaction length is of the order of 10^{-13} cm. The fraction of K^0 's changed into their antiparticles is $1 - \exp\{-\frac{l}{l_{cm}}\}$ or negligible.

In a very simple picture one would use the cross sections of the reaction $\bar{p}p \rightarrow \bar{\Sigma}^+ \Sigma^+$ and scale it by the ratio of the involved coupling constants. A suppression factor of about 10^{-12} is obtained. The question whether and how much the "double penguin" diagram of figure 1(a) might contribute would be very interesting and there is theoretical evidence that the penguin diagram might enhance flavor changing transitions in baryons [6]. In a more recent calculation [7] it was found that in next to leading order the $\Delta S = 2$ effective Hamiltonian gets a huge positive correction at low momentum transfer.

As a first approximation to an experiment to measure the reaction $pp \rightarrow \Sigma^+ \Sigma^+$ the experiment PS185 at LEAR/CERN can be taken. It has investigated hyperon production in proton-antiproton annihilations. The event signature for $pp \rightarrow \Sigma^+ \Sigma^+$ and $\bar{p}p \rightarrow \bar{\Sigma}^+ \Sigma^+$ is -besides the opposite charge for one of the particles- exactly the same. Near the reaction threshold the hyperons are very slow and thus should be clearly distinguishable from lighter charged particles by a dE/dx measurement. Using a threshold Cerenkov the slow Σ^+ could be retained while faster protons or pions would be vetoed. A further identification of the final state hyperons is possible by the decay $\Sigma^+ \rightarrow (p\pi^0 \text{ or } n\pi^+)$. The kink in the charged particle track

where the hyperon decays can only be faked by angle scattering and therefore a massless detector is desirable. A very good tracking resolution is required. A trigger would look for charged multiplicity 2 events. It could check coplanarity and use online dE/dx and Cerenkov information. The detection of neutrals (either π^0 or n) would help very much in the event identification. Background is coming mainly from prompt pion production (e.g. $pp \rightarrow p\pi^0$ or $pp\pi^0\pi^0$) and can be eliminated online by the recognition of decay kinks.

The threshold momentum for the reaction $pp \rightarrow p\Sigma^+K^0$ is at 2.24 GeV/c. The cross sections for this reaction are of the order of $30\mu\text{b}$. It has one decay kink and might fake a good event due to small angle multiple scattering of the proton. But this reaction is interesting in itself and thus might be a welcome present, as long as it can be identified unambiguously in the offline analysis.

The reaction $pp \rightarrow p\Sigma^+K^0$ also provides a " K^0 beam". By the propagation of the K^0 over a few mm a sizeable oscillation into $\bar{K}^0$ is obtained. An initial luminosity of $0.1\mu\text{b}^{-1}$ leads to a K^0 "beam" intensity of a few Hz. The mixing is around 5% inside a target of 1 cm^3 . The cross section $K^0p \rightarrow \pi^0\Sigma^+$ is about 1 mb at 300 MeV/c Kaon momentum. That would lead to about 1 event per day, but there are three protons involved in the reaction and that can be used in the trigger for identification.

To search in a wide mass region beam momenta up to several GeV/c might be used at a proton synchrotron like COSY. A liquid hydrogen target of $\frac{1}{4}\text{ cm}$ length can stand 10^7 protons per second. This leads to a luminosity of about $0.1 \cdot 10^{30}\text{ cm}^{-2}\text{sec}^{-1}$. With this basic luminosity count rates of 1 event per 2 days can be expected if a hypothetical resonance cross section is 50 pb. The ToF detector[8] at COSY would be an ideal instrument to detect the final state particles, because it provides particle identification via time of flight and dE/dx and has a massless volume outside the target. To measure dE/dx of the final state hyperons and to detect their decay kink an arrangement of thin μ strip detectors is conceivable. Due to the forward peaking event structure the detector system has highest acceptance even if only sensitive in the forward region. Nevertheless the reconstruction efficiency might be as low as 20%[5]. Most triggers will be caused by background from multi-pion production or single kaon production. A trigger based on dE/dx and charged multiplicities as well as being geometrically confined to the kinematically allowed regions will cut down raw rates. Requiring coplanarity and vetoing on charged particles outside the acceptance of the trigger will reduce furthermore. At this level an online decay-kink determination with a fast processor should be possible.

The possibility to study the reaction $pp \rightarrow \Sigma^+\Sigma^+$ has been discussed. It would be the first observation of the flavour changing $\Delta S = 2$ transition outside the neutral meson system. Expected rates are extremely small and there is only little hope to use this reaction channel to search for dibaryonic resonances. Only at known resonance positions one might expect some cross section, if

the strangeness content of that resonance is large.

Even though the expected cross sections are very small, there is much room for new and exciting physics. A nonresonant cross section would be a direct measurement of a flavour-changing neutral current. The proof of existence of the H^{++} would be a strong constraint to quark-gluon models. Its isospin 2 would be a unique identification as a nonstandard particle.

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Systematic study of the pn-CCD detector for the high resolution Bragg Spectrometer

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The spectroscopy of antiprotonic X-ray transitions is the only possibility to investigate the hadronic interaction at threshold. The measurement of the energy shifts and the line widths of $K\alpha$ and $L\alpha$ transitions in hydrogen and helium isotopes yields the corresponding complex scattering lengths and volumes. To perform this experiment at the antiproton storage ring LEAR of CERN, a specially optimized Bragg reflection spectrometer has been developed [1]. An essential component of this device is a two dimensional position sensitive detector with high position resolution and good energy resolution for efficient background reduction. For this purpose a CCD detector is ideally suited. Therefore, we have adapted to our requirements [2,3] a CCD, which was specially developed for X-ray astronomy applications by the MPE München [4,5].

Briefly, it consists of the CCD chip, the amplifier chip CAMEX and a timer chip TIMEX mounted on a Al_2O_3 carrier, which is connected to a copper frame cooled by a LN_2 reservoir. The whole unit is mounted inside a vacuum chamber in a way so that an adjustment in lateral direction and around a vertical axis through the surface of the CCD is possible. Thus an optimal position on the Rowland circle of a focussing spectrometer can be achieved.

The data are transferred serially to an ADC and event filtering unit residing in VME environment. For each pixel, a 32 bit word is formed that contains energy, position, and CCD identification. To handle the data flux which is of the order of 500 kByte/s, an online code has been developed which rejects all pixel rows except the ones with valid information. The corresponding software package for the event filter was developed on the basis of assembler and C, residing in the VME system [6]. Finally the net data are transferred via ethernet to the vax cluster where they could be stored or visualized. The system is described in detail elsewhere [2,3,6].

For the analysis, the following effects have to be dealt with:

- 1) Determination of the individual offset of each pixel.
- 2) Recognition of all pixels that are defect by means of a noise analysis.
- 3) Common mode corrections: This offset displaces the base line of a row of pixels in an arbitrary way.
- 4) Determination of the amplification of each of the 64 amplifier channels.
- 5) Correction for transfer efficiency.
- 6) Charge splitting: The charge created by one X-ray can be distributed into more than two neighbouring pixels. By checking the neighbouring pixels of a detected event the total energy can be reconstructed, but with decreased energy resolution.

1) and 2) must be performed before starting a measurement, because the parameters are already needed for the event filter.

The test measurements were performed with a collimated (200 μm wide) ^{55}Fe source. The following characteristic properties of the present CCD setup could be extracted:

exposure time:	150 ms
readout time:	12 ms
energy of Mn $K\alpha$	5895 eV
FWHM of sum spectrum	240 eV
spread of amplification	1,1%
transfer efficiency per pixel*	(999,5 $\pm$ 0,1) %
probability for charge splitting	(46 $\pm$ 4) %

* in our case, the charge is shifted through up to 200 pixels

A spectrum of ^{55}Mn $K\alpha$ and $K\beta$ for one pixel after application of the common mode correction is shown in fig. 1. In addition, only single hits are required. An energy resolution of 3.7% was achieved which is as good as for a small Si(Li) semiconductor detector.

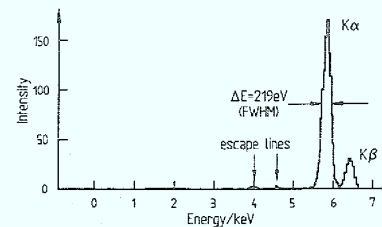


Fig. 1: Energy spectrum of the Mn K X-rays as obtained from a single pixel corrected for common mode and charge splitting.

Meanwhile, a new generation of pn CCDs is available [7]. It allows exposure times around 1s with drastically improved transfer efficiency which yields energy resolutions beyond the limit of good Si(Li) detectors.

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X-rays from Light Antiprotonic Atoms Measured with The Low-Energy Bragg Spectrometer

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The aim of LEAR experiment PS207 is the high resolution spectroscopy of Lyman and Balmer transitions of the antiprotonic hydrogen and helium isotopes to determine hadronic shifts and broadenings of the low-lying atomic hyperfine states [1]. In the present experiment, for the first time, antiprotonic X-rays have been measured using a crystal spectrometer.

The 105 MeV/c antiproton beam of LEAR was injected into the cyclotron trap and stopped to about 90% in hydrogen, deuterium, helium, and neon gas for 30 mbar hydrogen and pressures of equivalent stopping power. With the present beam intensities, about 10^9 antiprotons could be stopped per hour. As Bragg crystals spherically bent quartz and Si crystals of 100 mm diameter and a radius of curvature of about 3 m have been used [2]. The reflected X-rays were detected in a charge-coupled device (CCD) with a pixel size of $22.5 \mu\text{m} \times 22.5 \mu\text{m}$, which provides an excellent position resolution. In addition, the good energy resolution (about 180 eV at 1.7 keV) allow an efficient background rejection.

The width of the X-ray source provided by the stop volume in the center of the cyclotron trap was found to be ≤ 20 mm (FWHM) in diameter from a preliminary analysis. The measurement was done with the $L\alpha$ X-rays from antiprotonic hydrogen and the $5 \rightarrow 4$ transition of antiprotonic helium.

Fig. 1 shows the $L\alpha$ transitions of antiprotonic hydrogen and deuterium as obtained from the online analysis, where a cut is applied to the energy spectrum around the transition energy of 1.738 and 2.317 keV, respectively. An overall detection efficiency for one $L\alpha$ X-ray from hydrogen or deuterium of about $1 \cdot 10^{-8}$ per incoming antiproton was achieved for the whole setup. The absolute yield of the antiprotonic $L\alpha$ transition is about 50% both for hydrogen and deuterium.

The natural widths of the antiprotonic transitions are orders of magnitude smaller than the natural widths of fluorescence X-rays excited by means of X-ray tubes. Usually, fluorescence X-rays are the only possibility to investigate the performance of Bragg crystals. In the case of silicon and sulphur, the natural line widths of the $K\alpha$ transitions are

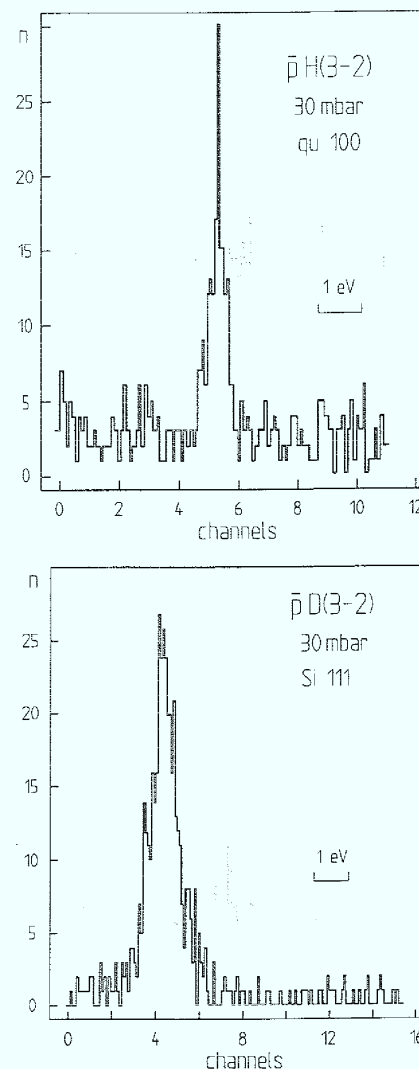


Fig. 1 Position spectra of the CCD showing the $L\alpha$ transitions of antiprotonic hydrogen [$E=1.7368$ keV] and deuterium [$E=2.3165$ keV]. The X-rays were reflected in first order by a spherically bent quartz or silicon crystal.

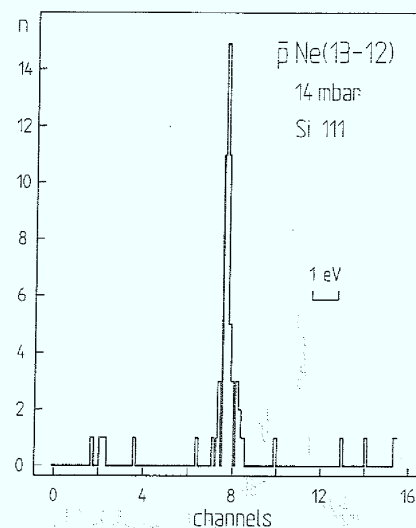
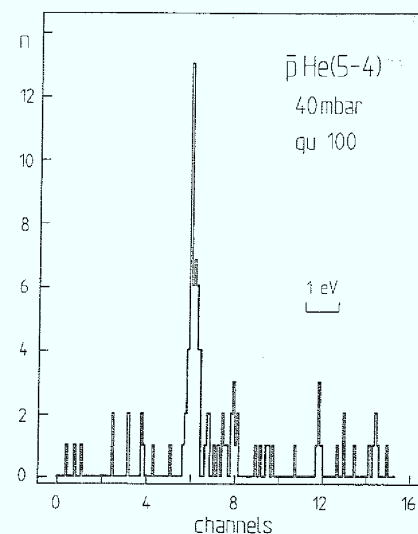


Fig.2 Position spectra of the CCD for the transitions $5 \rightarrow 4$ of antiprotonic ${}^4\text{He}$ [$E=1.7986$ keV] and $13 \rightarrow 12$ of antiprotonic ${}^{20}\text{Ne}$ [$E=2.4441$ keV]. These narrow transitions are used to determine the response functions of the apparatus for the measurements for the hydrogen isotopes.

0.52 and 0.77 eV and thus exceeding the width of the rocking curves by a factor of about 4 [3].

Ideally, the response function of the Bragg spectrometer is obtained from antiprotonic X-rays from atomic states not affected by strong interaction, finite size effects, or electron screening. The natural width of such X-ray transitions is due to a pure electric dipole transition and, therefore, only of the order of a few meV. Any broadening of the reflection width then only contains contributions of the rocking curve, the geometrical aberration, and the position resolution of the detector.

The response functions at Bragg angles close the setup for antiprotonic hydrogen and deuterium $L\alpha$ have been measured with the transitions $5g \rightarrow 4f$ from antiprotonic ${}^4\text{He}$ at 1.797 keV and $13 \rightarrow 12$ from antiprotonic ${}^{20}\text{Ne}$ at 2.444 keV at equivalent pressures corresponding to 30 mbar hydrogen (fig. 2). In the case of antiprotonic ${}^4\text{He}$, it was possible to observe also the parallel transition $5f \rightarrow 4d$. The yields of such parallel transitions provide new experimental information on a region of the atomic cascade which was accessible up to now only indirectly from cascade calculations.

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The aim of experiment R-94.01 presently being performed at the Paul-Scherrer-Institut (PSI) is to determine the mass of the negative pion to an accuracy of $< 3\text{ppm}$ by measuring characteristic X-rays of pionic and myonic atoms with a Bragg crystal spectrometer [1]. A myonic transition of almost the same energy as the pionic transition, i.e., of almost the same Bragg angle, will be used as calibration. Assuming CPT invariance, the myonic transition energy is calculable with the same accuracy as for the mass of the positive myon. If furthermore, the level energies are not affected by strong interaction, finite size effects, or electron screening from the remaining electrons in the exotic atom, the myonic transition provides a calibration almost one order of magnitude more accurate than the one available from measurements of fluorescence X-rays.

To suppress electron refilling from neighbouring target atoms, gas targets have to be used. Furthermore, the X-rays have to be emitted from those levels where the electron shell is already stripped off [2]. In the case of pionic nitrogen and myonic oxygen, various transitions exist which are accessible with a Bragg spectrometer (tab. 1). At a pressure below 10 bar, the probability that there is still one K electron left in the exotic atom is predicted to be a few per cent only [3].

	E / keV		E / keV
$\pi\text{N}(6-5)$	2.2023	$\mu\text{O}(6-5)$	2.1852
$\pi\text{N}(5-4)$	4.0546	$\mu\text{O}(5-4)$	4.0229
$\pi\text{N}(4-3)$	8.7618	$\mu\text{O}(4-3)$	8.6920

Tab. 1 Transition energies for pionic nitrogen and myonic oxygen as calculated from the Klein-Gordon and the Dirac equation. For the final analysis, the QED corrections have to be taken into account.

For a first measuring period, a specially designed Bragg spectrometer was installed in the πE5 area at PSI. The cyclotron trap was used to provide the X-ray source by stopping the pion beam and the myons from pion decay in the center of the target chamber of the trap. The stop density was optimized with Ne, N_2 , and O_2 gas filling by means of a Ge semiconductor detector for

recording pionic X-ray transitions. A stop rate of $3 \cdot 10^6 \pi/\text{s}$ was achieved at 2 bar Ne, which corresponds to a stop efficiency of about 2% for pions.

The stop distribution in the center of the target was measured with the $\pi\text{N}(4-3)$ transition (fig. 1). The X-rays were reflected by a spherically bent Si (111) crystal of 10 cm diameter and a radius of curvature of 3 m [4]. As position sensitive detectors, a multiwire proportional chamber and charge-couple devices (CCD) have been used.

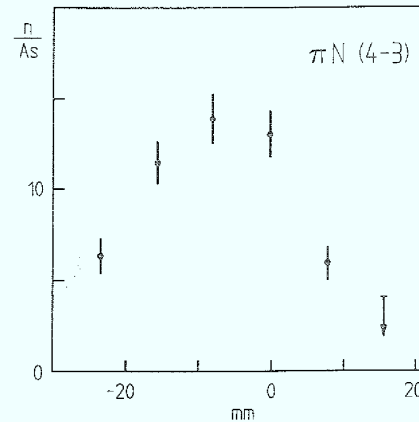


Fig. 1 Stop distribution of pions at 3 bar N_2 measured using the $\pi\text{N}(4-3)$ transition. The X-rays were reflected in forth order by a spherically bent Si (111) crystal and detected in a multiwire proportional chamber filled with a Xe/ CH_4 mixture (90%/10%) at 2 bar. The diameter of the window of the target chamber is 52 mm, i.e., the stop density distribution is uniform in the center of the cyclotron trap.

Negative pion beams produce a large neutron induced background, because each stopped pion produces on the average two or more fast neutrons after being captured in a nucleus. Therefore, a large effort has been spent to suppress the background by installing concrete and lead shielding around the crystal spectrometer. The optimization was done using the multiwire

proportional chamber reaching a peak to background ratio of 1:1 for the $\pi N(4-3)$ transition. A count rate of 40/hour was obtained in the maximum of the stop distribution. With this shielding and using the CCD as position-sensitive detector, a peak to background of about 30:1 was achieved (fig. 2).

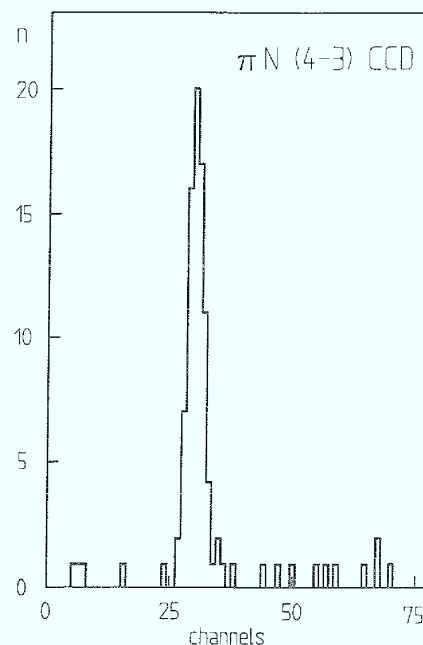


Fig. 2 Position spectrum of the sum of CCD 1-4 at a target pressure of 3 bar N_2 (measuring time: 6 hours, 1 channel = 225 μm or 323 meV).

A count rate of about 10/hour was obtained for the $\pi N(4-3)$ transition at 3 bar target pressure and using the CCD as position-sensitive detector. After a preliminary analysis a width of the reflection of 0.7 eV (FWHM) was obtained (fig. 2). The rate expected for the $\mu O(4-3)$ transition is an order of magnitude smaller. Only a few hours of running time were available for the myon measurement because of the breakdown of the production target E_1 and therefore no clear signal could be observed.

The feasibility study shows that assuming sufficient statistics of ≥ 1000 counts both for the myonic and the pionic transitions, an accuracy of ≤ 3 ppm should be achievable for the mass ratio. The measurements will be continued using the 5-4

transitions of pionic nitrogen and myonic oxygen. Here, the energy splittings for the pionic and the myonic transition due to the presence of a K electron are 0.5 and 1.1 eV, respectively, and can be resolved unambiguously with the crystal spectrometer.

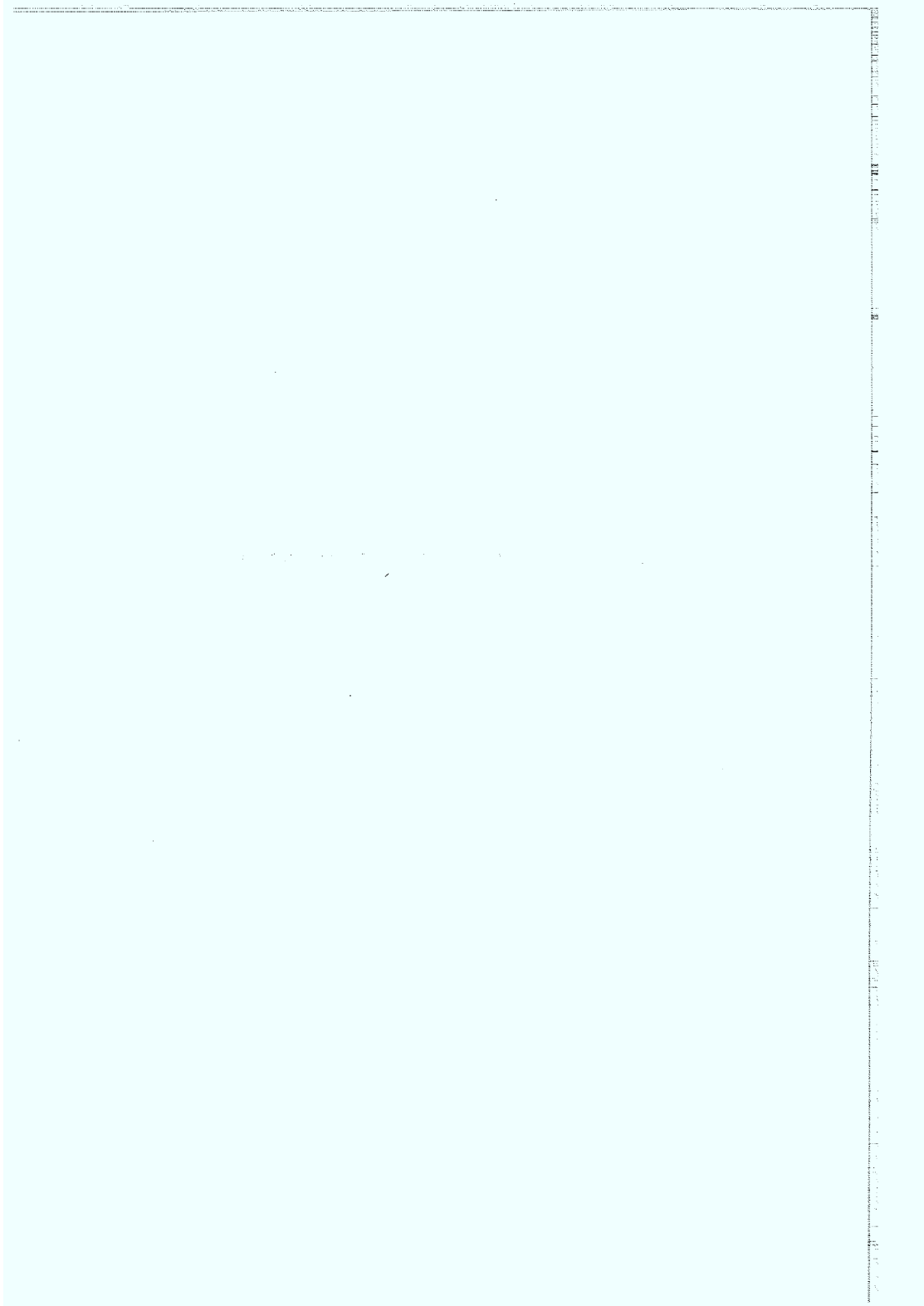
Received: 1990-01-15
 Fully revised: 1990-02-15
 Accepted: 1990-03-15
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2. NUCLEAR SPECTROSCOPY



Candidates for Linking Transitions out of the Superdeformed Band in ^{144}Gd

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We performed three experiments at the GASP spectrometer, Legnaro, in order to study superdeformation in the nucleus ^{144}Gd . All experiments were done with the reaction $^{48}\text{Ti}+^{100}\text{Mo}$ using as target a layer of three self-supporting thin foils with a total thickness of 1.1 mg/cm². The beam was provided by the Tandem XTU accelerator of the Laboratori Nazionali di Legnaro and the γ -rays were detected using the GASP array, an EURO-BALL Phase I spectrometer consisting of 40 high-efficiency anti-Compton spectrometers and an inner ball of 80 BGO scintillators which acts as sum-energy and γ -ray multiplicity filter. The first experiment was done at a beam-energy of 221 MeV whereas the second and third experiments were carried out with a beam energy of 215 MeV which led to a stronger population of ^{144}Gd .

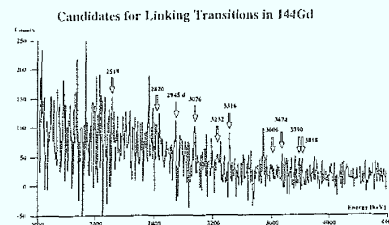


Figure 1: Sum of double gated coincidence spectra showing the SD band in ^{144}Gd . Gates have been set in the cube produced from triples data on all transitions of the band.

In all experiments events were collected when at least three suppressed Ge detectors and three inner ball detectors fired in coincidence. With such conditions the event rate was ≈ 6.5 kHz with a beam current of 4 pA and a 10 kHz singles rate in the Ge detectors. We sorted the calibrated events into RADFORD γ - γ -cubes, unfolding higher fold events into triples while setting proper gates on the BGO multiplicity and sum-energy and the time in order to reduce the background and to enhance the selected reaction-channel.

Analyzing the first experiment we could observe a superdeformed band in ^{144}Gd which shows for the first time a bandcrossing due to the alignment of a $N = 6$ proton pair [1]. In the second experiment it was possible to extend the superdeformed band by at least one transition (fig. 1), while the third experiment is currently analyzed. The predicted second bandcrossing at $\hbar\omega \approx 0.75$ MeV which is related to the occupation of the aligned $N = 7$ neutrons could not be seen in the first and second experiments. The third experiment with remarkable higher statistics aims to an extension of the superdeformed band up to or even above the predicted second bandcrossing.

A very interesting and still open question is the deexcitation of superdeformed states to normal deformed states.

The observed superdeformed band in ^{144}Gd shows the typical feature of decreasing intensity over a very few transitions at the low-spin end. The decay-out of the band occurs over three transitions and starts just at the band-crossing [1].

Several attempts to establish experimentally linking transitions between superdeformed and normal deformed states have been in vain. From the theoretical point of view [2] this is explained by the very large number of possible decay paths and therefore the extremely weak intensity of single transitions. Utilizing the fact that a significant part of the decay paths could go only through one intermediate state Atac et al. [3] summed up each two gamma energies in coincidence with the SD band in ^{143}Eu . It was possible to build up several two-step decay paths that fit into the previously known level scheme.

Using the data obtained from our $^{48}\text{Ti}(^{100}\text{Mo},4n)^{144}\text{Gd}$ experiment at GASP, we performed a similar search. Double gates on the members of the superdeformed band were set and the energies of coincident gammas were summed pairwise. About 350 million events with fold four or higher were used for this search.

In the produced spectrum all possible peaks related to sums of known superdeformed transition energies could be seen. Some peaks in the region from 2.5 to 3.5 MeV showing a proper width of ≈ 9 keV as predicted for decay transitions could be observed (fig 2).

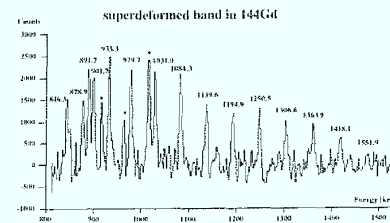


Figure 2: Sumpeaks related to two-step decay transitions in ^{144}Gd . The peaks in the left part of the spectrum are identified as sums of transitions between SD states. Three peaks are remarkable, these are the 2945 keV, 3076 keV and the 3316 keV peaks which show the predicted width of ≈ 10 keV.

It was possible to fit eleven possible two-step transitions into the level scheme with good accuracy. Especially three clear peaks, including one double peak, are considered as linking transitions. The results from the second experiment confirm the obtained peaks and energies from the first experiment [4].

Using this information it was possible to establish an absolute energy assignment for the superdeformed band (fig. 3).

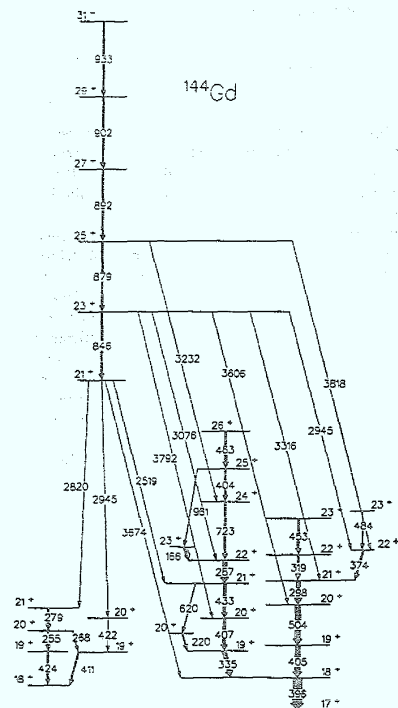


Figure 3: Part of the level scheme of ^{144}Gd [5] which shows the proposed connection between the SD-band and the normal deformed states.

Spins of 19, 20 or 21 $\hbar$ for the lowest SD level are possible considering the selection rules assuming an E1 character for most of the decay transitions. The assignment with a spin of 21 $\hbar$ is preferred as it does not make use of transitions from lower to higher spins. This result is in good agreement with theoretically predicted spins [2] in the $A \approx 140$ mass region.

Since the sum peaks, representing the connecting transitions, obtained in the first two experiments have not a sufficient statistical significance the described search will be repeated with the data obtained in the third experiment of high statistics on ^{144}Gd .

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Proton and Neutron Band Crossing in the Superdeformed Nucleus ^{145}Gd

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The experimental investigations of the $Z=64$ isotopes $^{144,146-150}\text{Gd}$ [1-7] proved the existence of one or several superdeformed bands in each of these nuclei. For the interpretation of the features of the superdeformed bands, like band crossings and the identical band phenomenon, it is important to complete the systematics and to study superdeformation in ^{145}Gd .

High-spin states in ^{145}Gd were populated in the $^{114}\text{Cd}(^{36}\text{S},n)$ reaction at a beam energy of 182 MeV. The beam was delivered by the Tandem XTU accelerator of the Legnaro National Laboratory and the γ -rays were measured with the GASP array which consists of 40 large escape-suppressed Ge detectors and an inner ball of 80 bismuth germanate (BGO) scintillators acting as a γ -ray multiplicity and sum-energy filter. The target consisted of a self-supporting foil of ^{114}Cd enriched to 99.2 % with a thickness of 1.07 mg/cm². The recoil velocity was $v/c = 2.5$ %. Events were collected when ≥ 3 escape-suppressed Ge detectors and > 8 BGO scintillators detected a γ -ray. The rate of such events was ≈ 8 kevents/s for a beam current of 3 particle-nA. In total 2.3 Giga events have been collected. The nucleus ^{145}Gd was produced the strongest through the $5n$ channel. The remaining events belong mainly to the $4n(^{146}\text{Gd})$, $6n(^{144}\text{Gd})$, $p5n(^{144}\text{Eu})$, and $\alpha 4n(^{142}\text{Sm})$ channels. The ^{145}Gd channel was enhanced in the sorting procedure by requiring that at least 15 BGO detectors of the inner ball fired. Furthermore, by placing an appropriate window on the time spectrum, only prompt events have been accepted in the sort. After unpacking $1.3 \cdot 10^9$ three-fold events were available for analysis in a fully symmetrized cube.

In order to search for superdeformed bands a γ - γ matrix has been produced. In this matrix strong ridges with an energy separation of 54 keV in the energy range from ≈ 700 to ≈ 1500 keV have been found. Also a second and third ridge could be identified. This indicates the presence of pronounced superdeformed bands. Discrete superdeformed transitions have been searched for in the matrix through gating channel by channel. In this way three strong superdeformed bands have been identified. In order to increase the selectivity the cube has been analysed subsequently. One of the superdeformed bands has been assigned to ^{142}Sm [8]

The other two superdeformed bands are shown in fig. 1. They were obtained by setting double gates on the strongest transitions of the superdeformed bands. Both bands can clearly be assigned to ^{145}Gd since the strong transitions of 794.9, 924.6 and 1552.6 keV between normal deformed states in ^{145}Gd [9] can clearly be seen, whereas γ -transitions belonging to the neighbouring nuclei disappear. The strongest band, consisting of 18 members, is considered to be the yrast superdeformed band. The weaker first excited superdeformed band consists of 17 transitions.

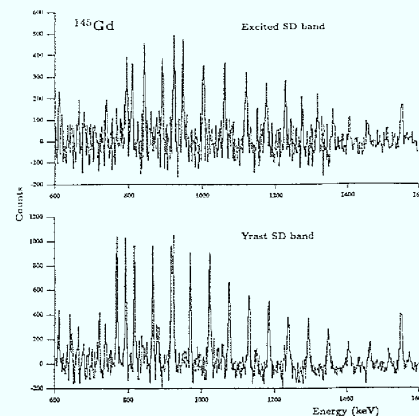


Figure 1: The yrast and excited superdeformed bands in ^{145}Gd .

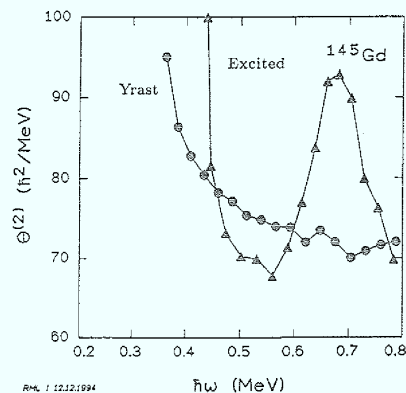


Figure 2: The dynamic moments of inertia as function of the rotational frequency for the superdeformed bands in ^{145}Gd .

The dynamic moments of inertia of the superdeformed bands in ^{145}Gd are plotted as function of the rotational frequency in fig. 2. It can be seen that the yrast superdeformed band starts with a dynamic moment of inertia of $95 \text{ h}^2/\text{MeV}$ which decreases smoothly to a value of $\approx 71 \text{ h}^2/\text{MeV}$ similarly to the values observed in the neighbouring nuclei $^{144,146,147,148}\text{Gd}$ [1-4,7] outside the respective band crossing regions. The first excited superdeformed band, however, shows two subsequent band crossings at

rotational frequencies of ≈ 0.4 and 0.68 MeV, respectively. This is the first time that two band crossings have been observed in one superdeformed band. It should be stressed that the pronounced low-frequency crossing is similar to that observed at a rotational frequency of 0.45 MeV in ^{144}Gd whereas the second band crossing has similar features as that found in ^{146}Gd at 0.65 MeV.

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Study of Superdeformation in ^{146}Gd

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In ^{146}Gd , two superdeformed bands are known [2,3], which are amongst the weakest SD bands found in this mass region. The earlier investigations had been performed using mainly sum spectra with poor statistics. The aim of the experiment we report here was to study these bands with sufficient statistics in order to prove by means of individual gates the mutual coincidence relations of the band transitions and to determine the transition energies more accurately. In addition a search for further excited SD bands in ^{146}Gd was planned.

Excited states in ^{146}Gd were populated by the $^{102}\text{Ru}(^{48}\text{Ca},4n)^{146}\text{Gd}$ reaction at a beam energy of 203 MeV. The beam was supplied by the tandem Van de Graaff accelerator at the Nuclear Structure Facility, Daresbury. The target consisted of two self supporting ^{102}Ru foils of about 0.5 mg/cm² thickness mounted with a separation of approximately 1 mm. Gamma rays were detected with the EURO-GAM I spectrometer, comprising 45 large volume Compton suppressed germanium detectors. In total $7 \cdot 10^8$ events with an unsuppressed fold ≥ 6 were recorded.

To check the coincidence relations of the transitions and to determine the transition energies precisely, all events with a G γ -fold ≥ 3 (fold ≥ 4) were sorted into $2k \times 2k$ matrices for both bands with the additional condition that at least

one (two) of the coincident transitions belong to the SD band.

Fig. 1 shows examples of gated spectra where in each case the gate was set on an individual SD transition in the single gated matrices. The background was removed by subtracting the product of the matrix projections normalized to the total number of counts from the raw matrix.

These spectra show for both SD bands that the transitions are mutually coincident and that one is really observing single cascades. This was not clearly demonstrated up to now especially for band 2 because of the poor statistics of the single gated spectra obtained previously [3].

The transition energies of both bands were obtained from the sum spectra in the double gated and single gated matrices (see fig. 2). Band 1 could be extended by the 1498.5 keV and the 153.6 keV transition. Band 2 was extended by the 1582.0 keV transition. The band crossing in band 1 reported in [2] at $\hbar\omega = 0.65$ MeV can be clearly confirmed.

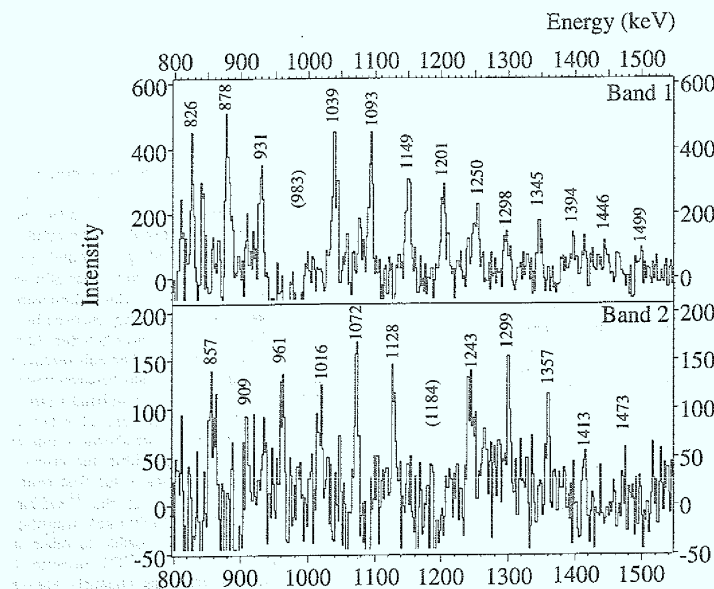


Figure 1: Spectra from single gates on the 983 keV and 1184 keV transitions of band 1 (upper part) and band 2 (lower part), respectively.

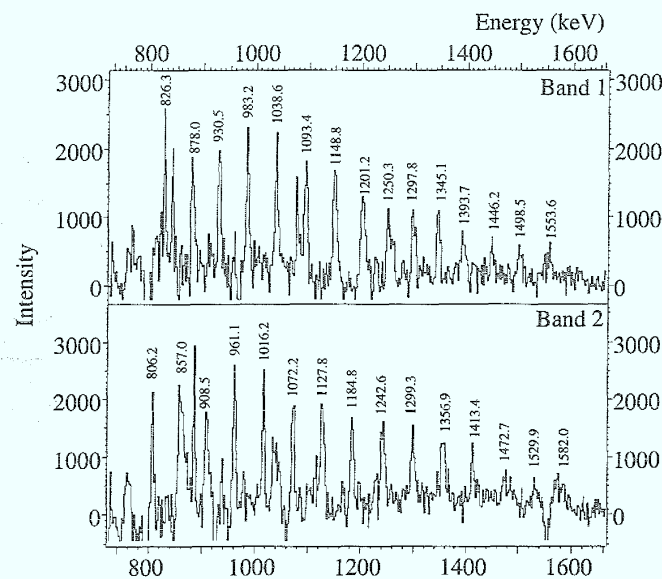


Figure 2: Sum spectra from the 8 cleanest gates of the two superdeformed bands in ^{146}Gd

In order to check that the two bands really belong to ^{146}Gd we compared single and double gated spectra as shown in fig. 3. In the single gated spectrum from a sum gate on band 1 (two fold data, upper part) the stronger transitions of ^{145}Gd (labeled by arrows) also appear in coincidence with the band, while in the sum spectrum of double gated spectra with gates on band 1 (three fold data, lower part) only transitions from ^{146}Gd remain in coincidence with the band. ^{145}Gd is the second strongest reaction channel in our data. The gates were chosen in a way that they did not contain any transitions from the known discrete level schemes of ^{145}Gd and ^{146}Gd . The elimination of ^{145}Gd by the double gating procedure indicates that the remaining peaks correspond to real coincidences and not to the background structure which contains transitions from ^{145}Gd . An assignment to an Europium isotope can be excluded as the small cross section for Eu nuclides in our experiment would imply a SD-intensity of 20 % for this reaction channel which is very unlikely. The situation for band 2 is similar and confirms that this band also belongs to ^{146}Gd . From the coincidence analysis it was found that the normal deformed 27^- and 26^+ levels are most strongly populated from SD states.

To search for new superdeformed bands a grid search method in two [4] and three dimensions was used. For the three dimensional search events were sorted with Ge-fold ≥ 3 into a $1\text{ k} \times 1\text{ k} \times 1\text{ k}$ cube. To remove the background, the generalization of the Palameta-Waddington-formula [5] to three dimensions [6] and various approximations of it were used. In all cases the quality of the three dimensional spectra was much better as could be easily checked by the two known bands. All transition distances between 30 and 70 keV were examined in 2-keV steps. The only discovered structure is a group of 9 equidistant peaks between 959 and 1369 keV with about 1/8 of the intensity of band 1 (see fig. 4). The transitions at 1006, 1065, 1124 and 1176 keV can be observed in individual gates and are shown in this way to be coincident with each other. Within the errors all energies are identical to the energies of the first band of ^{147}Gd [7], but the strongest peaks of the ^{147}Gd -band, namely the 697, 745, 795, 848, and 900 keV transitions, could not be observed. It was not possible to relate this band unambiguously to ^{146}Gd or ^{147}Gd because the contaminants from various reactions channels are very strong. From the intensity it can be estimated that in our experiment the lower limit for detecting new bands is at about 1/10 of the intensity of the first band, and the intensity of any additional band in ^{146}Gd should be below that limit.

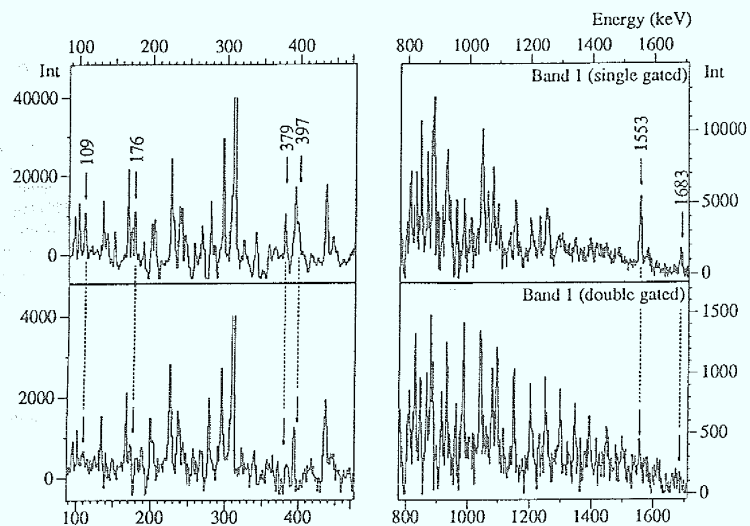


Figure 3: Single gated (upper part) and double gated (lower part) sum spectra of band 1. In the double gated spectrum only transitions of ^{146}Gd are present. The transitions of ^{145}Gd have disappeared (see arrows).

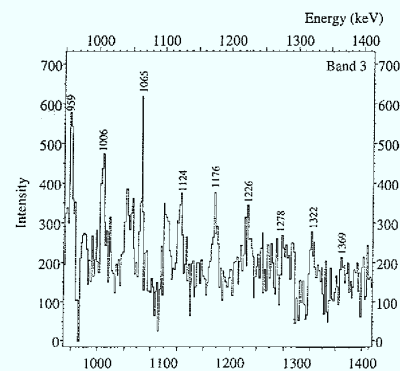


Figure 4: Sum spectrum of band 3.

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In the nuclei above ^{146}Gd the valence protons fill the $\pi h_{11/2}$ shell, and their β -decays are dominated by transformation of an $h_{11/2}$ valence proton to the empty $h_{9/2}$ neutron shell. This Gamow-Teller transition is the only allowed β -decay for the nuclei with 82 or more neutrons. For doubly-even parent nuclei, where a paired $h_{11/2}$ proton will be transformed, the decay proceeds as $(\pi h_{11/2}^2) 0^+ \rightarrow (\pi h_{11/2} \nu h_{9/2}) 1^+$ to essentially one single low-lying 1^+ state in the odd-odd daughter, and for an odd-N parent, with an $f_{7/2}$ spectator valence neutron, the decay mode will be the same but now populates 3qp-states of $(\nu f_{7/2} \times 1^+)_{5/2^-, 7/2^-, 9/2^-}$ character at ≈ 2 MeV excitation in the odd-Z daughter, see Fig. 1.

The situation changes drastically for even-Z parent nuclei with $N \leq 81$, where now the odd valence neutron is either an $s_{1/2}$ or an $h_{11/2}$ neutron hole; both states undergo β -decay. The dominant GT-decay is still $(\pi h_{11/2}^2) 0^+ \rightarrow (\pi h_{11/2} \nu h_{9/2}) 1^+$, but here proceeding to the three-nucleon configurations $(\pi h_{11/2} \nu h_{9/2} \nu h_{11/2}^{-1})_{9/2^-, 11/2^-, 13/2^-}$ or $(\pi h_{11/2} \nu h_{9/2} \nu s_{1/2}^{-1})_{1/2^+, 3/2^+}$, both involving a neutron excitation across $N=82$ and therefore are not expected to occur below ~ 4 MeV excitation in the odd-Z $N=82$ daughter. Moreover, the additional GT transitions $\pi h_{11/2} \rightarrow \nu h_{11/2}$ and $\pi s_{1/2} \rightarrow \nu s_{1/2}$ can now also contribute to the decay, albeit with reduced strength since only a single neutron hole is open. These transitions however proceed to the daughter ground states where the larger phase space will cause increased decay intensity. Moreover, they will not give rise to γ -rays in the daughter nucleus, which makes it difficult to investigate these β -branches in a conventional on-line mass separator experiment using γ -spectroscopic techniques.

In the present study we have investigated the $0^+ \rightarrow 1^+$ GT-decays of the two $\nu s_{1/2}^{-1}$ and $\nu h_{11/2}^{-1}$ isomers of $^{147}_{66}\text{Dy}_{81}$, with the principal aim to populate above-yrast states in the

one-proton nucleus $^{147}_{65}\text{Tb}_{82}$. Of particular interest here are the level multiplets formed by the couplings of the ^{147}Tb valence proton to the low-lying octupole- and quadrupole vibrational excitations of the $^{146}_{64}\text{Gd}_{82}$ are at 1.6 and 2.0 MeV. These ^{147}Tb particle-phonon states cover the spin range from $1/2$ to $17/2$. Most of them occur more than 1 MeV above the yrast line, and even a concurrent dedicated in-beam experiment at the Coulomb barrier with ^6Li beam [1] failed to populate ^{147}Tb above-yrast states with $I < 5/2$. The levels may however be excited in secondary γ -decay following ^{147}Dy GT decay, which essentially can excite all spins up to $15/2$.

Our original ^{147}Dy β -decay experiments [2] utilized HI-compound evaporation production with 4.5 MeV/u ^{58}Ni beam on ^{92}Mo at the GSI on-line mass separator where standard $\gamma\gamma$ - and γX -coincidences and timed γ - and p-singles spectra were measured. Later, additional data were obtained with the high yields of the ISOLDE and ISOCELE II mass separators; a miniorange electron spectrometer, Compton-suppressed γ -detectors, and a cooled Si(Li)-detector for proton measurements were employed in these experiments.

These data gave the decay schemes of Figs. 2 and 3. Firm attribution of the γ -rays to the respective parent activity is obtained from their decay characteristics. While the $11/2^-$ decay shows the pure 56s half-life of the ^{147}Dy M4 isomer, a consecutive 56s $\rightarrow$ 49s decay curve identifies the transitions following $1/2^+$ decay. Else, coincidence data and energy relations were used to build the level schemes.

The two ^{147}Dy decays both reveal quite sharp GT-resonances in ^{147}Tb , which lie 2.33 and 2.42 MeV below their respective $\nu s_{1/2}^{-1}$ and $\nu h_{11/2}^{-1}$ parent states of ^{147}Dy . We note that the $\pi h_{11/2} \nu h_{9/2} \nu h_{11/2}^{-1}$ resonance at 4.70 MeV excitation contains the $h_{9/2}$ - $h_{11/2}$ neutron particle-

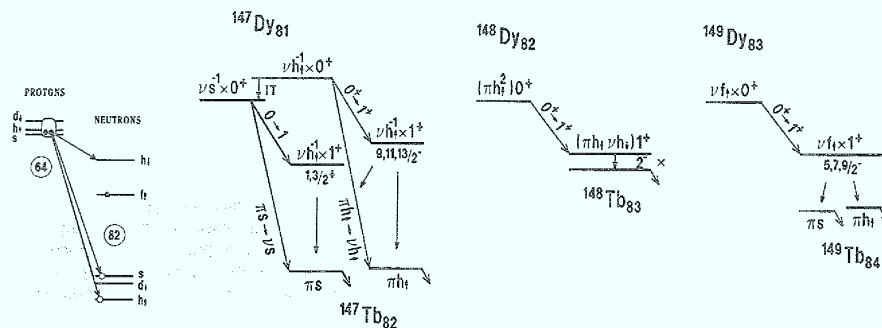


Fig. 1: Schematic GT-decays of even-Z nuclei above Gd with 81 to 83 neutrons. The single-particle diagram shown to the left includes the three GT-transitions occurring in these decays.

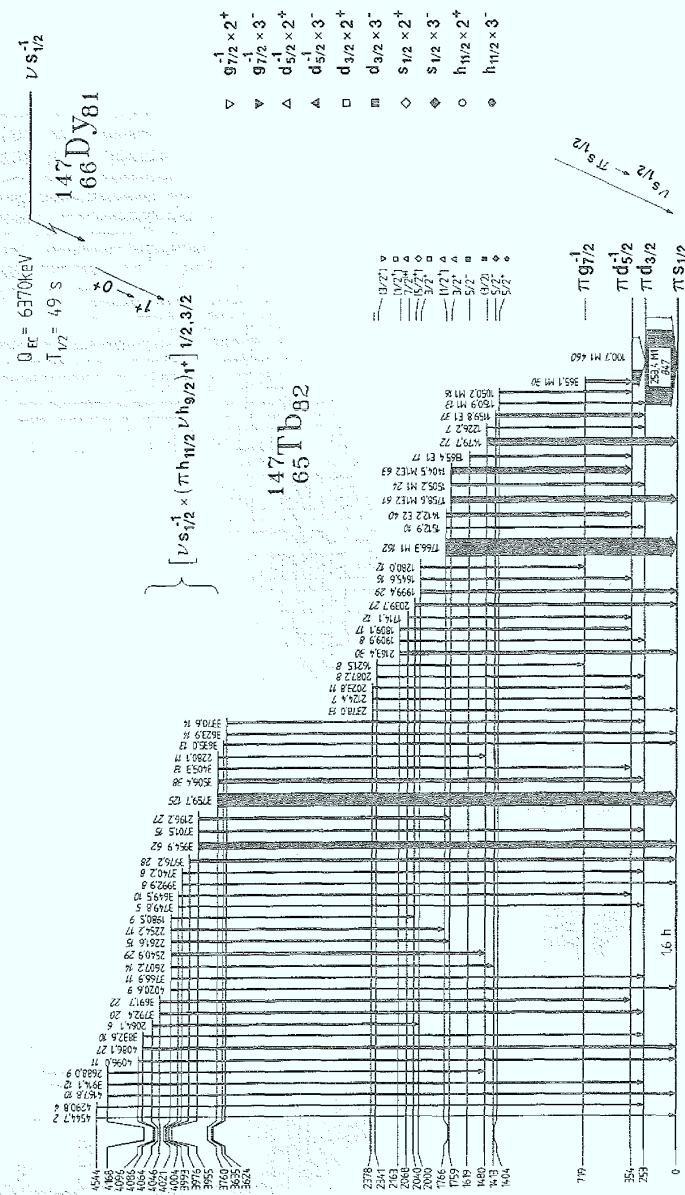
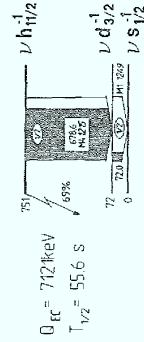


Fig. 2 and 3: Decay schemes of ^{147}Dy $1/2^+$ and ^{147}Dy $11/2^-$. Proposed spin-parities and configurations are specified, where the results of a concurrent in-beam study [1] were also considered. Total transition intensities are given, normalized to 1000 ^{147}Dy $11/2^-$ M4 gamma-transitions (1215 M4-decays). Multipolarities extracted from conversion electron data are specified.



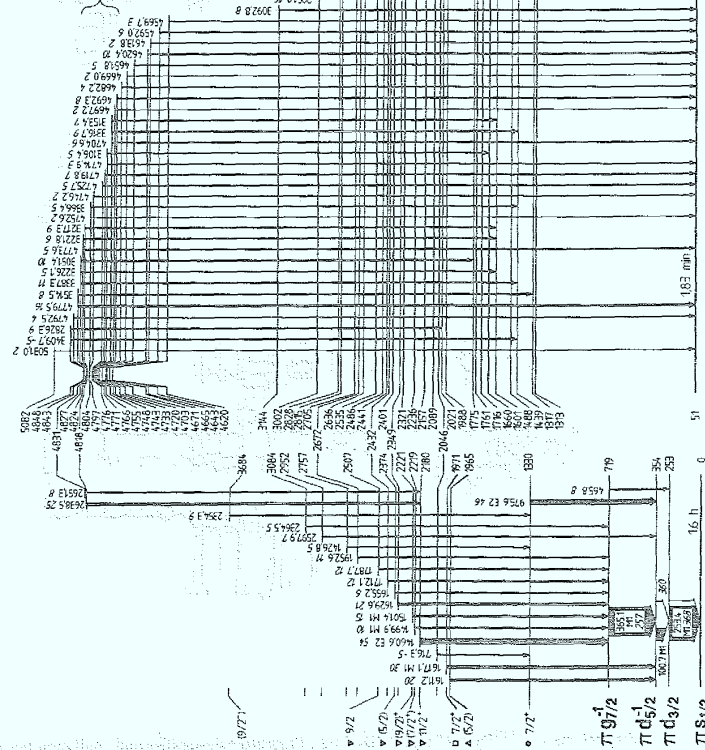
147Dy81
66Dy81

$Q_{EC} = 712 \text{ keV}$
 $T_{1/2} = 55.6 \text{ s}$

$\nu h \frac{1}{2} \times (\pi h \frac{1}{2} \nu h \frac{1}{2} \nu h \frac{1}{2})$

$9/2, 11/2, 13/2$

147Tb82
65Tb82



hole excitation across $N=82$ and thus specifies the splitting of the neutron $i=5$ spin-orbit pair except for contributions from residual interactions.

Finally, we mention that the precise level energies of low-spin resonance states together with the energies of the discrete lines in the proton spectrum have determined the ground state mass of ^{147}Tb [3].

The figure also includes the proposed configuration assignments for the ^{147}Tb levels up to 2.35 MeV. In many cases, angular distribution information from the in beam study [1] was vital to assign the level spins, but several states with $I=1/2$ and $3/2$ could only be observed in the present β -decay experiment. Altogether, the data from the two experiments identified 37 of the total expected 47 states arising from octupole- and quadrupole-vibrations built on the five low-lying single proton states in ^{147}Tb . Of particular interest is the now completely identified $\pi h_{11/2}$

octupole septet where the deviations from the harmonic energy of the individual septet members are found [1] to be in good agreement with fundamental theoretical predictions.

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An interesting structural aspect in one-valence particle nuclei are the energies of the levels formed by the couplings of the valence particle to the low-lying collective excitations in the core nucleus. In ^{146}Gd these are the octupole and quadrupole phonons at 1.58 and 1.97 MeV, and one expects in ^{147}Tb associated particle $\times$ phonon multiplets built on the low-lying single proton states.

Up to now ^{147}Tb levels above 1 MeV were only investigated in yrast experiments using compound evaporation reactions [1,2,3], and therefore the only known particle $\times$ phonon levels are the 1.317 MeV $15/2^+$ and the 2.089 MeV $17/2^+$ members of the $\pi h_{11/2} \times 3^-$ octupole septet. This septet is of particular interest but so far the members with the spins $13/2^+$ and below were only tentatively known (e.g. quoted in [4]). Identification of these states is difficult, since they all lie far above the yrast line. Moreover, at nearby energies the 3^- and 2^+ phonon multiplets built on the four other single proton states $\pi s_{1/2}$, $\pi d_{3/2}$, $\pi d_{5/2}$ and $\pi g_{7/2}$ are also expected, where naturally for the proton hole states the 1678 keV 2^+ level of ^{148}Dy is the quadrupole phonon. This gives rise to altogether 47 particle $\times$ phonon levels with unperturbed energies below 2.41 MeV. As configuration-selective excitation of these states is excluded, reasonably reliable configuration assignments are only possible in cases where at least level spin and parity are determined conclusively from experiment.

Since ^{147}Tb is not stable, nuclear reaction studies of these above-yrast states are essentially possible only in a light ion induced fusion-evaporation reaction where little

angular momentum is transferred to the compound nucleus. The lightest projectile to form ^{147}Tb is ^6Li on ^{144}Sm followed by evaporation of three neutrons. We have used this reaction with 29 MeV beam energy, practically at the Coulomb barrier, where the three strongest exit channels are ^{145}Eu , ^{147}Tb , and ^{148}Tb , with relative formation cross sections of about 3:2:1. Gamma- γ coincidences were measured with the Kölner Würfel Compton-suppressed hexadetector array, and conversion electron data were obtained at the Orsay tandem accelerator with the same reaction using a magnetic electron transport lens in combination with a $\text{Si}(\text{Li})$ spectrometer. These measurements have identified ^{147}Tb levels up to 4 MeV excitation and $I=27/2$. Of vital importance for classification of the weakly-fed lower spin states were data from a concurrent study of ^{147}Tb populated in the β -decay of its two ^{147}Dy $1/2^+$ and $11/2^+$ parent activities [5] which identified $1/2^-$ and $3/2^-$ states that are not accessible in $^6\text{Li}, 3n$ and confirmed many of the levels populated in beam.

The in-beam experiment established below 4 MeV in ^{147}Tb 92 new in addition to the previously known 13 states. Many of these are also excited in ^{147}Dy β -decay where furthermore 20 low spin states not seen in $^6\text{Li}, 3n$ were observed. A more complete discussion of these β -decay results is given in [5,6]. In the present note we only regard the ^{147}Tb particle $\times$ phonon excitations which should all lie below ~ 2.6 MeV. These results are shown in Fig. 1, except for three states close to 2.4 MeV firmly assigned to 3qp configurations [6]. In the figure the levels are grouped into columns for the specific particle-vibration

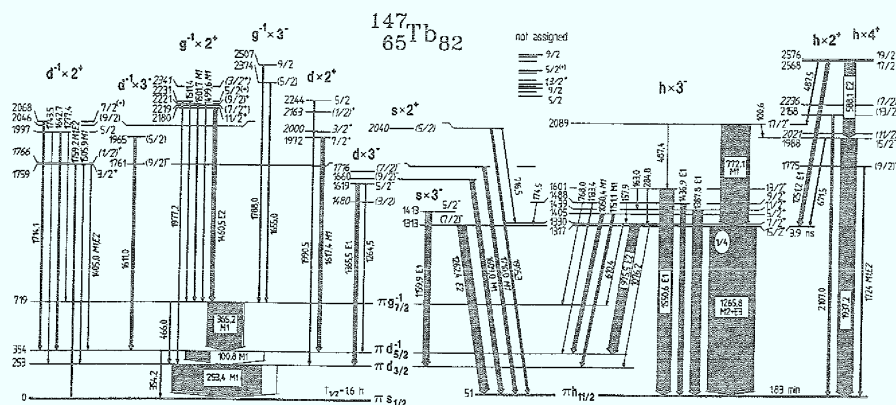


Fig. 1: Particle $\times$ phonon states in ^{147}Tb with transition intensities shown as observed in the ($^6\text{Li}, 3n$) experiment. Results based exclusively on β -decay data are given in italics.

multiplets as specified above. We also note that the decay γ -rays are not shown in the figure for particle $\times$ phonon multiplet members with low spin observed exclusively in β -decay [5] and for the 12 states below 2.6 MeV which remain unassigned.

An important result is the clear-cut identification of the five low-spin $\pi h_{11/2} \times 3^-$ septet members. They are populated with at most 3.4% of the $3n$ -exit channel (and similarly per ^{147}Dy parent decay). But nevertheless in all cases spins and parities are conclusively determined from the observed decay branches, the angular distributions, the conversion coefficients for six high-energy γ -transitions, and moreover from the observation of four intra-septet γ -rays competing with high-energy dipole deexcitation. From these results, and the observed level energies, the configuration assignment for all seven members of this octupole multiplet is firm, which apparently cannot similarly be claimed for all other assignments proposed [6].

The complete $\pi h_{11/2} \times 3^-$ octupole septet is shown in the yrast diagram of Fig. 2. It was already noted [1] that the large energy shifts for the $17/2^+$ and $15/2^+$ members are caused by the exchange of the $\pi h_{11/2}$ valence particle with the dominant $\pi h_{11/2} \pi d_{5/2}^{-1}$ component in the ^{146}Gd octupole phonon as characterised by the diagram in Fig. 2. There, the complete septet is compared with the energy shifts predicted for the low spin members from this process. In agreement with theory, one notes that the shifts of the five lower spin members from the octupole energy are distinctly smaller than for the $15/2^+$ and $17/2^+$ states, and even their energy staggering occurs as expected. But they all lie 130 to 300 keV too low. This additional shift may partly arise from the interaction of the $h_{11/2}$ valence proton with the $h_{11/2}$ proton in the phonon when they are coupled to the lower spins 8^+ , 6^+ , 4^+ , 2^+ , or 0^+ .

In summary, numerous above-yrast states have been identified in the one-proton nucleus ^{147}Tb from $(^6\text{Li}, 3n)$ - and β -decay experiments, and the data have largely characterised their spins and parities. Up to 2.6 MeV the states are in accord with the levels expected from the couplings of the five single-proton states of the 50 to 82 shell with the octupole and quadrupole phonons of the core. Particle-vibration configuration assignments are

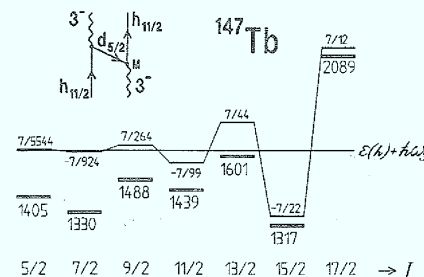


Fig. 2: The observed $\pi h_{11/2} \times 3^-$ septet states (thick lines) of ^{147}Tb , with E^* in keV, compared with energies calculated from the exchange diagram shown in the insert, calibrated with the measured 772 keV $17/2^+$ to $15/2^+$ separation. The fractions are the geometrical factors, which specify the relative energy shift for each spin.

proposed for most of the observed levels; of the two lowest-lying multiplets, $\pi h_{11/2} \times 3^-$ and $\pi s_{1/2} \times 3^-$ all members are unambiguously assigned. Such data provide fundamental information on fermion-boson coupling, and are vital input data for quantitative analysis of complex multiparticle $\times$ phonon structures in neighbouring nuclei as it was shown - using less firm data - e.g. for the three-nucleon $\times 3^-$ couplings in ^{149}Tb [4].

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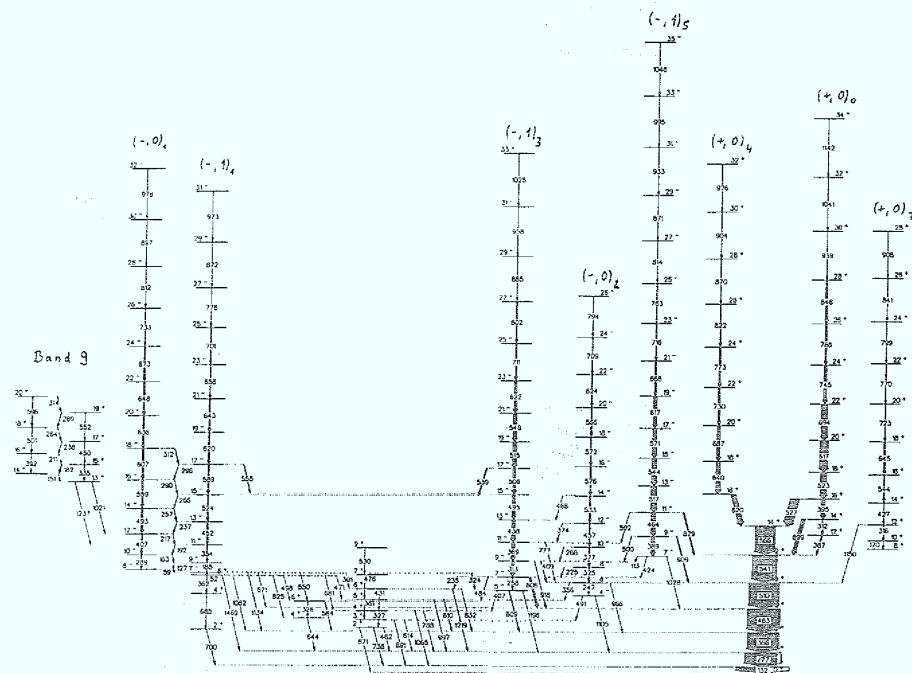
Study of Band Structures in ^{180}Os

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We reinvestigated ^{180}Os because of the experimental evidence for ridge structures [1,2] indicating the existence of a superdeformed shape in this nucleus and the observation of a "trumpet" ridge between 1000 and 1300 keV indicating a strong decrease of the dynamic moment of inertia [2]. Bands in this mass region corresponding to a large deformation were theoretically predicted a long time ago [3-5]. The occurrence of the trumpet ridge is related to a change from a prolate to a strongly triaxial shape ($\gamma \approx -90^\circ$) [6]. Prior to this work the level scheme of ^{180}Os has been studied by Draconulis et al. [7-10] and Lieder et al. [11,11-14].

triple (Ge-Ge-Ge) and higher fold coincidence events were collected. Different matrices were sorted: γ - γ , γ -time, γ -BGO sum-energy, γ -BGO, BGO sum-energy-BGO fold using the SOPHIE sort program [16]. In the analysis a narrow energy dependent time gate was set on the Ge energies and gates were put on BGO fold (≥ 8) and high BGO sum-energy. A total number of 1 Giga triple events were sorted into a Radford-cube and were used in the final analysis. The relation of ^{180}Os to the ^{181}Os and ^{182}Os yields was 70%, 20% and 10%, respectively.

The level scheme of ^{180}Os is shown in fig. 1. The placement of the transitions and levels are based on coincidence



spins. A few remarks pertinent to each band will be given subsequently.

The $(-,0)_1$ and $(-,1)_1$ bands: The band head is the well studied isomer [8,14,17] with a half-life of 26 ns. The new transition of 185.6 keV belonging to the $(-,1)_1$ sequence was identified as a result of the analysis. In the previous measurement [8] using ^{16}O beam on ^{168}Er targets, due to the presence of strong 185.6 keV line originating from Coulomb excitation the placement of this transition was hindered. Three new transitions: 812, 897, and 978 keV in the $(-,0)_1$ band and three new lines in the $(-,1)_1$ band: 778, 872, and 973 keV extending the bands up to spins of 32^- and 31^- , respectively were identified.

The $(-,1)_3$, $(-,0)_2$ and $(-,1)_5$ bands: The observation of the new 185.6 keV transition allows to place the two new transitions of 539.2 and 555 keV between the $(-,1)_1$ and $(-,1)_3$ bands. Four new transitions in band $(-,1)_3$: 802, 885, 958, 1028 keV up to spin 33^- can be seen in fig. 1. The very weakly populated $(-,0)_2$ band has been extended by only two new transitions of 709 and 794 keV up to the spin 26^- . The $(-,1)_5$ band was extended up to spin 35^- with the five new transitions: 814, 871, 933, 995 and 1046 keV.

The positive parity bands: The $(+,0)_4$ band, previously assigned [14] as $(+,0)_6$ band shows four new transitions: 821, 870, 904 and 927 keV up to the level with spin 32^+ . New transitions belonging to band $(+,0)_6$ previously assigned as the $(+,0)_4$ band in ref. [14] were found. The three new transitions: 939, 1042 and 1142 keV extend the band to spin values of 34^+ .

The $(+,0)_7$ band was extended with three new transitions of 799, 841 and 908 keV up to spin 28^+ . The transition 1150 keV which deexcites the band to the 10^+ ground state level appears clear in the coincidence spectrum.

Band 9: A weakly populated strongly-coupled band starting at 3706 keV and deexciting through the 1021 and 1237 keV to the $(-,0)_1$ and $(-,1)_1$ bands, respectively, has been found. We used the DCO ratios of the deexciting transitions to assign the spin I of the band head. Within the limits of our experiments $t_{1/2} \geq 5$ ns no half-lives were measured for the 1021.5 and 1236.8 keV transitions.

The 7^- isomeric state at 1928 keV in ^{180}Os was assigned by [8,11,14,17] as a state with the single configuration $(\nu 7/2^+ [633] \uparrow \otimes \nu 7/2^- [514] \downarrow)$ [8]. Taking into account the experimental g_K factors determined by the branching ratios from the transitions between the two bands $(-,0)_1$ and $(-,1)_1$ built on the isomer the same two-quasiparticle neutron configuration with a two quasiprotone admixture of at least 40% was assigned in ref. [14]. The moment of inertia parameter $\Delta E(I \rightarrow I-1)/2$ vs. the spin I gives a strong support for the configuration assignment of the band head. The placement of the new 185.6 keV transition results into a new spin assignment for the $(-,0)_1$ and $(-,1)_1$ bands. The spin increases by one unit relative to the values in refs. [8,14] and the signatures exchange. This fact gives an excellent overlap with the moment of inertia vs. I plot of the 7^- band in ^{178}W [18] which is definitely assigned as the neutron configuration $7/2^+ [633] \otimes 7/2^- [514]$. The configuration of the isomer is reflected in the structure of its rotational bands. The aligned angular momentum which is shown in fig. 2 support the suggested neutron configuration.

An Interaction strength estimation is obtained by using the two level approximation [19] assuming that the interaction matrix element has a linear spin dependence and mixes only states of the same spin. A value of $V = 18 \pm 0.3$ keV

was obtained for the interaction strength of the 17^- levels from the experimental branching ratios for the 17^- level of the $(-,1)_3$ band at 3982 keV. Using this value the calculated branching ratio for the transitions deexciting the 17^- level at 4032 keV of the $(-,1)_1$ sequence gives $\lambda = 555/589_{th} = 0.29 \pm 0.2$ which is in good agreement with the experimental branching ratios of 0.35 ± 0.10 and 0.22 ± 0.1 from spectra gated on the 620.4 and 643.3 keV transitions, respectively.

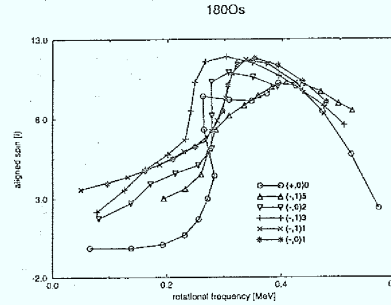


Figure 2: Plots of the experimental aligned angular momenta i as function of the rotational frequency $\hbar\omega$. The parameters of the Harris expansion of the reference band are $\Theta_0 = 24.8 \hbar^2/\text{MeV}$ and $\Theta_1 = 91.1 \hbar^4/\text{MeV}^3$.

The interaction of bands and bandcrossings in ^{180}Os has been analysed in the framework of the cranked shell model [20]. The aligned angular momenta i of these bands are displayed in fig. 2 as function of the rotational frequency $\hbar\omega$. Several band crossings can be seen. In the region $0.26 \leq \hbar\omega \leq 0.30$ the $(+,0)_6$ bands and all side bands, except for the $(+,0)_4$ and $(+,0)_7$ bands display the crossings. The rotational alignment of an $i_{13/2}$ quasineutron pair is considered to be responsible for these band crossings. The low-frequency anomalies observed in the yrast band in ^{180}Os and in several nuclei around [7,8,11,21] are explained [6] by a transition from a partly aligned prolate $\nu_{13/2}^2$ configuration to an aligned $\nu_{13/2}^2$ S-band at a triaxial deformation. A similar explanation for the complex behaviour have been proposed earlier in ref. [11] where it is related to the feature that, due to the large ϵ_4 deformation, a change from a Fermi alignment to rotation alignment [22,23] occurs. A multiple alignment mechanism or shape co-existence have been proposed in refs. [21,24-26].

The $(-,0)_2$ and $(-,1)_3$ bands show signature splitting already for the lowest observed frequencies (cf. fig. 2). These sequences start with aligned angular momenta of $1.8 \hbar$ and $3.6 \hbar$, respectively and acquire values of $4.7 \hbar$ and $6.7 \hbar$ at a rotational frequency of $\hbar\omega = 0.28$ MeV and 0.26 MeV. The features of these bands are similar to the predicted octupole bands by Vogel et al. [27]. Negative parity bands of $\nu_{13/2}^2 \otimes \nu_{7/2}^-$ and $\pi_{11/2} \otimes \pi_{5/2}^-$ configurations should according to Vogel have the character of an octupole-vibrational band at low spins with the octupole angular momentum aligned along the axis of rotation but transform into two-quasiparticle bands at larger rotational frequency.

The sizeable signature splitting and the rapid alignment are features observed also for the low-lying $(+,+)_1$ and $(+,-)_1$ bands in ^{181}Os [28]. It is reasonable, therefore, to assign the lowest-lying $i_{13/2}$ quasineutron orbitals $v(+,+)_1$ and $v(+,-)_1$ having $\Omega^\pi = 9/2^+$ at $\omega = 0$ to the bands $(-,1)_3$ and $(-,0)_2$. The second quasineutron in the configurations may be the $v(-,+)_1$ orbital with $\Omega^\pi = 7/2^+$ at $\omega = 0$ originating from the $7/2^+ [514]$ Nilsson orbital. The signature splitting observed between the $(-,1)_3$ and $(-,0)_2$ bands can only be understood for a triaxial nuclear shape if $v(+,+)_1(-,+)_1$ and $v(+,-)_1(-,+)_1$ configurations are assumed, respectively.

The $(-,0)_1$ and $(-,1)_1$ bands do not show a signature splitting (cf. fig. 2) except for the highest levels. They were assigned to have predominately a singlet $v7/2^+ [633] \otimes v7/2^+ [514]$ configuration, with the quasineutrons $v(+,\pm)_2$ and $v(-,\pm)_1$ which originate from the $i_{13/2}$ quasineutron orbital with $\Omega^\pi = 7/2^+$ and the $f_{7/2}$ quasineutron orbital with $\Omega^\pi = 7/2^+$ at $\gamma = 0^\circ$ and $\omega = 0$ [8,11,14]. The crossings of the two bands have alignment gains of $5.5\hbar$ at frequencies of 0.30 MeV and probably the same origin as those of the $(-,0)_2$ and $(-,1)_3$ bands.

In the energy correlation matrices a ridge-structure formed like a trumpet is observed which indicates a decrease of the dynamic moment of inertia [2]. In fig. 2 the aligned spin vs. the rotational frequency is shown. One can notice a decrease of the alignment at higher frequencies which reflects a decrease of the dynamic moment of inertia. This behaviour of the rotational bands at high spins can explain the observation of the "trumpet" ridge. It is probably due to a change from a prolate to a strongly triaxial deformation as predicted by Wyss et al. [6].

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Discrete Spectroscopy in ^{190}Au at High-Spins*

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High-spin states in ^{190}Au were populated through the reaction $^{176}\text{Yb}(^{19}\text{F},5n)^{190}\text{Au}$ at a beam energy of 107 MeV. The beam was provided by the VICKSI accelerator of the HMI, Berlin. The target consisted of a stack of two ^{176}Yb foils of 500 $\mu\text{b}/\text{cm}^2$ thickness placed ≈ 0.5 mm apart. The γ -radiation was detected by the OSIRIS multidetector array, consisting of 12 Compton-suppressed Ge detectors and 48 BGO elements used as a sum-energy and multiplicity filter.

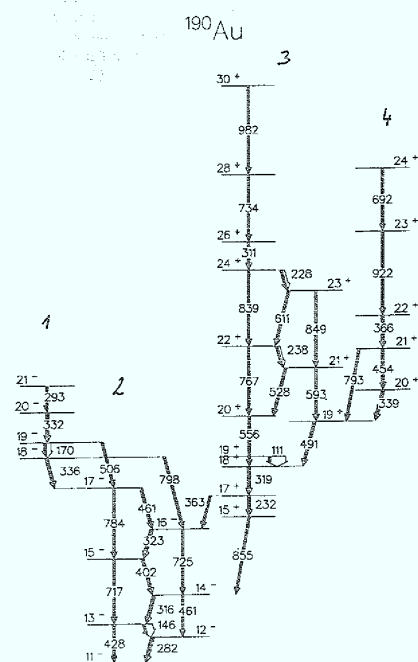


Figure 1: Partial level scheme of ^{190}Au

The partial level scheme resulting from the analysis of the coincidence matrix is shown in fig. 1. The most intense γ rays belong to the rotational band (labelled as 2) built on the 11⁻ oblate [$\nu_{1/2}^{-1} \otimes \pi h_{7/2}^{-1}$] isomeric state. A new structure (labelled as 1), similar to one observed in ^{188}Au [1], feeds into the 11⁻ band at spins 16⁻ and 17⁻. The energy levels grouped into band 3 of fig. 1 form a more irregular structure based on the 15⁺ [$\nu_{1/2}^{-1} \otimes \pi h_{7/2}^{-1} \otimes \nu_{1/2}^{-1} \otimes \nu_{1/2}^{-1} = \nu p_{3/2}$ or $\nu_{5/2}^{-1}$] four particle-hole state. The last previously known transition in this band has 556 keV (from ref. [2]). The new $\Delta I = 2$ sequence 593 and 849 keV is connected with the sequence 536, 767 and 839 keV through the M1 transitions 528, 238, 611 and 228 keV. A weak band labelled as 4 in fig. 1 feeds to the 19⁺ state of band 3. Many γ -rays have not been placed in the level scheme until now indicating the existence of different structures deexciting through complicated paths. Spin-parity assignments result from the analysis of the DCO ratios of the γ -transitions and the level systematics for the neighbouring nuclei.

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* Supported by DFG under contract number 436 BUL 113/75/0 (S)

Introduction

A new concept for the digital processing of signals from high-resolution semiconductor detectors has been introduced previously [1]. It involves the reconstruction of the original charge distribution of the detector signal from the preamplifier output pulses with the help of a newly developed digital Moving Window Deconvolution (MWD) algorithm. This concept ensures that a true ballistic measurement of the event charge can be performed and ballistic deficit effects are eliminated. A hardware implementation has been used to study the basic constraints of its technical realization [2]. A prototype of a Pulse Processing Analog to Digital Converter (PPADC) employing a software implementation of the algorithm on a Motorola 96002 Digital Signal Processor (DSP) has been developed in collaboration with industry in order to explore the benefits of the new concept for applications in high-resolution spectroscopy [3]. It was demonstrated that with respect to throughput and ballistic deficit- and noise-suppression the PPADC is able to approach the theoretical optimum in signal processing performance.

However, for multi-detector arrays like EUROBALL not only the performance of the signal processing channels but also the costs, power consumption and size are important design criteria. The EUROBALL electronics will use the VXI standard (VME-Bus Extension for Instrumentation) since it allows to integrate analogue and digital electronics on the same hardware platform. Eight complete high-resolution signal processing channels including local trigger and readout should be integrated on one D-size VXI-card and the power consumption should not exceed 75 Watts per card. In order to satisfy these requirements an attempt was made to reduce the size and power consumption of the PPADC drastically with respect to the

prototype design. It was achieved by integrating a major part of the digital signal processing algorithm into silicon.

The ASIC functional description

An Application Specific Digital Signal Processor (ASDSP) has been designed to perform the basic MWD algorithm. The circuit contains three identical deconvolution sections, two sections for dc-stability, a digital delay-line section, an input section for base-line compensation and an output stage with additional oversampling (decimation) functions (fig. 1).

Each deconvolution section consists of a 16x16 bit multiplier, 32-bit accumulator, 32-bit 4 to 1 multiplexor for rescaling the accumulated result and an output 24-bit adder. Each section receives the current data from its data input register, multiplies it by a locally stored coefficient, accumulates the product, rescales the accumulated result by a preset factor and adds it to the current data. The output result is sent to the next section via a pipeline register.

The moving window function is provided by the digital delay-line differentiating section, which is placed upstream of the deconvolution stages. The section contains a 512 x 16 bit FIFO which serves as an adjustable digital delay, a 9-bit counter to preset the delay from 1 to 512 clock cycles and a 16-bit subtractor. The section subtracts the delayed from the current data and the result is fed to the first deconvolution stage.

The long-term dc-stability is provided by two additional identical stages placed between the deconvolution sections. Each stage contains a 32-bit accumulator, rescaling multiplexor, a register to save the accumulated and rescaled result and an output subtractor. The accumulation interval is presettable up to 2 of power 24

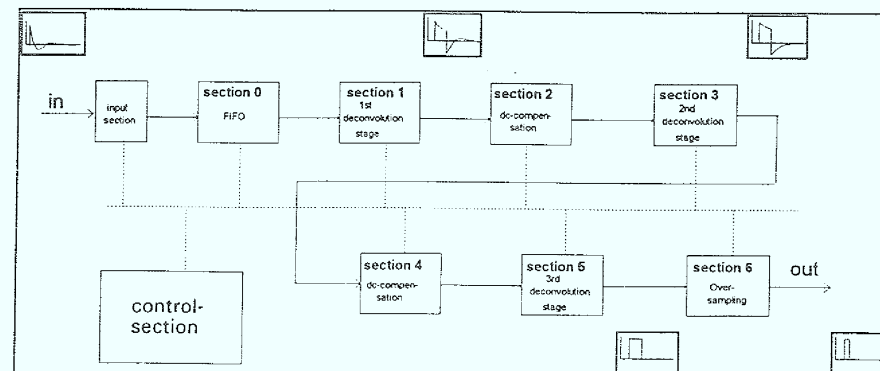


Figure 1: Block diagram of the ASDSP

clock cycles. The stage accumulates all input data during the accumulation interval, rescales the result and saves it as a reference for the next interval. The output data of the stage is calculated as a difference between the current data and that reference.

The input stage consists of a 16-bit data input register, two 16-bit base-line registers and a 16-bit subtractor. The base-line registers can be loaded with selected data from the input data stream or through the coefficient port with a reference value fed-back from the output of the circuit. The output data of the stage is calculated as a difference between the current data and the preloaded base-line value. The output section provides oversampling (decimation) for up to 256 clock cycles, which reduces the output rate by the same factor and thus allows slower devices to be placed downstream of the circuit.

The coefficient and control registers are arranged in a 32x8 bit internal memory. The access to this memory is provided by a separate coefficient port which contains 8-bit data, 5-bit address and load enable lines. Six memory locations are preserved as selective reset commands for the six sections of the circuit. Two memory locations are used as a 16-bit configuration register. Each bit of this register controls a specified path of the circuit. Each path can be open for normal data flow or closed, which can be used to define the function of each stage (normal, transparent, modified) and thus to set-up the configuration of the whole circuit.

The ASDSP accepts 25 MSPS 16-bit data flow and produces 16-bit output data at the same rate or 24-bit data at the reduced rate. It performs three multiplications, six shifts for rescaling and seven arithmetic operations in one clock cycle which is equivalent to 550 Mega-operations per second. Several ASDSP can be placed in cascade for more complex data processing. A sequential rounding is applied to minimize the error accumulation. The separate coefficient port and the individual access to each coefficient register allows an on-line update of the parameters which permits an adaptive processing.

The chip has a size of 52 sqmm and is placed in 84-pin CLCC package. The power consumption is expected to be below 1W at 25 MHz data rate. Eight power and eight ground pins are distributed around the package to reduce the switching noise.

The first samples of the ASDSP chip have been already delivered and their test is planned for the beginning of 1995.

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II. Theoretical Nuclear Physics

3. MEDIUM AND HIGH ENERGY PHYSICS

11. "Theological" Technical

A. INDIAN LITERATURE

THESE

Present microscopic hyperon-nucleon (ΛN , ΣN) interactions based on either meson [1, 2] or effective quark-gluon exchange [3] contain an exchange of a fictitious scalar meson, ' σ ', with a mass of about 600 MeV, as an important source for the required intermediate-range attraction. However, the interpretation of the σ -exchange process is quite different in the existing models. The Nijmegen group views σ -exchange as generated by genuine scalar meson exchange, to which $SU(3)$ relations are imposed (as for the $\{0^-\}$ and $\{1^-\}$ meson nonets). In contrast, in the framework of the Bonn-Jülich potentials [2, 4] the exchange of a sharp-mass σ ($m_\sigma = 550 \text{ MeV}$) has to be thought of as a parameterization of correlated 2π -exchange in the S-wave ($J^P = 0^+$) channel.

The characteristic features of the σ -contribution, e.g. the relative strengths in the several baryon-baryon (NN , ΛN , ΣN , $\Lambda\Lambda$, etc.) channels, are quite different in the existing models. The ratios of σ coupling constants, $g_{\sigma\Lambda\Lambda}/g_{\sigma NN}$ and $g_{\sigma\Sigma\Sigma}/g_{\sigma NN}$, for instance, differ by up to a factor of about 3 between the various models. Such variations provide drastic changes in interactions and calculated observables. Unfortunately, none of these YN models can be ruled out empirically since the few existing YN cross section data (with large uncertainties) can always be reasonably well reproduced by adjusting open parameters. Consequently we feel that the only chance to make progress consists in providing a reliable model for σ exchange. In order to eliminate the gross uncertainty and arbitrariness induced by this phenomenological parameterization, we [5] have evaluated explicitly the correlated 2π -exchange contribution, along the same lines as done by our group recently for the nucleon-nucleon (NN) system [6].

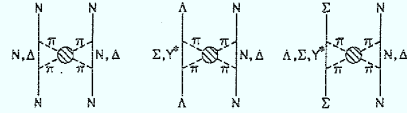


Fig. 1: Correlated 2π exchange processes in the NN and YN channels.

Fig. 1 shows our dynamical model for correlated 2π exchange. The evaluation of the diagrams in Fig. 1 has been done in two steps: First the $N\bar{N}(\Lambda\bar{\Lambda}, \Sigma\bar{\Sigma}) \rightarrow 2\pi$ amplitudes have been determined in the pseudophysical ($t \geq 4m_\pi^2$) region, with N, Δ (Λ, Σ, Y^*) exchange, cp. Fig. 2. Corresponding coupling constants in the transition interaction have been taken from the Bonn NN potential [4] and the Jülich YN models [2]. Note that this cannot be done for the parameters of the form factors (i.e. the cutoff masses) since the form factors now act in a quite different kinematic regime: here, the exchange-baryon is the essential off-shell particle. For the $NN \rightarrow \pi\pi$ case, the parameters can be fixed independently ($\Lambda_{NN\pi} = 1.9 \text{ GeV}$, $\Lambda_{N\Delta\pi} = 2.1 \text{ GeV}$) by using quasi-empirical informa-

tion obtained by Höhler et al. [8] from analytic continuation of πN (and $\pi\pi$) data. Since such information is not available for the $Y\bar{Y}$ ($Y = \Lambda, \Sigma$) channels, we make the reasonable assumption that $\Lambda_{\Lambda\Sigma\pi} = \Lambda_{\Sigma\Sigma\pi} = \Lambda_{NN\pi}$ and $\Lambda_{\Lambda Y^*\pi} = \Lambda_{\Sigma Y^*\pi} = \Lambda_{N\Delta\pi}$. The dynamical model [7] we use for the $\pi\pi \rightarrow \pi\pi$ amplitudes is likewise based on the meson exchange framework (and therefore has a well defined off-shell behavior) and involves the coupled channels $\pi\pi, K\bar{K}$. Driving terms for $\pi\pi \rightarrow \pi\pi$ consist of pole and exchange diagrams with $f_0(1400)$, $\rho(770)$ and $f_2(1274)$ intermediate states. The coupling $\pi\pi \rightarrow K\bar{K}$ is provided by $K^*(892)$ exchange.

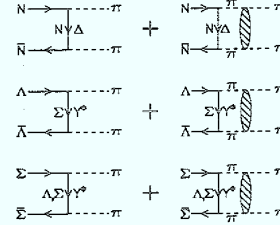


Fig. 2: Our dynamical model for the $B\bar{B} \rightarrow \pi\pi$ amplitudes.

Based on the $B\bar{B} \rightarrow \pi\pi$ amplitudes in the scalar-isoscalar channel (f_+^0) a dispersion integral together with unitarity has then been used to derive the interaction terms of Fig. 1:

$$T_{BN \rightarrow BN}^{(\sigma)}(t) \sim \int_{4m_\pi^2}^{\infty} \frac{\rho_{BN}(t')}{t' - t} dt' \quad (1)$$

In the spectral functions

$$\rho^{BN} \equiv f_+^0(B\bar{B})f_+^{0*}(N\bar{N}) - f_{+,B}^0(B\bar{B})f_{+,B}^{0*}(N\bar{N}) \quad (2)$$

the product of Born terms (see Fig. 2) has to be subtracted in order to isolate the correlated contribution.

Fig. 3 shows the resulting spectral functions ρ_{NN} , $\rho_{\Lambda N}$, $\rho_{\Sigma N}$. For the NN case, there is a sizable discrepancy compared to the quasi-empirical result [8]. However, as demonstrated in Ref.[6], this has only a minor effect on the actual NN potential contribution. We mention that, in the vector-isovector t-channel (corresponding to ' ρ '-exchange), our $N\bar{N} \rightarrow \pi\pi$ amplitudes [6] are in almost perfect agreement with the quasi-empirical result [8]. Obviously both hyperon-nucleon spectral functions are much smaller than the NN one. This can be traced to differences in the π coupling constants and isospin factors as well as in the masses of baryons exchanged.

It is instructive to parameterize the results in terms of effective σ -exchange, i.e.

$$\frac{1}{\pi} \int_{4m_\pi^2}^{\infty} \frac{\rho_{BN}(t')}{t' - t} dt' \equiv g_{\sigma BB}(t) \frac{1}{m_\sigma^2 - t} g_{\sigma NN}(t) \quad (3)$$

with $m_\sigma = 550 \text{ MeV}$. Such a parameterization does not involve any approximation since the σ -coupling con-

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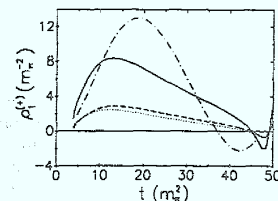


Fig. 3: Spectral functions $\rho^{\sigma N}$ as function of t in the pseudophysical region. The solid line denotes the results for the NN case derived from our dynamical model [6] whereas the dashed (dotted) line is obtained for the AN (ΣN) channel. The dash-dotted line shows the results obtained from the quasi-empirical $NN \rightarrow \pi\pi$ amplitudes [8].

stants are allowed to become t -dependent. Quite remarkably there is almost no t -dependence [5]; obviously conventional σ -exchange (with t -independent coupling constants) is indeed able to account for the actual situation provided the correct strength is used. At $t = 0$, we obtain the following ratios[5]: $g_{\sigma NN}/g_{\sigma NN} = 0.33$, $g_{\sigma \Sigma\Sigma}/g_{\sigma NN} = 0.28$. Therefore, the strength of the scalar-isoscalar component arising from correlated 2π -exchange (and by this the intermediate range attraction in the YN channels) is much reduced compared to the NN case, to an extent much larger than in any of the present YN models.

Before this scalar-isoscalar 2π -exchange piece is actually put into our YN models, the role of correlated $K\bar{K}$ exchange remains to be investigated: Especially for the strange baryons the contribution of the $B\bar{B}' \rightarrow K\bar{K}$ (Born) amplitudes (cf. Fig.2) might appear to be not negligible for the calculation of spectral functions.

Furthermore, the same evaluation of correlated 2π -exchange should be done for the vector-isovector (' ρ ') part of the $B\bar{B}'$ interaction in order to provide similar constraints as for the scalar-isoscalar component. (Especially the tensor coupling constants of the ρ differ widely in the various models). Corresponding calculations are under way.

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The role of correlated two-pion exchange in K^+N scattering

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An important ingredient of our meson exchange model for the K^+N interaction [1] is the contribution arising from σ - and ρ -exchange: σ -exchange provides the dominant part of the intermediate range attraction and ρ -exchange determines to a large part the isospin dependence of the interaction. Despite their importance, however, these pieces have been treated so far in a very rough way: In both cases a sharp mass has been used, i.e. their appreciable decay widths have not been taken into account. This simple model in terms of σ - and ρ -exchange essentially stands for the correlated 2π -exchange contribution, in the $J^P = 0^+$ (σ -) and $J^P = 1^-$ (ρ -) channel.

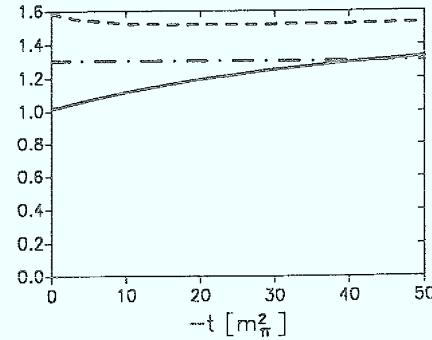


Fig. 1: Effective coupling constants for the σ -channel as a function of $-t$. The dashed dotted line shows the result used in the updated version of Ref. [1] (with $m_\sigma = 0.6$ GeV), the solid (dashed) line shows the result for correlated 2π -exchange with $m_\sigma = 0.6(0.75)$ GeV.

We now derive a microscopic model for correlated 2π (and $K\bar{K}$)-exchange between kaon and nucleon, in the scalar-isoscalar (σ) and vector-isovector (ρ) channel. For that reason, we have first constructed a model for the reaction $N\bar{N} \rightarrow K\bar{K}$ with intermediate 2π ($K\bar{K}$) states, based on a transition in terms of baryon ($N, \Delta, \Lambda, \Sigma$) exchange and a realistic $\pi\pi \rightarrow K\bar{K}$ resp. $K\bar{K} \rightarrow K\bar{K}$ amplitude [2, 3]. The contribution in the s channel is then obtained by performing a dispersion relation over the unitarity cut:

$$f_{\pm}^J(t) = \frac{1}{\pi} \int_{4m_\pi^2}^{\infty} \frac{\text{Im} f_{\pm}^J(t')}{t' - t} dt', \quad (1)$$

with the helicity (f -) amplitudes defined in [4]. (For convergence reasons, see Ref. [4], the dispersion relation in the σ -channel has to be performed for $\frac{f_{\pm}^0(t)}{p_t^2}$, with p_t being the nucleon momentum in the t channel.)

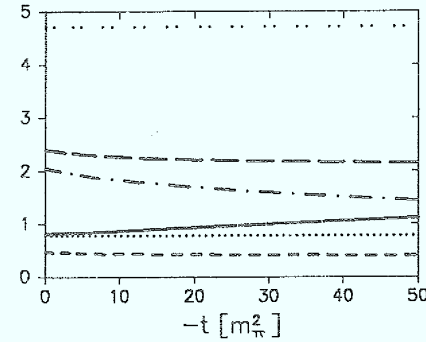


Fig. 2: Effective coupling constants for the ρ -channel as a function of $-t$. The dotted (double dotted) line shows the result for g_1 (g_2) used in the updated version of Ref. [1]. The dashed (long dashed) line shows the results for g_1 (g_2) for correlated 2π -exchange using the second method, whereas the solid (dashed dotted) line shows the results g_1 (g_2) using the first method.

Concerning ρ -exchange it was pointed out in Refs. [2, 5] that there is a considerable uncertainty in the results. The reason is that one can perform the dispersion relation directly for the helicity amplitudes or alternatively, one ([6]) can first construct combinations $\Gamma_{1,2}(t)$ corresponding to vector (Γ_1) and tensor (Γ_2) coupling amplitudes and then perform the dispersion integral. Indeed, the additional t -dependence in Γ_i apart from the t -dependence provided by the helicity amplitudes $f_{\pm}^J$ causes the results to be substantially different.

It is instructive to parametrize the result by sharp mass σ - resp. ρ -exchange with appropriate t -dependent coupling constants, which can be compared with those used in our former model. For the σ -channel the result (with $m_\sigma = 0.6$ GeV) is shown in Fig.1 together with the result of an updated model of Ref. [1] (with the same value of m_σ). Obviously our model prediction for the correlated 2π -exchange in the scalar channel has approximately the same magnitude as the σ -strength used before, which was phenomenologically adjusted to empirical K^+N data. But the sizable t -dependence of our new result demonstrates clearly that it cannot be well approximated by sharp mass σ -exchange with $m_\sigma = 0.6$ GeV. Since it grows with $(-t)$, the contribution is shorter ranged. Therefore, the effective mass exchanged should be higher. Indeed, if we use $m_\sigma = 0.75$ GeV the result also shown in Fig.1 will

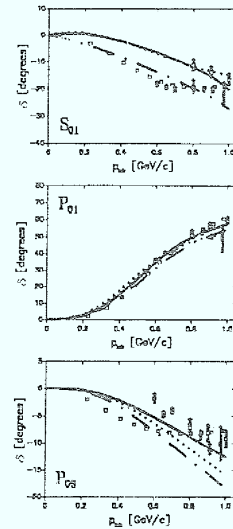


Fig. 3: S- and P- wave phase shifts for isospin 0. The dotted line shows the result for the updated version of Ref. [1]. The dashed dotted line shows the results if the dispersion integral is performed for the Γ amplitudes, the solid line, if it is performed for the f amplitudes. The empirical data is taken from Refs. [7] (empty circles), [8] (full circles) and [9] (empty boxes).

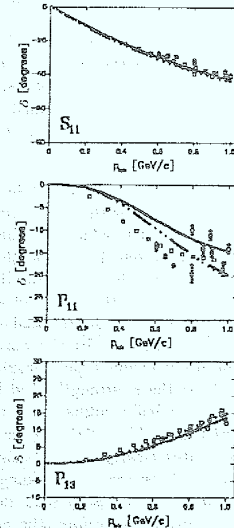


Fig. 4: The same as in Fig. 3 for isospin 1.

have almost no t dependence.

For the ρ channel the results for the vector (g_1) resp. tensor (g_2) coupling, based on $m_\rho = 769$ MeV are shown in Fig. 2, again together with the values used in the updated model of Ref. [1]. We can see that there is a remarkable difference between the two alternatives for doing the dispersion relation in the ρ -channel. The second choice (dispersion integral over the Γ amplitudes) provides larger tensor (g_2) but smaller vector coupling (g_1). Consequently the ratio g_2/g_1 characterizing the tensor to vector coupling ratio of the (effective) ρ to the nucleon, $g_{NN\rho}^T/g_{NN\rho}^V$ is much larger (~ 5) for the second method than for the first (~ 2.5). In addition, the results based on the second choice are almost t -independent; therefore, they can be well identified with an exchange of an (effective) ρ meson with the empirical mass. On the other hand, the first method yields a result with a non negligible t -dependence, with opposite behaviour for g_1 and g_2 , which, given the conventional Lagrangians, a common ρ mass cannot be assigned to. The results differ strongly from the values used before in our phenomenological ρ -exchange amplitude, which were determined from SU(3) relations. Only the ad hoc introduction of rather soft formfactors into the latter amplitude (which is required to fit the empirical K⁺N data) brings the result in rough agreement with those derived from our microscopic model.

We now want to compare our correlated 2π -exchange models with empirical data. Therefore, we have to perform an off-shell extrapolation. In addition we take into account short range terms (generated partly by conventional ω -exchange) and box diagrams involving $\pi\rho$ - and $\rho\rho$ -exchange developed before. After adjusting free vertex parameters we achieve a satisfactory description of the empirical situation, of the same overall quality as obtained before using phenomenological sharp mass σ - and ρ -exchange. As an example we show in Figs. 3 and 4 the results for empirical S- and P-wave phase shifts.

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Experimental results from LEAR for $\bar{N}N$ annihilation into two mesons (for references see [1, 2]) provide clear evidence for the dynamical suppression of transitions which are in principle allowed by the conservation of fundamental quantum numbers. Such dynamical selection rules yield valuable information about the transition process and should impose in principle severe restrictions on the theoretical description of the annihilation mechanism, e.g. in terms of conventional baryon exchange or explicit quark-gluon exchange. Unfortunately, things are very much obscured by the presence of initial ($\bar{N}N$) and final (meson-meson) state interactions. Although it was often argued that the consideration of relative ratios should minimize these effects it turned out (see e.g. [3, 4]) that there is a strong sensitivity to whether and even which kind of initial state interaction is included. The latter does not drop out even if ratios from the same partial wave are considered. We will demonstrate that also, at least in some specific channels, meson-meson correlation effects have a considerable influence on the explanation of the experimental data.

The most prominent example for such a dynamical selection rule is the so-called “ $\rho\pi$ ” puzzle: Already in the sixties, $\rho\pi$ branching ratios in liquid hydrogen have been determined [5], being essentially the same for all charge combinations. The $\rho^0\pi^0$ combination, being a pure isospin-zero state, can be produced only from the ($^3S_1, I=0$) protonium state whereas the charged combinations can also be generated from the ($^1S_0, I=1$) state. Since all S -states in protonium should be populated with about the same probability the annihilation from the 1S_0 state in the $\rho\pi$ channel is obviously strongly suppressed. In P -waves, however, such a suppression of annihilation from $I=1$ states cannot be seen in the recent results of the ASTERIX experiment [6].

In Table 1, present empirical information about these ratios is compared with theoretical results, which we obtained recently [3, 4] using a model for the transition process based on baryon (N, Δ) exchange. An initial state interaction is either completely neglected (Born) or included (A(BOX) [3], D [4]). Both $\bar{N}N$ interactions use as elastic

Table 1: Ratios of branching ratios “at rest” for the annihilation $\bar{p}p \rightarrow \rho^\pm \pi^\mp$. The experimental values are obtained from Ref. [6]. The theoretical values (as relative cross sections at $p_{lab} = 100 MeV/c$) are based on a baryon-exchange transition model, without final state interaction. The initial state interaction is either neglected, too (Born) or taken from model A(BOX) [3] and D [4].

	Born	A(BOX)	D	Exp.
$\bar{p}p(^1S_0, I=1) :$	1	1	1	1
$\bar{p}p(^3S_1, I=0) :$	3.45	2.29	4.67	35 ± 16
$\bar{p}p(^1P_1, I=0) :$	1	1	1	1
$\bar{p}p(^3P_{1,2}, I=1) :$	5.28	9.10	2.95	0.84 ± 0.22

part the G -parity transformed (full) Bonn NN potential [7]. In model A(BOX), annihilation has been accounted for by a simple phenomenological optical potential. In model D [4], the annihilation is split into two parts: processes with two intermediate mesons (all possible combinations of $\pi, \eta, \rho, \omega, a_0, f_0, a_1, f_1, a_2, f_2, K, K^*$) are described microscopically in terms of baryon-exchange processes. Their strength has been adjusted in a consistency procedure to the empirical information about the annihilation. The remaining part (three-meson channels etc.) is taken into account by a phenomenological optical potential of similar form as used in model A(BOX) [3].

The results of Table 1 indeed confirm that there is a considerable sensitivity to which kind of initial state interaction is used. Still, for the S -state ratio, all theoretical results quoted there are far off the experimental values. There are numerous calculations based on alternative transition models in the literature. For example, the model of Maruyama et al. [8] using the A2 mechanism and the 3P_0 vertex, achieves a value of 1:10-18 for the $^1S_0/^3S_1$ ratio. Thus it appears that this quark model might be superior in comparison to conventional baryon-exchange, at least in this sector. However, this conclusion is definitely premature, for the following reason: So far, correlation effects in the outgoing meson-meson channel have been neglected completely; their inclusion in the description of the transition process might well lead to different theoretical values in Table 1.

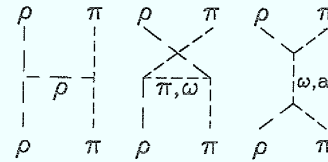


Fig. 1: Driving terms used to construct our $\rho\pi \rightarrow \rho\pi$ amplitude.

Recently we have constructed a dynamical model for the $\rho\pi \rightarrow \rho\pi$ amplitude [9] based on both s - and t -channel meson exchange diagrams, see Fig. 1, acting as driving terms in a scattering equation. Parameters (coupling constants and cutoff masses in formfactors) have been partly taken from other investigations, partly adjusted to empirical information about $a_1, \omega \rightarrow \rho\pi$ decay.

First we include the $\rho\pi$ interaction (with all parameters prefixed) in our calculation in a DWBA-type approach in the final state only, cp. Fig. 2(a,b). In a second step, we perform a $\bar{N}N, \rho\pi$ coupled channels calculation; in this way, it is also included in the initial state, see Fig. 2(c). In order to avoid double counting the $\rho\pi$ channel has to be removed from the $\bar{N}N$ annihilation part in the coupled channels calculation. This can only be done in a well defined way for the consistent microscopic annihilation model D.

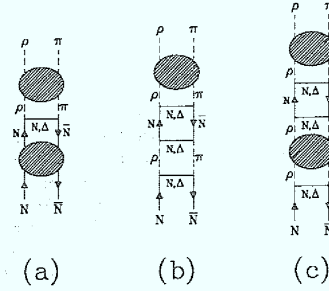


Fig. 2: $\bar{N}N \rightarrow \rho\pi$ transition model of the present paper including the $\rho\pi$ interaction (a). In the DWBA-type approach based on the $\bar{N}N$ interaction model D[4], processes (b) are included whereas (c) is not. In an alternative full coupled channels (CC) calculation, the latter is also taken into account.

Table 2: Branching ratios “at rest” for annihilation into $\rho\pi$. The data are taken from Ref.[1]. The first column denotes the results (relative cross sections at $p_{lab} = 100 MeV/c$) based on the baryon-exchange transition model with model D as initial state interaction, but without final state interaction. In the second (third) column, the meson-meson interaction is included in a DWBA calculation (as final state interaction) or in a full $\bar{N}N, \rho\pi$ coupled-channel (CC) calculation.

$\bar{p}p \rightarrow$	no FSI	DWBA	CC	EXP.
$\rho^+ \pi^- + \rho^- \pi^+$	2.32	2.50	2.73	3.4 ± 0.2
$\rho^0 \pi^0$	0.85	1.08	1.19	1.4 ± 0.1

We start by demonstrating the effect of the $\rho\pi$ interaction on the branching ratios at rest for the $\rho^+ \pi^- + \rho^- \pi^+$ and $\rho^0 \pi^0$ annihilation channels (Table 2). The consideration of the $\rho\pi$ interaction increases both branching ratios. Its inclusion in both the initial and final state (CC) leads to a result in good agreement with experiment[1].

Let us now look again at the relative branching ratios from specific $\bar{N}N$ partial wave states (cp. Table 3). The inclusion of the meson-meson interaction changes the $^1S_0/^3S_1$ ratio (corresponding to π/ω quantum numbers respectively) drastically: The change is so large that, in a DWBA calculation, the result nearly coincides with the empirical value; the full coupled channel calculation (which includes the $\rho\pi$ interaction also in the $\bar{N}N$ annihilation potential, cp. Fig. 2(c)) leads to a somewhat smaller ratio well within the experimental error bars. On the other hand, the modification of the P -state ratios is quite small. While the $^1P_1/^3P_1(l'=0)$ ratio is in agreement with experiment, annihilation from the $\bar{N}N(^3P)$ state into the $\rho\pi$ system with relative orbital angular momentum $l'=2$ provides the largest contribution of all P -states, whereas in the experiment annihilation from this state appears to be negligible. Whether an extension of our $\rho\pi$ interaction model or improved data can solve this problem remains to be seen.

Table 3: Ratios of branching ratios “at rest” for the annihilation $\bar{p}p \rightarrow \rho^\pm \pi^\mp$. The experimental values are obtained from Ref.[6]. The first column denotes the results based on the baryon-exchange transition model with model D as initial state interaction, but without final state interaction. In the second (third) column, the meson-meson interaction is included in a DWBA calculation (as final state interaction) or in a full $\bar{N}N, \rho\pi$ coupled-channel (CC) calculation. l' denotes the orbital angular momentum of the $\rho\pi$ system.

	no FSI	DWBA	CC	Exp.
$\bar{p}p(^1S_0, I=1):$	1	1	1	1
$\bar{p}p(^3S_1, I=0):$	4.67	32.4	23.9	35 ± 16
$\bar{p}p(^1P_1, I=0):$	1	1	1	1
$\bar{p}p(^3P_{1,2}, I=1):$	2.95	3.02	2.95	0.84 ± 0.22
$\bar{p}p(^1P_1, I=0):$	1	1	1	1
$\bar{p}p(^3P_1)(l'=0):$	0.54	0.68	0.64	0.8 ± 0.2
$\bar{p}p(^1P_1, I=0):$	1	1	1	1
$\bar{p}p(^3P_{1,2})(l'=2):$	2.41	2.34	2.31	0.04 ± 0.015

In any case and in more general terms, meson-meson correlation effects have to be included and treated in a consistent way before one can seriously address the question about the relevant transition mechanism in $\bar{N}N$ annihilation.

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The work presented here is based on our recently developed meson-exchange model for pion-nucleon interaction [1]. This model accounts for the scattering data in the elastic region as well as low-energy parameters and the $\pi N \Sigma$ term.

A microscopic model of πN scattering is definitely needed for investigating the spectrum of the nucleon excited states. Here the essential question is whether the experimentally observed baryon resonances really contain a genuine three-quark 'excited state' contribution or whether they can be understood purely in terms of the underlying meson-baryon dynamics. In this report, we address the Roper resonance, i.e. the first excited state with the quantum numbers of the nucleon.

The P_{11} partial wave is one of the most interesting partial waves in πN scattering. It is a challenge for models of πN interaction to describe quantitatively the repulsion at low energies followed by the attraction leading to the Roper resonance at higher energies. Within the model presented here, the rise of the P_{11} phase shift in the elastic region of the interaction can be reproduced without any contribution from a genuine N^* particle. In Fig. 1 (a) we have plotted the

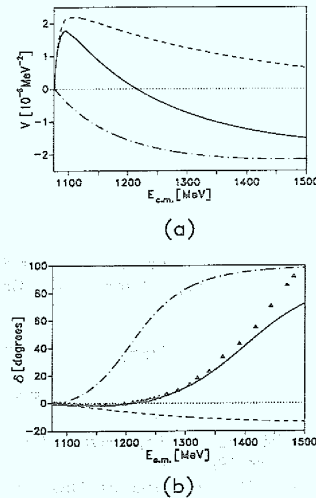


Figure 1: (a) On-shell potentials in the P_{11} partial wave as function of the c.m.s. energy. The dashed line arises from the direct nucleon pole; the dash-dotted line belongs to the non-pole part of the potential; the solid line gives the total potential. (b) Resulting phase shifts after iteration of the corresponding contributions to the potential. Empirical data points are taken from Ref. [2].

on-shell potential (i.e. the Born term) in this partial wave split into its pole part, which is generated by the direct nucleon pole, and the non-pole part. This non-pole part of the interaction is dominated by correlated 2π exchange in the ρ channel. Whereas the pole part of the potential is repulsive the non-pole part is attractive over the whole energy range investigated here. This attraction is so strong that the iteration of the non-pole part of the potential even leads to resonant behavior in this partial wave at relatively low energies ($\simeq 1.3 \text{ GeV}$, cf. Fig. 1 (b)). If this attraction is counterbalanced by the direct nucleon pole contribution the data (in the elastic region) can be described quantitatively.

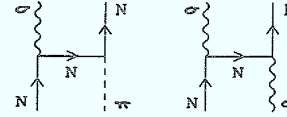


Figure 2: Additional contributions to our interaction model.

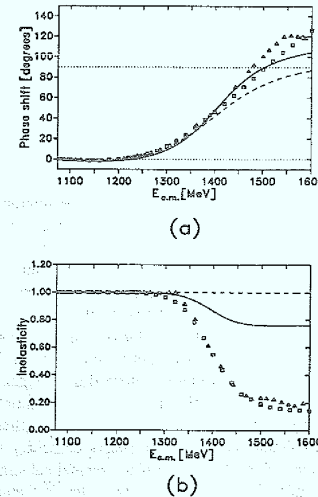


Figure 3: The $\pi N - \sigma N$ coupled channel calculation (solid lines) compared to the one-channel calculation (dashed lines). Squares (triangles) denote the empirical information from Ref. [5] (Ref. [2]).

For energies above 1.3 GeV, pion-nucleon interaction gets inelastic. According to the 'Review of Particle Properties' [3], the branching ratio of the Roper resonance into the $\pi\pi N$ channel is 30-40%. Here

the dominant contributions are the $\pi\Delta$ channel (20–30%) and the σN channel (5–15%), where σ denotes an isoscalar $\pi\pi$ S -wave state.

As a first step, we have investigated the coupled channel system $\pi N - \sigma N$, where in our model the additional contributions shown in Fig. 2 are taken into account.

The (effective) σNN coupling is taken from Ref. [4]. Self-energy contributions of the σ are implemented in our calculation. As one can see from Fig. 3, inclusion of these additional processes leads to a resonant behavior in the P_{11} partial wave. However, not surprisingly, the model accounts only for part of the inelasticity. Here the additional effect of the reaction channel $\pi\Delta$ remains to be investigated.

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With increasing experimental information about the different members of the meson spectrum it becomes more and more important to develop a consistent understanding of the observed mesons from a theoretical point of view. For the low lying pseudoscalar, vector and tensor mesons this has been done quite successfully within the framework of the simple quark model assuming the mesons to be quark-antiquark ($q\bar{q}$) states grouped together into nonets. For the scalar mesons, however, several questions still remain to be answered, most of them being related to the nature of the experimentally observed mesons $f_0(975)$ and $a_0(980)$.

In a long standing controversial discussion the $f_0(975)$ has been described as, for example, a conventional $q\bar{q}$ meson [1], a $K\bar{K}$ molecule [2] or a multiquark state [3]. In order to discriminate between different models Morgan and Pennington [4, 5] recently investigated new data on the decay $J/\psi \rightarrow \phi\pi\pi/\phi K\bar{K}$ [6, 7] concluding that a conventional Breit-Wigner (i.e. $q\bar{q}$) structure is most probable for the $f_0(975)$. They claim that the $\pi\pi/K\bar{K}$ amplitude is characterized by two poles near the $K\bar{K}$ threshold on the $[bt]$ (second) and $[bb]$ (third) sheet (see Ref. [8] for sheet structure and notation). In the Jülich model of the $\pi\pi$ interaction [9] the $f_0(975)$ appears to be a pure $K\bar{K}$ bound state generated dynamically by vector-meson exchange. Such a state has only one nearby pole on the $[bt]$ sheet, a structure seemingly disfavored by the results of Ref. [4]. We have therefore fitted this data with the $\pi\pi/K\bar{K}$ amplitude resulting from our meson-exchange model and obtained a description comparable in quality with those of the alternative models discussed above. From this result, which is shown in Fig. 1, a $K\bar{K}$ bound state structure for the f_0 is definitely not disfavored.

What turns out to be quite different in our analysis is the rather large mass (interpreted as the pole position) of the $f_0(975)$ ($m_{f_0} = 1015$ MeV, $\Gamma_{f_0} = 30$ MeV) which is a consequence of its $K\bar{K}$ bound state nature in combination with a rather strong coupling between $\pi\pi$ and $K\bar{K}$ channels. While there has been considerable discussion in the literature concerning the width of the $f_0(975)$ [5, 10, 11], there seems to be a consensus that the mass lies between 0.97 and 1.0 GeV. Interestingly, our value for the mass is somewhat larger than this due to the strong influence of the $K\bar{K}$ threshold.

Closely related to the $f_0(975)$ is the $a_0(980)$ which is observed in the $\pi\eta$ channel. As with the $f_0(975)$, the nature of the $a_0(980)$ is still under discussion, especially since the available experimental information about $\pi\eta$ scattering is rather poor compared to $\pi\pi$ scattering [12]. In Ref. [2] the $f_0(975)$ and $a_0(980)$ emerge from a non-relativistic quark model calculation to be degenerate $K\bar{K}$ bound states. However,

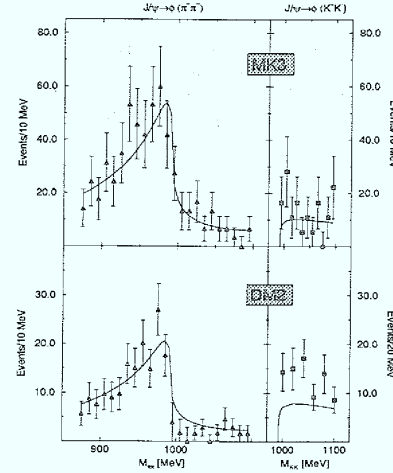


Fig. 1. Fit to $J/\psi \rightarrow \phi\pi\pi/\phi K\bar{K}$ data. The upper panel shows the fit to data of Ref. [7] and the lower panel to Ref. [6].

while the Jülich model of the $\pi\pi$ interaction leads to a $K\bar{K}$ bound state structure for the $f_0(975)$ the situation turns out to be rather different for the $a_0(980)$. The $K\bar{K}$ interaction generated by vector-meson exchange, which for isospin $I = 0$ is strong enough to generate a bound state, is much weaker for $I = 1$ making a degeneracy of $a_0(980)$ and $f_0(975)$ impossible.

It is therefore one aim of our work to address the question of whether the close relation between $a_0(980)$ and $f_0(975)$ expressed by their similar properties (masses, proximity to $K\bar{K}$ threshold, etc.) can be understood in the framework of our meson-exchange model. By extending the $\pi\pi$ model consistently to the $\pi\eta$ system we find that this is indeed possible. We obtain a model where both the $a_0(980)$ and $f_0(975)$ originate from the coupling to the $K\bar{K}$ channel. However, the underlying structure turns out to be quite different. In contrast to the $f_0(975)$, the $a_0(980)$ would not appear to be a $K\bar{K}$ bound state but a dynamically generated threshold effect with a relatively broad pole on the $[bt]$ sheet. We find poles responsible for the observed $a_0(980)$ and $f_0(975)$ at complex energies ($\text{Re}E, \text{Im}E$) = (991, ± 101) MeV and (1015, ± 15) MeV respectively. Despite such different pole positions, the observed position and width of the structure in the cross sections turn out to be rather similar. This is directly

attributable to the proximity of the $K\bar{K}$ threshold at 991 MeV. Fig. 2 shows our result for the $\pi\eta$ cross section leading to $m_{a_0} \simeq 990$ MeV and $\Gamma_{a_0} \simeq 100$ MeV, in reasonable agreement with experimental results [12].

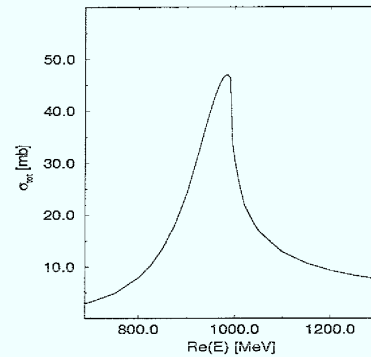


Fig. 1. The $\pi\eta$ cross section for our full model.

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Experimental results on antinucleon-nucleon ($\bar{N}N$) annihilation reactions of the type $\bar{N}N \rightarrow \pi X$ from the ASTERIX collaboration [1] have provided evidence for a new state AX , with a mass $M = 1565 \pm 10 \text{ MeV}$, width $\Gamma = 170 \pm 20 \text{ MeV}$ and quantum numbers $J^{PC}(I^G) = 2^{++}(0^+)$. Based on a detailed analysis including not only mass and width but also production rates and decay branching ratios, Dover, Gutsche and Faessler [2] have made the case that this new state should be interpreted as a quasinuclear $\bar{N}N$ bound state. However, as stressed by the authors themselves, their arguments rely heavily on the existence of two further quasinuclear states with $J^{PC}(I^G) = 1^{--}(0^-)$ at about 1250 MeV and $0^{++}(0^+)$ at 1100 MeV ; in fact, they propose key experiments to provide conclusive evidence above the sparse hints in the present data.

The purpose of this work is twofold: First, we provide further support for the interpretation of the AX as a quasinuclear bound state. Second, we show that the arguments of Ref. [2] are too restrictive; namely, it is demonstrated that this interpretation does not unambiguously require the existence of the additional states mentioned above.

As discussed in Refs. [2, 3] the above mentioned states 0^{++} , 1^{--} and 2^{++} should be considered as members of an $I = 0$ rotational band, which is generated by the coherent action of attractive tensor forces in the $\bar{N}N$ interaction. This spectrum is therefore realized if the tensor forces are the by far dominant components in the elastic $\bar{N}N$ interaction compared to other possible contributions (of central, spin-spin and spin-orbit type). Indeed, this seems to be the case for the Dover-Richard potential, c.p. Refs. [2, 3], based on the Paris potential [4]. In this model, the inner part of the elastic interaction is cut off at a distance r_0 (chosen in Ref. [2] to be about 0.8 fm in order to bring the predicted $2^{++}(0^+)$ state at the experimental AX position) and set to a constant value.

However, (as we will demonstrate below) the characteristics of the spectrum is completely changed if the Bonn NN potential is used as starting point and the inner region of the elastic interaction is treated differently.

The following calculations are based on an elastic $\bar{N}N$ interaction (V_{el}) taken to be the G -parity transform of the full Bonn NN potential [5]. Note that V_{el} is completely prescribed in terms of meson exchanges since we do not introduce an additional cutoff on top of the form factors fixed already in the NN system.

The $\bar{N}N$ bound state spectrum (for states with $M > 0.5 \text{ GeV}$) calculated from V_{el} by the solution of a bound-state equation is shown in figure 1 in the conventional notation $(2I+1)(2S+1)L_J$; numerical values are given in Table 1. Quite interestingly, without any adjustments of parameters the model predicts a bound state in the $^{13}PF_2$ state at 1564 MeV , i.e. just at the position of the empirical $AX(1565)$ with the same quantum numbers, and thus supports the interpretation of the AX as quasinuclear state. The empirical width of this state, however, can not be described in the present study, because we do not consider

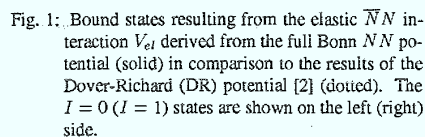
Table 1: $\bar{N}N$ Bound states predicted by the elastic $\bar{N}N$ interaction V_{el} derived from the full Bonn NN potential. In the upper half, the isospin $I = 0$ states are given, in the lower part the ones for isospin $I = 1$.

$\bar{p}p$ -state	$M [\text{MeV}/c^2]$
$^{11}S_0$	1380
$^{13}SD_1$	1847
$^{13}SD_1$	1754
$^{13}SD_1$	829
$^{11}P_1$	1461
$^{13}P_0$	1854
$^{13}P_0$	1026
$^{13}PF_2$	1564
$^{31}S_0$	1844
$^{31}S_0$	1002
$^{33}SD_1$	1875
$^{33}SD_1$	1307
$^{31}P_1$	1683
$^{33}P_1$	1518
$^{33}P_0$	1664
$^{33}PF_2$	1825

the annihilation part of the $\bar{N}N$ interaction.

On the other hand, our model does obviously not result in the rotational band structure $^3P_0(0^{++})$, $^3SD_1(1^{--})$, $^3PF_2(2^{++})$ for the $I = 0$ states, suggested in Ref. [2]. In our case, the 3SD_1 state lies lowest and additional spin singlet states ($^1S_0, ^1P_1$) appear. Therefore, in our model the $I = 0$ tensor forces are not so dominant as anticipated by Dover et al. We have quite naturally a central part due to 2π - and ω -exchange acting strongly in the inner region of the $\bar{N}N$ interaction. Only if we arbitrarily cut off the inner part of our interaction and furthermore use only π and ρ exchange (which generate predominantly tensor forces), we obtain a similar rotational band structure as suggested in Ref. [2].

Thus, our main point is the following: In case future experimental results cannot confirm the candidate members of the $I = 0$ rotational band at the positions discussed in Ref. [2] to be $\bar{N}N$ bound states, the interpretation of the $AX(1565)$ as a quasinuclear bound state is not necessarily ruled out. Anyway, apart from the elastic $\bar{N}N$ interaction V_{el} , annihilation is important in the $\bar{N}N$ system: processes like $\bar{N}N \rightarrow M_1 M_2$ or $\bar{N}N \rightarrow M_1 M_2 M_3$ contribute to the $\bar{N}N$ interaction. Although their main effect is to provide hadronic widths to the $\bar{N}N$ states, second iterations of these processes will lead to an additional real part of the $\bar{N}N$ interaction which then might cause state dependent shifts of masses of the bound states, depending on the particular annihilation model.



References:

Due to increasing computational efforts the description of the nucleon-nucleon interaction within the Skyrme model has been improved tremendously over the past years [1]. Multi-soliton configurations provide a semiclassical picture for interacting nucleons and leads e.g. to a quite successful description of the long range part of the N-N potential [2]. This will increase our knowledge of the structure of the non-perturbative regime of QCD since the Skyrme model is to be expected to be the approximate large N_c expansion model of the strong interaction at low energies [3].

Most of the recent results were based on static numerical calculations. However, only a time dependent analysis of the Skyrme model can give us informations about the hidden dynamics of two solitons. Here the tools are again numerical lattice methods which are able to tackle the highly nonlinear process of interaction between solitons.

To perform these lattice calculations, the Skyrme Lagrangian is usually written in the form

$$\mathcal{L} = \frac{f_\pi^2}{2}(\partial_\mu \phi_\alpha)^2 - \frac{1}{4e_s^2}(\partial_\mu \phi_\alpha)^2(\partial_\nu \phi_\beta)^2 + \frac{1}{4e_s^2}(\partial_\mu \phi_\alpha \partial_\mu \phi_\beta)^2 + m_\pi^2 f_\pi^2 (\phi_4 - 1) \quad (1)$$

where ϕ_α ($\alpha = 1, 4$) is the nonlinear pionfield constraint to the chiral circle ($\phi_\alpha^2 = 1$). The parameters of the model are the pion decay constant f_π , the Skyrme parameter e_s and the pion mass m_π . The Euler-Lagrange equations of the fields are calculated to

$$\partial_t \phi_\alpha = M_{\alpha\beta}^{-1} \pi_\beta \quad (2)$$

$$\partial_t \pi_\alpha = \partial_i (C_{\alpha\beta} \partial_i \phi_\beta) \quad (3)$$

where π_α is the canonical field momentum to ϕ_α . $M_{\alpha\beta}$ and $C_{\alpha\beta}$ are functions of the fields and their derivatives [1]. These nonlinear, coupled partial differential equations can be solved on a space-time lattice via a finite difference method. The necessary initial and boundary conditions for these equations are defined by the physical process one wants to analyse.

The description of the real particles like the nucleons follows from a semiclassical quantization method [4]. A first step in this procedure is to obtain all the different configurations for the zero and nearly zero modes of a chosen soliton sector. These configurations serve later on as background fields in the quantization process. For the two soliton sector of our model, energy and symmetry considerations lead to 12 different zero and nearly modes of the solitonic field [1]. Neglecting here the zero modes (they are related to symmetric field configurations and have to be discussed separately), one ends up with 5 relativ

degrees of freedom. They are the 3 spatial coordinates $R = R(r, \phi, \theta)$ for the separation axis and the 2 isospin coordinates $C = C(\alpha, \beta)$ for rotations of the solitons relative to each other in isospin space. A transformation of these coordinates into time dependent quantities results in the classical collective coordinate Hamiltonian

$$H_c = 2M_s + \frac{1}{2} \dot{Q}_k M_{kl} \dot{Q}_l + V(Q) \quad (4)$$

which can be quantized by using the canonical commutator relations between the collective coordinates and their momenta. Here M_s is the classical mass of one soliton, $Q_k = (R, C)$ are the collective coordinates, M_{kl} is the metric or mass tensor of the system and $V(Q)$ is the classical potential energy of two interacting solitons. This is the model Hamiltonian one has to study to get informations about the nucleons and their interactions.

All the relevant terms in the Hamiltonian H_c are functions of the soliton fields $\phi_\alpha = \phi_\alpha(Q)$. The main work is to determine the corresponding fields to all possible configurations Q and to calculate simultaneously the masses and energie in equation (4). The separation R we fix on the lattice as the mean baryonic rms radius while for the relative isospin orientation C of the system we use a special symmetry relation between the two solitons [1].

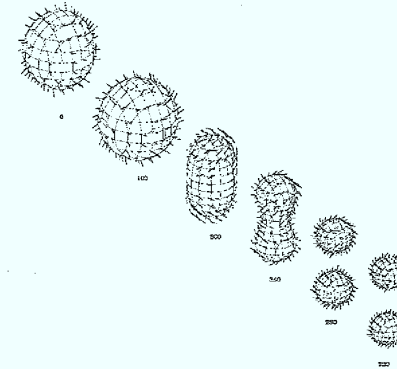


Fig. 1. Contourplot of the baryon number density for different time steps. The arrows are the isospin vectors at the chosen baryon number density.

To obtain a complete set of configurations [5] we start from the excited baryon number $B=2$ Hedgehog solution where we had calculated all its unstable modes previously [6]. Evolving this system in time we see a splitting of the Hedgehog to a lower potential energy configuration (Fig. 1). The $B=2$ soliton is moving in

the internal collective coordinate space and starts to split and spin in isospin space.

Unfortunately, due to the large initial energy, our numerical method is not capable to handle the high velocity rotations of the isospin which occur at some spatial points during the fission of the Hedgehog. After reaching a maximal kinetic energy the total energy on the lattice starts to increase in a few time steps dramatically and the procedure becomes numerical instable. We have tried a set of numerical methods but the evolution of the soliton configuration still broke down at the same critical point. This might be related to the problems considered in [7]. The equation of motion for the Skyrme model could run from a stable hyperbolic into a unstable non-hyperbolic partial differential equation for large kinetic energies. We will investigate this problem further. At this point our suggestion is to use an additional viscosity term in the equation of motion. This term will damp the small amplitude fluctuations of the field configurations which are responsible for these instabilities. Since this results also in a decrease of the total energy one has to turn off the damping terms after the critical point.

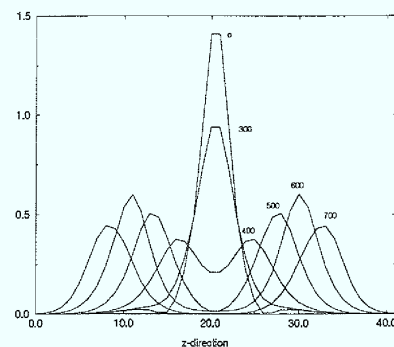


Fig. 2. Cut in the baryon number density along the separation axis for different time steps t

Introducing viscosity allows to determine of the collective coordinates as functions in time. Additionally, one can consider the baryon number density as in Fig. 2. It seems that not only the mentioned collective coordinates are important for the evolution of the solitonic field. From timestep 400 to 600 in the figure above the individual solitons experience a strong compression mode and will therefore start to vibrate. This has been seen before in soliton scattering calculations [8] and is due to the fact that the shape defining potential of the soliton is rather weak. We expect even stronger excitation of this vibrational modes when the paths the solitons take are in the vicinity of the toroidal solution. There the solitons merge into a axial symmetric configuration. When returning into individual solitons they will be very excited, too [9].

Determination of the metric M_{kl} in terms of the collective coordinates R and C is therefore a difficult task since we have to introduce more degrees of freedom to describe the complete motion of the solitons on the lattice. This will increase our numerical expense, but we hope to control this problem in the future and to gain an improved picture of the nucleon-nucleon interaction in the Skyrme model.

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It is anticipated that future experiments at the Brookhaven Relativistic Heavy Ion Collider (RHIC) will create a highly excited, but nearly baryon-free central zone. At later stages of the collision this zone will contain a dense meson gas, primarily consisting of pions with temperatures of the order of the pion mass. The equation of state and the transport properties of the pion gas are therefore of great current interest [1, 2, 3, 4]. However, when calculating thermal properties of an interacting pion gas it has been noticed [5] that the commonly employed quasiparticle approximation (QPA) becomes unreliable for temperatures above $T \approx 150$ MeV, i.e. in the vicinity of T_c , the chiral restoration temperature.

Within the Matsubara formalism we derive exact, off-shell expressions for the pion selfenergy and the uncorrelated two-pion propagator. The latter enters the scattering equation describing the $\pi\pi$ interaction in the meson exchange picture [6, 7]

$$M_{\pi\pi} = V_{\pi\pi} + V_{\pi\pi} G_{\pi\pi} M_{\pi\pi}, \quad (1)$$

where $M_{\pi\pi}$ denotes the invariant scattering amplitude.

At moderate gas densities it is sufficient to treat medium modifications at the one- and two-body level rendering a Brueckner-type description. Leaving the pseudopotentials $V_{\pi\pi}$ temperature independent, which should be a good approximation for temperatures well below the masses of the exchanged mesons, the temperature dependence of $M_{\pi\pi}$ is solely induced by the thermal two-pion propagator, $G_{\pi\pi}^T$. In terms of the thermal single-pion propagator, D_π , it can be evaluated by means of a Matsubara sum [8, 9]:

$$G_{\pi\pi}^T(\Omega_\lambda, \vec{k}_1, \vec{k}_2) = T \sum_{z_\nu} D_\pi(z_\nu, \vec{k}_1) D_\pi(\Omega_\lambda - z_\nu, \vec{k}_2), \quad (2)$$

with discrete frequencies

$$z_\nu = i\omega_\nu, \quad \Omega_\lambda = i\omega_\lambda, \quad \omega_\nu = 2\pi\nu T, \\ \nu, \lambda \text{ integer}. \quad (3)$$

Using standard techniques of analytical continuation one finds two contributions to the imaginary part of $G_{\pi\pi}^T$:

$$ImG_{\pi\pi}^{T,A}(E_+, k) = - \int_0^E \frac{d\omega}{\pi} \\ \times ImD_\pi(\omega_+, k) ImD_\pi(E - \omega_-, k) \\ \times [1 + f^\pi(\omega) + f^\pi(E - \omega)] \\ ImG_{\pi\pi}^{T,B}(E_+, k) = - \int_E^\infty \frac{2d\omega}{\pi} \\ \times ImD_\pi(\omega_+, k) ImD_\pi(\omega_+ - E, k) \\ \times [f^\pi(\omega - E) - f^\pi(\omega)], \quad (4)$$

where $\omega_\pm = \omega \pm i\eta$ etc. and $f^\pi(\omega) = (\exp[\omega/T] - 1)^{-1}$. In the zero-temperature limit contribution A reproduces the imaginary part of the vacuum two-pion propagator, whereas B vanishes. The latter arises solely from thermal excitations. The real part of $G_{\pi\pi}^T$ is calculated via a dispersion integral [10]:

$$ReG_{\pi\pi}^T(E, k) = -\mathcal{P} \int_0^\infty \frac{dE'^2}{\pi} \frac{ImG_{\pi\pi}^T(E'_+, k)}{E^2 - E'^2}. \quad (5)$$

In a second step we evaluate the pion selfenergy, Σ_π , which enters the single-pion propagator

$$D_\pi(\omega_+, k) = \frac{1}{(\omega_+)^2 - m_\pi^2 - k^2 - \Sigma_\pi(\omega_+, k)}. \quad (6)$$

At the two-body level Σ_π is expressed in terms of the forward scattering amplitude as

$$\Sigma_\pi(z_\nu, \vec{k}) = - \int \frac{d^3p}{(2\pi)^3} T \sum_{z_{\nu'}} D_\pi(z_{\nu'}, \vec{p}) \\ \times M_{\pi\pi}(z_\nu + z_{\nu'}, \vec{k}, \vec{p}). \quad (7)$$

Using the spectral representations of $M_{\pi\pi}$ and D_π one arrives at [11]

$$\Sigma_\pi(\omega_+, \vec{k}) = \int \frac{d^3p}{(2\pi)^3} \frac{dE'}{\pi} \frac{d\omega'}{\pi} \\ \times ImD_\pi(\omega'_+, \vec{p}) ImM_{\pi\pi}(E'_+, \vec{k}, \vec{p}) \\ \times \frac{f^\pi(\omega') - f^\pi(E')}{\omega_+ + \omega' - E}. \quad (8)$$

For numerical evaluation of eq. (8) we transform the $M_{\pi\pi}$ -amplitude into the two-pion CMS identifying $E^2 \equiv s = (\omega \pm \omega')^2 - (\vec{k} + \vec{p})^2$:

$$M_{\pi\pi}(\omega \pm \omega', \vec{k}, \vec{p}) = M_{\pi\pi}(\sqrt{s}, q, q) \quad (9)$$

with $q = q(\omega, \omega', k, p, \cos\Theta)$ being the CMS momentum of both in- and outgoing pions. For $ImM_{\pi\pi}$ we take into account all partial waves up to $J = 2$ and all isospins. The real part of Σ_π is calculated by means of a dispersion relation including a subtraction at $\omega = 0$:

$$Re\Sigma_\pi(\omega, k) = -\mathcal{P} \int_0^\infty \frac{d\omega'^2}{\pi} \frac{Im\Sigma_\pi(\omega'_+, k)}{\omega^2 - \omega'^2} \frac{\omega^2}{\omega'^2}. \quad (10)$$

The subtraction is dictated by the fact that $Re\Sigma_\pi(\omega, k) \rightarrow 0$ for $\omega \rightarrow 0$. Eqs. (1), (4), (5), (8) and (10) form the selfconsistent set of equations which is solved by numerical iteration.

The results for the imaginary part of the $\pi\pi$ M-amplitude are displayed in Fig. 1. With increasing temperature we find in the σ -channel ($JI=00$) a considerable enhancement over the vacuum curve for

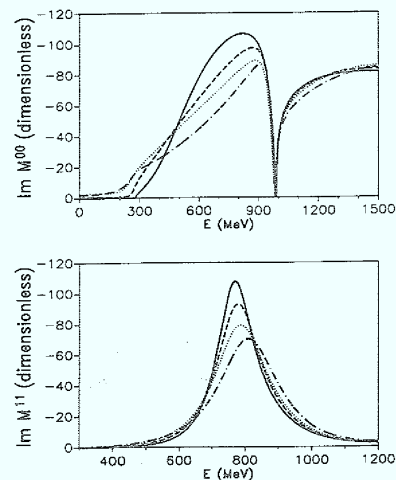


Figure 1: The imaginary part of the $\pi\pi$ M-amplitude in σ ($J=0$) and ρ ($J=1$) channel in a hot pion gas at $\mu_\pi=0$. The dashed and dotted lines are the selfconsistent results of the full off-shell calculations at $T=150, 200$ MeV, the dashed-dotted lines represent the selfconsistent QPA results at $T=200$ MeV. The full lines correspond to the vacuum amplitude.

CMS energies below $E \approx 450$ MeV. Apart from the highest temperature ($T=200$ MeV) there is, however, no significant strength below $2m_\pi$. For comparison, results within the QPA are also given (in contrast to ref. [5] the 'polestrength' $z_k = (1 - \partial \text{Re} \Sigma_\pi(\omega, k) / \partial \omega^2|_{s_k})^{-1}$ is not included since it leads to a violation of the normalization condition for the 1-pion spectral function). For $T=200$ MeV some accumulation of subthreshold strength is found which is almost entirely due to thermal excitations (the 'B'-contribution in eq. (4)). For the ρ -channel ($J=1$) both the full and the QPA results show a broadening of the resonance, increasing with temperature due to the thermal motion of the gas particles. The QPA results clearly overestimate the thermal broadening. A comment concerning the QPA calculations of refs. [5, 12] is in order. When expanding around the quasiparticle pole, in next-to-leading order a complex polestrength z_k arises. This results in a strong enhancement of $\text{Im} M_{\pi\pi}^{11}$ near the resonance peak. Such a strong enhancement is not confirmed by the full calculations presented here. In the low-energy region, the full calculations do show an enhancement over the QPA results, however.

In conclusion, we wish to stress that proper off-shell calculations beyond the QPA are necessary to obtain reliable results for the pion selfenergy and the in-medium $\pi\pi$ amplitudes of a hot pion gas at temperatures beyond 150 MeV. These quantities could

e.g. serve as an input for transport models simulating the time evolution of ultrarelativistic heavy-ion collisions [13] near freeze-out.

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Temperature Dependence of the Chiral Condensate from an Interacting Pion Gas

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The aim of relativistic heavy-ion collisions with center energies $E_{cm} \geq 20$ GeV/A is to create a quark-gluon plasma at early stages of the collision. It is predicted that the formation of the plasma is accompanied by a chiral restoration transition which, for two light flavors, is most likely second order [1]. To obtain an expression for the chiral condensate ratio $\langle \bar{q}q \rangle_T / \langle \bar{q}q \rangle_0$ in terms of pionic degrees of freedom we can use the fact that, according to the Feynman-Hellmann theorem, $\langle \bar{q}q \rangle_T$ is related to the derivative of the free-energy density (the thermodynamic potential) $\Omega(T)$ with respect to the bare quark mass m as

$$\langle \bar{q}q \rangle_T = \partial \Omega(T) / \partial m. \quad (1)$$

Denoting the difference in free energy density as $\Omega(T) = \Omega(T) - \Omega(0)$ and using the Gell-Mann Oakes Renner relation [2]

$$m_\pi^2 f_\pi^2 = -m \partial \Omega(0) / \partial m \quad (2)$$

one immediately derives that

$$\frac{\langle \bar{q}q \rangle_T}{\langle \bar{q}q \rangle_0} = 1 - \frac{1}{f_\pi^2} \frac{\partial \Omega(T)}{\partial m_\pi^2} \quad (3)$$

This expression is identical to that of ref. [3] if one ignores correction to f_π and m_π of order m .

In ref. [3] chiral perturbation theory, including three loops, has been used to evaluate $\Omega(T)$. As the authors point this is equivalent to a virial expansion in powers of T up to power T^8 . We pursue a different route by calculating $\Omega(T)$ from a non-perturbative method, well-known for phonon-phonon interactions in condensed matter physics [4, 5]. While less systematic from the point of view of a gradient expansion of the effective QCD Lagrangian it has the advantage of treating the vicinity of the phase transition more properly.

For the evaluation of the free energy $\Omega(T)$ we start from well-known expression for interacting boson systems at finite temperature [4, 5]. In terms of the pion Green's function [6]

$$D(\xi, k, T) = (\xi^2 - m_\pi^2 - k^2 - \Pi(\xi, k, T))^{-1} \quad (4)$$

the thermodynamic potential is given by

$$\begin{aligned} \Omega(T) = & \Omega'(T) - \frac{3}{2} \int \frac{d^3k}{(2\pi)^3} \int \frac{d\xi}{\pi} \\ & \text{Im} \left\{ \ln[-D^{-1}(\xi, k, T)] \right. \\ & \left. + D(\xi, k, T) \Pi(\xi, k, T) \right\} \end{aligned} \quad (5)$$

where $\Omega'(T)$ is the sum of all contributions from 'skeleton diagrams' representing the perturbation expansion of Ω . These are evaluated by using full single-particle Green's functions rather than bare ones. Furthermore

$$\delta \Omega / \delta \Pi = 0 \quad (6)$$

and

$$\delta \Omega' / \delta D = \Pi \quad (7)$$

(see refs. [4, 5, 6]). Since the physical pion mass enters the propagator D , the self energy $\Pi(\xi, k, T)$ is the difference

$$\Pi(\xi, k, T) = \tilde{\Pi}(\xi, k, T) - \tilde{\Pi}(\xi, k, 0), \quad (8)$$

whereby the infinite contributions to the mass operator $\tilde{\Pi}(\xi, k, 0)$ due to vacuum fluctuations are removed from $\Pi(\xi, k, T)$. Consequently also $\Omega(T)$ is a finite difference $\Omega(T) = \Omega(T) - \Omega(0)$ and no cut-off's are needed.

The key quantity in eq. (5) is the self energy Π . To calculate it we use the Weinberg Lagrangian [7]. To lowest order in the coupling constant $1/f_\pi$ it takes the form

$$\begin{aligned} \mathcal{L} &= \mathcal{L}^0 + \mathcal{L}' \\ \mathcal{L}^0 &= \frac{1}{2} \partial_\mu \pi \partial^\mu \pi + \frac{1}{2} \pi^2 m_\pi^2 \\ \mathcal{L}' &= \lambda (-\partial_\mu \pi \partial^\mu \pi + \frac{1}{2} \pi^2 m_\pi^2) \pi^2 \end{aligned} \quad (9)$$

where $\lambda = f_\pi^{-2}/4$ ($f_\pi = 93$ MeV). Here $\mathcal{L}^0$ is free pion Lagrangian and $\mathcal{L}'$ denotes the $\pi\pi$ -interaction. The second term in $\mathcal{L}'$, which violates chiral symmetry for physical pions, is introduced according to ref. [8]. The Weinberg Lagrangian reproduces the current algebra results for the $\pi\pi$ scattering lengths, and has been successful in describing low-energy $\pi\pi$ scattering [9]. The interaction Lagrangian $\mathcal{L}'$ is irreducible and describes a point-like pion-pion interaction. Hence, by definition, it suffers no medium modification.

Given eq. (8), the pion self energy $\tilde{\Pi}$ contains single-pion and three-pion contributions

$$\tilde{\Pi}(\xi, k, T) = \tilde{\Pi}(\xi, k, T)^{1\pi} + \tilde{\Pi}(\xi, k, T)^{3\pi}. \quad (10)$$

Because of eq. (8), the difference $\Pi(\xi, k, T)$ only includes intermediate virtual states with energies $E \leq T$. Because of resonances, the three-pion contribution should reduce to diagrams involving heavy mesons in the intermediate states. Since their masses are larger than the temperature range of interest we only include the first part of eq. (10) which is equivalent to the Hartree approximation. In this case, the quantity $\Pi(\xi, k, T)$ takes the form, see ref. [10]

$$\begin{aligned} \Pi^{1\pi}(\omega, k, T) &= -10d\lambda m_\pi^2 + 6d\lambda(\omega^2 - k^2) \\ &\quad + 6d\lambda \tilde{m}_\pi^2 \end{aligned} \quad (11)$$

and pion propagator is given by

$$D(\omega, k, T) = \gamma(\omega^2 - k^2 - \tilde{m}_\pi^2)^{-1}, \quad (12)$$

It contains a temperature dependent effective mass

$$\tilde{m}_\pi^2 = \frac{1 - 10\lambda d(T)}{1 - 12\lambda d(T)} m_\pi^2 \quad (13)$$

and a residue $\gamma = (1 - 6\lambda d(T))^{-1}$. The quantity $d(T)$ is determined by [10]

$$d(T) = \frac{1}{2\pi^2} \int_0^\infty \frac{dk k^2}{\omega_k} \frac{\chi(\omega_k)}{1 - 6\lambda d(T)} \quad (14)$$

$$\chi(\omega_k) = \left[\exp(\omega_k/T) - 1 \right]^{-1}$$

with

$$\omega_k^2 = k^2 + \tilde{m}_\pi^2; \quad (15)$$

It should be noted that, for any temperature, $d(T)$ never reaches the value $1/12\lambda$.

In the Hartree approximation, Ω' takes the form

$$\Omega' = \frac{3\lambda}{4} \int \frac{dk}{(2\pi)^3} \int \frac{dp}{(2\pi)^3} \int_{-\infty}^\infty \frac{d\xi}{\pi} \chi(\xi) \text{Im} D(\xi, k, T) \int_{-\infty}^\infty \frac{d\eta}{\pi} \chi(\eta) \text{Im} D(\eta, p, T) (-10m_\pi^2 + 6(\xi^2 - k^2) + 6(\eta^2 - p^2)). \quad (16)$$

The expressions in eq. (5), eqs. (11-15) and eq. (16) determine $\Omega(T)$ in the Hartree approximation and can be used to find the chiral condensate ratio via eq. (3).

The thermodynamic potential Ω depends on the free pion mass m_π , both, explicitly and through the self energy Π . Due to the stationarity condition (7) the m_π -dependence of Π does not influence the calculation of $\partial\Omega/\partial m_\pi^2$ and we obtain

$$\frac{\partial\Omega}{\partial m_\pi^2} = \frac{3}{2} \int \frac{dk}{(2\pi)^3} \int_{-\infty}^\infty \frac{d\xi}{\pi} \chi(\xi) \text{Im} D(\xi, k, T) - \frac{30\lambda}{4} \int \frac{dk}{(2\pi)^3} \int \frac{dp}{(2\pi)^3} \int_{-\infty}^\infty \frac{d\xi}{\pi} \chi(\xi) \text{Im} D(\xi, k, T) \int_{-\infty}^\infty \frac{d\eta}{\pi} \chi(\eta) \text{Im} D(\eta, p, T) \quad (17)$$

which can be expressed in terms of $d(T)$ as

$$\frac{\partial\Omega}{\partial m_\pi^2} = \frac{3}{2} d(T) - \frac{30\lambda}{4} d^2(T). \quad (18)$$

where $d(T)$ is the solution of the integral equation (14). The result is displayed in Fig. 1.

As expected, one observes a decrease of $\langle \bar{q}q \rangle_T / \langle \bar{q}q \rangle$ with temperature which, for the physical values of m_π MeV and λ , reaches, however, a constant asymptotic value of ~ 0.7 at high temperature. This is easily understood from eq. (18), recalling the behaviour of the quantity $d(T)$ for $T \rightarrow \infty$, as mentioned above. Our result is quite different from the

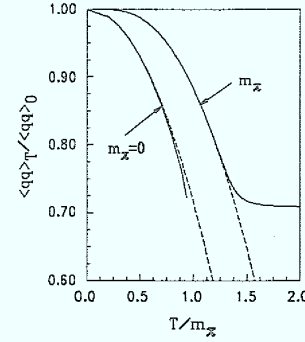


Fig. 1: Temperature dependence of the chiral condensate from an interacting pion gas. The full lines denote the results with $\lambda = f_\pi^2/4$ both for the physical pion mass and in the chiral limit. The dashed lines give the corresponding results with $\lambda/4$.

chiral perturbation theory [3]. In that work the condensate ratio monotonically drops with temperature. One should recall here that our calculations are carried out in the Hartree approximation summing a selected, but infinite, number of loops. In such a self-consistent scheme the strong pion-pion interaction first enhances $\partial\Omega/\partial m_\pi^2$ with T growth, but then it prevents a further increase due to non-linearities.

To demonstrate the non-perturbative effects we have repeated the calculation artificially increasing f_π by a factor of two, i.e. reducing the coupling constant by a factor of 4. In this case perturbation theory should be more appropriate. One finds that, for low values of T the reduction, is the same as with the full coupling, but then continues to higher T . Perhaps accidentally these results are very close to those of ref. [3].

True chiral restoration only takes place in the chiral limit $m_\pi \rightarrow 0$. We have studied this limit as well and the result is also shown in Fig. 1. In this case the temperature-dependent effective mass $\tilde{m}_\pi(T)$ remains zero and the pion propagator becomes

$$D_0 = \gamma(\omega^2 - k^2)^{-1}. \quad (19)$$

with the residue γ being of the same form as for $m_\pi \neq 0$. The solution for $d(T)$ in eq. (14) now reduces to the very simple form

$$d_0(T) = \frac{1 - \sqrt{1 - 24\lambda j(T)}}{12\lambda}, \quad (20)$$

$$j(T) = \int_0^\infty \frac{k^2 dk \chi(\omega_k^0)}{2\pi^2 \omega_k^0}, \quad \omega_k^0 = k.$$

One concludes immediately that, for temperatures where $j(T) > 1/24\lambda$ there no longer exists a real solution for $d_0(T)$, i.e. the residue γ becomes imaginary. Physically this means that, above the critical temperature, T_c , where $j(T_c) = 1/24\lambda$ no stable single-pion states exist. This we interpret as the signature of the chiral phase transition. At this point the chiral condensate remains *finite*, however, due

to non-perturbative effects. Reducing the coupling strength to $\lambda/4$ (dashed line) the condensate drops to zero again.

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**In-Medium $\pi - \pi$ Correlations:
To condense or not to condense**

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It is believed that the exchange of two correlated pions with the quantum numbers $I=J=0$ (the "sigma-meson channel") accounts for the major part of the nuclear attraction. Two-pion exchange thus constitutes a crucial ingredient for nuclear saturation. On the other hand we know from pion-nucleus scattering phenomenology that pions are strongly renormalized through Δ -hole ($\Delta-h$) and nucleon-hole ($N-h$) excitations. One may therefore ask what influence this pion renormalization has on pion-pion correlations in the medium. We have investigated this question recently in a series of papers [1, 2, 3] with rather spectacular results. It turns out that, after the inclusion of $N-h$ states [3], a collective state at low energy appears which, for slightly compressed nuclear matter, moves to zero energy indicating an instability with respect to pion pair condensation. By comparing different phenomenological models for the $\pi\pi$ interaction [4, 5] it was shown in [6] that this instability is generic and does not depend on the particular choice of kinematics or the off-shell behavior. The pion pair instability has drastic consequences for the nuclear saturation mechanism and one would like to ensure that mechanisms for providing stability have not been overlooked.

A major shortcoming of phenomenological $\pi - \pi$ interactions lies in fact that they usually do not comply with the soft pion theorems. Close inspection of these theorems indicates that the interaction is, most likely, repulsive below the $2m_\pi$ threshold. It is known [6, 7] that the so-called Σ commutator is related to the invariant π - π scattering amplitude as:

$$M^{00}(s = m_\pi^2, t = 0, u = m_\pi^2) = \frac{\Sigma_{\pi\pi}}{f_\pi^2}, \quad (1)$$

with $\Sigma_{\pi\pi}$ given by

$$\Sigma_{\pi\pi} = \langle \pi^i(\mathbf{k} = 0) | [Q_5, [Q_5, H]] | \pi^i(\mathbf{k} = 0) \rangle, \quad (2)$$

where Q_5 denotes the axial charge, H is the strong interaction Hamiltonian and f_π the usual pion decay constant. The Σ commutator can be evaluated for a Hamiltonian which explicitly breaks chiral symmetry yielding $\Sigma_{\pi\pi} = m_\pi^2$ and thus,

$$M^{00}(s = m_\pi^2, t = 0, u = m_\pi^2) = \frac{m_\pi^2}{f_\pi^2}. \quad (3)$$

At the physical threshold this result should be contrasted with the Weinberg result [8] which gives:

$$M^{00}(s = 4m_\pi^2, t = 0, u = 0) = -7 \frac{m_\pi^2}{f_\pi^2}. \quad (4)$$

This strongly suggests that somewhere below $\sqrt{s} = 2m_\pi$ the interaction turns from attraction to repulsion. In order to implement such properties we should

base our considerations on chirally invariant models for $\pi - \pi$ scattering.

In the linear sigma model, for instance, the invariant Born amplitude is given by

$$M_B^{I=0}(s, t, u) = \frac{m_\sigma^2 - m_\pi^2}{f_\pi^2} \times \left(3 \frac{s - m_\pi^2}{s - m_\sigma^2} + \frac{t - m_\pi^2}{t - m_\sigma^2} + \frac{u - m_\pi^2}{u - m_\sigma^2} \right) \quad (5)$$

satisfying eq. (3) to order m_π^4 . We also see that for $s = m_\pi^2, u = t = 0$, which can be reached in sub-threshold π - π scattering, the amplitude is even more repulsive, namely

$$M^{00}(s = m_\pi^2, t = 0, u = 0) = 2 \frac{m_\pi^2}{f_\pi^2} \left(1 - \frac{m_\pi^2}{m_\sigma^2} \right). \quad (6)$$

In order to construct a π - π interaction from the tree-level amplitude (4) to be used as an input in the scattering equation we assume that $t, u \ll m_\sigma^2$ and take the Mandelstam variables as

$$\begin{aligned} s &= E^2, \\ t &= 2m_\pi^2 - 2\omega_q\omega_{q'} + 2\mathbf{q}\mathbf{q}', \\ u &= 2m_\pi^2 - 2\omega_q\omega_{q'} - 2\mathbf{q}\mathbf{q}', \end{aligned} \quad (7)$$

with obvious notation. We then obtain the following separable π - π interaction:

$$V^{00}(E, q, q') = \frac{m_\sigma^2 - m_\pi^2}{f_\pi^2} \times \left(3 \frac{E^2 - m_\pi^2}{E^2 - m_\sigma^2} + \frac{4\omega_q\omega_{q'} - 2m_\pi^2}{m_\sigma^2} \right) v(q)v(q'). \quad (8)$$

where

$$v(q) = \left(1 + \frac{q^2}{64m_\pi^2} \right)^{-1} \quad (9)$$

is a phenomenological form factor to render intermediate state integrations finite. Adjusting the sigma mass at $m_\sigma = 1 \text{ GeV}$ yields reasonable agreement with the experimental δ_0^0 phase shifts below the $K\bar{K}$ threshold and gives $a^{00} = 0.23 m_\pi^{-1}$ for the scattering length. One should mention that, due to the separability of eq. (6), the scattering length obeys $a^{00} \propto m_\pi$ after iteration. Implementing this model for the in-medium corrections of the two-pion propagator we obtain the mass distribution (the imaginary part of the $\pi\pi$ invariant amplitude) shown in fig. 1.

It is seen that the repulsion below threshold, inherent in chiral models, prevents a strong accumulation of strength in the region $E_{\pi\pi} < 2m_\pi$. In order to judge to which degree these results are generic we have repeated the calculation within the gauged nonlinear sigma model. Unfortunately, in this case the Born term is not separable. Furthermore, when applying the Blankenbecler-Sugar reduction for the relativistic scattering equation, the chiral limit $a^{00} \rightarrow 0$

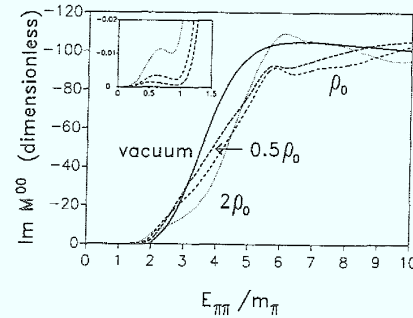


Figure 1: In-medium results for the imaginary part of the invariant M^{00} amplitude calculated in the linear σ model. Solid curve: vacuum result, long-dashed: $\rho = 0.5\rho_0$, short-dashed: $\rho = \rho_0$, dotted: $\rho = 2\rho_0$. Insert: magnified low-energy region.

as $m_\pi \rightarrow 0$ is destroyed. To preserve it we have applied a subtraction procedure to the 2π propagator:

$$\tilde{G}_{\pi\pi}(s, k) = G_{\pi\pi}(s, k) - \text{Re}G_{\pi\pi}(0, k) \quad (10)$$

which, in fact, can also be interpreted as a redefinition of the original interaction $V \rightarrow \tilde{V}$ with

$$\tilde{V} = V - V \text{Re}G_{\pi\pi}(0) \tilde{V}. \quad (11)$$

This procedure ensures a linear dependence of the scattering length on the pion mass and no instability occurs at finite density. There is, however, a considerable amount of strength in the region $m_\pi < E_{\pi\pi} < 2m_\pi$. The above scheme can also be applied to the purely phenomenological interactions such as the Jülich model [5]. In this case little strength is obtained in the subthreshold region and the strength above $2m_\pi$ becomes strongly quenched (see fig. 2). We should note that the subtraction procedure (10) is somewhat ad hoc and that other prescriptions for maintaining the chiral limit of a^{00} might not prevent the instability [9].

In conclusion, we have demonstrated here that in-medium effects on π - π scattering are probably strongly reduced when using chiral models. However, it is rather difficult to maintain chiral symmetry beyond the tree level. Symmetry preserving scattering equations from RPA like formalism are therefore currently investigated. Such efforts deem necessary to keep quantitative control over the in-medium π - π correlations which are of prime importance for the nuclear EOS.

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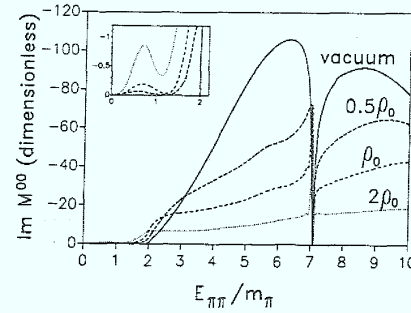


Figure 2: In-medium M^{00} amplitude calculated with the extended Jülich model: full line: vacuum results, long-dashed line: $\rho = 0.5\rho_0$, short-dashed line: $\rho = \rho_0$, dotted line: $\rho = 2\rho_0$.

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The weak production of strangeness in the reaction $NN \rightarrow \Lambda N$ is one of the processes which could be measured at COSY in Jülich. In this context it is important to know which production rates one can expect in order to decide about the feasibility of such an experiment.

We study the process $NN \rightarrow \Lambda N$ in the meson exchange framework in a DWBA calculation. We start with the one pion exchange illustrated in fig. 1(a). The weak $\Lambda N\pi$ vertex (indicated by a circle in figs. 1) is represented by an effective Hamiltonian

$$H_\pi^w = ig_w \bar{\Psi}_N (1 - \lambda \gamma_5) \vec{\tau} \Psi_\Lambda \vec{\phi}_\pi + h.c. \quad (1)$$

with coupling constants $g_w = 0.23 \times 10^{-6}$ and $\lambda = 6.9$ fixed by the free $\Lambda \rightarrow N\pi$ decay [1]. Ψ_N and $\vec{\phi}_\pi$ are the nucleon and pion fields, respectively, while $\Psi_\Lambda = \begin{pmatrix} 0 \\ \psi_\Lambda \end{pmatrix}$ is the Λ spurion field used to enforce the empirical $\Delta T = \frac{1}{2}$ rule.

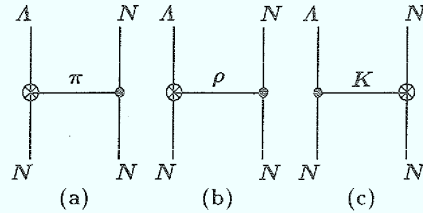


Fig. 1. (a) π exchange contribution to the reaction $NN \rightarrow \Lambda N$. (b) ρ exchange contribution. (c) K exchange contribution. The $\otimes$ indicates the weak vertex.

The ρ meson exchange contribution to the $NN \rightarrow \Lambda N$ process is depicted in Fig. 1(b). The weak $\Lambda N\rho$ vertex is described by the effective Hamiltonian

$$H_\rho^w = \bar{\Psi}_N [f_1 \gamma^\mu + i \sigma^{\mu\nu} k_\nu \frac{f_2}{2m_N} + g_1 \gamma^\mu \gamma_5] \vec{\tau} \Psi_\Lambda \vec{\phi}_{\rho,\mu} + h.c. \quad (2)$$

However, now it is not possible to determine the weak coupling parameters from experiments since heavier mesons cannot participate in the free decay of the Λ . Theoretical estimations have to be used. The weak $\Lambda N\rho$ coupling has been calculated by G. Nardulli using a pole model and information from weak nonleptonic and radiative decay [2] and we adopt his values for f_1 , f_2 , and g_1 in the present work.

For the (pseudoscalar) K -exchange contribution, shown in fig. 1(c), the weak NNK vertex is given by an effective Hamiltonian similar to eq. (1). Again theoretical estimates have to be utilized for the weak coupling parameters. Here we take the values derived in Ref. [3] using $SU(6)_W$ by which the weak NNK

vertex can be related to the experimentally known $\Lambda N\pi$ and $\Sigma N\pi$ couplings.

The parameters at the strong $NN\pi$, $NN\rho$ and ΛNK vertices (coupling constants and form factors) are taken over from our NN [4] and YN [5] interaction models, respectively. We use the same form factors at the strong and corresponding weak vertices. The interaction in the ΛN channel is taken from Ref. [4] (model A), for the NN interaction we utilize the model OBEPF from Ref. [6].

Cross sections predictions for $np \rightarrow \Lambda p$ are shown in fig 2. Since the total np cross section at energies around the ΛN threshold is about 30 mb the branching ratio for weak Λ production amounts to $R = \frac{\sigma_{np \rightarrow \Lambda p}}{\sigma_{np}} \approx 10^{-13}$ in our model calculation. Note the strong enhancement of the production cross section at $p_{lab} \approx 1150 \text{ MeV/c}$ (which is close to the ΣN threshold). This enhancement is due to a resonance in the ΛN interaction induced by the strong coupling to the ΣN channel [5].

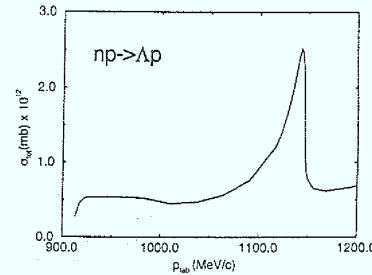


Fig. 2. Cross section predicted for the reaction $np \rightarrow \Lambda p$ process.

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NN potentials which explicitly include the delta isobar (Δ) degrees of freedom represent one of the ways how 3-body forces can be introduced into the $3N$ problem. The effect of such a $3N$ force on the $3N$ bound state has been studied thoroughly by the Hannover group and others [1, 2]. Solutions of the Faddeev equations in the $3N$ continuum are, however, still lacking. Here the large number of possible couplings between partial waves in each channel introduced by the Δ together with the more complex singularity structure of the integral kernel requires an exceedingly huge computer memory as well as a large CPU time.

A possible way to overcome these difficulties consists in the expansion of the interaction model into separable form because for such a finite-rank representation the $3N$ equations are considerably simplified. For that purpose an expansion method proposed by Ernst, Shakin, and Thaler (EST) has proven to be a powerful tool [3]. With only a few terms in the expansion the 3H binding energy but also nucleon-deuteron scattering observables could be obtained for the Paris NN potential which agree within 1% with results using the Paris model directly [4, 5].

We study separable expansions of realistic NN models with Δ degrees of freedom aiming at their application in $3N$ continuum calculations. We choose the so-called Hannover model [1] as our model potential. This model has two major advantages. Firstly only the virtual excitation of a single Δ isobar is considered. As a consequence the number of channels occurring in the NN and in the $3N$ problem is still manageable. Secondly, the Hannover model is constructed in such a way that it exactly coincides with the Paris NN potential for $E = 0$ and is (approximately) phase-equivalent to it for energies not too far from the scattering threshold. This property is important since it allows to isolate effects of the Δ isobar excitation in the $3N$ system via a comparison with results for the Paris model [1]. In the isospin singlet partial waves which do not couple to the $N\Delta$ channel the Hannover model is identified with the Paris potential.

According to EST a rank N separable representation for an interaction model V is given by [3],

$$V_N = \sum_{i,j=1}^N V | \psi(E_i) \rangle \lambda_{i,j} \langle \psi(E_j) | V, \quad (1)$$

$$\sum_{j=1}^N \lambda_{i,j} \langle \psi(E_j) | V | \psi(E_k) \rangle = \delta_{i,k}, \quad (2)$$

where the expansion functions $| \psi(E_i) \rangle$ are solutions of the Lippmann-Schwinger equation for V at arbitrary but fixed energies E_i (bound states, scattering states, or even off-shell solutions). For

NN models involving $N\Delta$ channels the functions $| \psi(E_i) \rangle$ are vectors with one or two components in the NN channel and up to four components in the $N\Delta$ channel. Accordingly, V and V_N are properly generalized matrices [3]. Note that, in general, there are several independent solutions for a particular energy, corresponding to the different choices for the boundary condition of the incoming wave, and each of them can be chosen as an expansion function. For the selection of the approximation energies E_i we can rely upon our experience with the Paris model. We use again those energy points which were found advantageous for the pure NN case and add now a few more expansion functions corresponding to solutions with incoming $N\Delta$ states. The chosen approximation energies and boundary conditions are specified in table 1. With this separable representation (NN) phase shifts are very accurately reproduced for $E_{Lab} \leq 300 MeV$ for all partial waves $J \leq 2$ [6]. Note that the isospin singlet partial waves are identical to the Paris potential and therefore the separable representations can be taken over from Ref. [4, 5]. Informations about the used approximation energies can be found there.

Here we report a first application of this separable representation to Faddeev calculations of the triton binding energy. The results are compiled in table 2 and compared with values obtained by applying the Hannover model directly [7]. It can be seen that there is nice agreement between the two calculations. We are presently working on the extension of our code to the $3N$ scattering problem.

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Table 1: Energies and boundary conditions selected for the separable expansion of the isospin triplet partial waves. The number indicating the boundary condition (b.c.) corresponds to the order in which the partial waves are listed in the heading. $N\Delta$ partial waves are given in brackets.

$^1S_0(-^5D_0)$					
Energy (MeV):	0	100	300	-100	-100
b. c.:	1	1	1	1	2
$^3P_0(-^3P_0)$					
Energy (MeV):	10	10	-50	-150	
b. c.:	1	2	1	1	
$^3P_1(-^3P_1 - ^5P_1 - ^5F_1)$					
Energy (MeV):	10	10	10	10	-50 150 300 -50
b. c.:	1	2	3	4	1 1 1 3
$^3P_2 - ^3F_2(-^3P_2 - ^3F_2 - ^5P_2 - ^5F_2)$					
Energy (MeV):	10	10	10	10	10 10 75 175
b. c.:	1	2	3	4	5 6 1 2
$^1D_2(-^5D_2 - ^5S_2 - ^5G_2)$					
Energy (MeV):	10	10	10	10	-50 150 300
b. c.:	1	2	3	4	1 1 1

Table 2: Results for the triton binding energy (in MeV).

NN input	Hannover model	separable representation
$^1S_0 - ^5D_0, ^3S_1 - ^3D_1$	-6.825	-6.841
$J \leq 1$	-7.101	-7.101
$J \leq 2$	-7.696	-7.658

Sumrule for the apparent Coulomb reacceleration of the neutrons in the electrodisintegration of the deuteron

A.Bianconi^{1,2)}, S.Jeschonnek, N.N.Nikolaev³⁾, B.G.Zakharov³⁾

In the recent experiment [1] an intriguing velocity shift between the neutrons and ${}^3\text{Li}$ fragments in the Coulomb dissociation of ${}^{11}\text{Li}$ nuclei was observed. The observation was qualitatively interpreted as an effect of classical deceleration of ${}^{11}\text{Li}$ and acceleration of ${}^3\text{Li}$ in the Coulomb field and led to the discussion of the Coulomb reacceleration as a clock for nuclear reactions.

We have investigated the contribution of final state interaction (FSI) between the fragments of electrodisintegration into the shift of the average velocity of fragments from the beam velocity[2]. We consider the shift of the velocity of neutrons in the simplest case of the Coulomb dissociation of deuterons at large momentum transfer squared Q^2 . The deuteron is particularly suited for our investigation as it is the simplest example for the mechanism of reacceleration by FSI and as there is a realistic wave function calculated from the Bonn potential [3] available and FSI in ${}^2\text{H}(e, e'p)n$ is well understood [4].

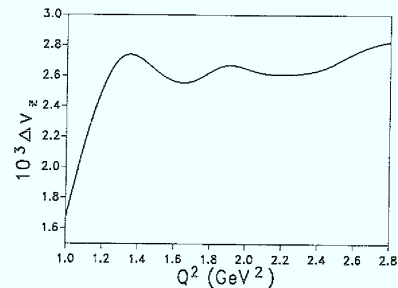


Fig. 1: The shift of the average velocity of neutrons Δv , from the velocity of deuterons in the electrodisintegration of deuterons vs. Q^2 .

We have shown that the final state p - n interaction in the reaction of electrodisintegration of the deuteron at large Q^2 in a static external field leads to the apparent reacceleration of neutrons. The shift of the neutron velocity from the velocity of the deuteron beam is related to the quantum-mechanical forward-backward asymmetry of the missing momentum distribution in the ${}^2\text{H}(e, e'p)n$ scattering. We have derived a sumrule for the velocity shift of the neutron. The calculated velocity shift is shown in Fig.1. The quantal physics of the apparent reacceleration by FSI is simple and both the acceleration and deceleration effects are possible depending on the sign of the real part of the p - n forward scattering amplitude.

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How to measure the pion structure function at HERA

H. Holtmann, G. Levman¹, N.N. Nikolaev, A. Szczurek², and J. Speth

Up to now the only feasible method to extract the pion structure was the πN Drell-Yan production. The disadvantages of this method are that the attainable luminosity is low and that only the valence part of the pion structure function at rather large x ($\gtrsim 0.2$) can be studied. An extension of our knowledge of the pion structure function is possible by using virtual pions around the nucleon as targets in deep inelastic scattering. These pions arise naturally as a consequence of the pion-nucleon coupling which leads to an admixture of the πN Fock state in the light-cone nucleon. Interaction of high-energy projectiles (nucleons, pions, leptons) with the virtual pion of the πN Fock state of the proton is a typical stripping reaction, in which the momentum distribution of the spectator nucleon reflects the momentum distribution in the πN (meson-baryon) Fock state. For instance, the differential cross section of the inclusive reaction $\pi p \rightarrow nX$ takes the form

$$\frac{d\sigma(\pi p \rightarrow nX)}{dzdp_1^2} = \frac{2}{3} f_{N/p}(z, p_1^2) \sigma_{\pi\pi}(s_{\pi\pi}), \quad (1)$$

where $f_{N/p}(z, p_1^2)$ is the flux of nucleons with the transverse momentum p_1 and the fraction z of the proton's light-cone momentum. The $\pi\pi$ total cross section enters at the center of mass energy squared $s_{\pi\pi} = (1-z)s_{\pi p}$, and the factor $\frac{2}{3}$ comes from the isotopic Clebsch-Gordan coefficients.

The salient feature of Eq. (1) is that the flux $f_{N/p}(z, p_1^2)$ as well as its counterpart $f_{\Delta/p}(z, p_1^2)$ for the $p \rightarrow \Delta$ inclusive fragmentation, is a universal property of the πN , $\pi\Delta$ Fock state, which does not depend on the projectile. Capturing on the remarkable success of the pion exchange model in hadronic reactions, we expect that the semi-inclusive reactions

$$a) \quad ep \rightarrow e'nX \quad b) \quad ep \rightarrow e'\Delta X \quad (2)$$

in the properly chosen kinematical domain also will be dominated by the pion exchange (Sullivan [1]) mechanism. If this is the case, then the straightforward generalization of Eq. (1) to semi-inclusive deep inelastic scattering is

$$\frac{d\sigma(ep \rightarrow e'nX)}{dx dQ^2 dz dp_1^2} = \frac{2}{3} f_{N/p}(z, p_1^2) K(x, Q^2) F_2^{e\pi}(x_\pi, Q^2), \quad (3)$$

where $F_2^{e\pi}(x_\pi, Q^2)$ is the structure function of the pion; $x_\pi = x/(1-z)$ is the Bjorken variable in the electron-pion deep inelastic scattering with the obvious kinematical restriction $0 < x < 1-z$, and $K(x, Q^2)$ is the standard kinematical factor

$$K(x, Q^2) = \frac{4\pi\alpha^2}{Q^4} \frac{1}{x} \left[1 - y + \frac{y^2}{2} \right], \quad y = \frac{Q^2}{xs}, \quad (4)$$

assuming for the sake of simplicity $2xF_2^{e\pi}(x) = F_2^{e\pi}(x)$. Knowing all kinematical variables and trusting the theoretical prediction for $f_{N/p}(z, p_1^2)$, one

can invert Eq.(3) and determine the pion structure function from the experimentally measured semi-inclusive cross section. In the HERA experiments, one can go down to the region of very small $x_\pi \gtrsim 10^{-4}$. This will be an enormous expansion of the studied kinematical region compared to the πN Drell-Yan experiments, which can not go much below $x_\pi \sim 0.1$ [2]. Besides, such a determination of the pion structure function at HERA will allow to study the scaling violations in the pion structure function in a broad range of (x_π, Q^2) , which hardly is possible in the Drell-Yan experiments.

From the purely experimental point of view, the semi-inclusive reaction $ep \rightarrow e'nX$ is being studied already by the ZEUS collaboration, which has installed a test forward neutron calorimeter (FNC) to complement its leading proton spectrometer [3]. This FNC was tested with neutrons from inclusive proton-beam gas interactions, and an excellent agreement between the measured spectra and the pion-exchange predictions was found.

Our principal task is to find the kinematical domain in which the semi-inclusive reaction $ep \rightarrow e'nX$ is dominated by the pion exchange contribution. The semi-inclusive production of neutrons with $z \sim 0.8$ turns out to be the optimal kinematical domain, and it also corresponds to the domain in which the semi-inclusive cross section is largest.

Above we have considered the *elastic* fragmentation $p \rightarrow n$. One can also consider the excitation of the proton into resonances and continuum states, which then will decay into the observed neutron. The prominent excitation channel is the production of Δ states, which is also dominated by the same pion (and ρ) exchange mechanism. Another type of background is the diffractive excitation of the proton into the prominent N^* resonances and/or the low mass $N\pi$, $N\pi\pi$, ... continuum states, whose decays can produce fast neutrons with large z . Such diffractively produced states carry $\approx 100\%$ of the protons momentum. The underlying mechanism is the pomeron ($\mathbb{P}$) exchange process, in which the photon interacts with the pomeron. We have shown that the background from the diffractive excitation of protons is negligibly small. Furthermore, the pion and pomeron exchange reactions differ by the rapidity distributions of secondary hadrons: in the pion-exchange reaction the rapidity gap between the observed neutron and the hadronic debris from the $\gamma^*\pi$ collision is small, $\Delta\eta \sim \log(1/(1-z)) \sim 1$, whereas in the diffractive pomeron-exchange reaction there will be a large rapidity gap between the decay products of N^* and hadrons from the $\gamma^*\mathbb{P}$ interaction, which will allow to experimentally separate the two reaction mechanisms.

In the suggested mechanism of the semi-inclusive neutron production, the differential cross section

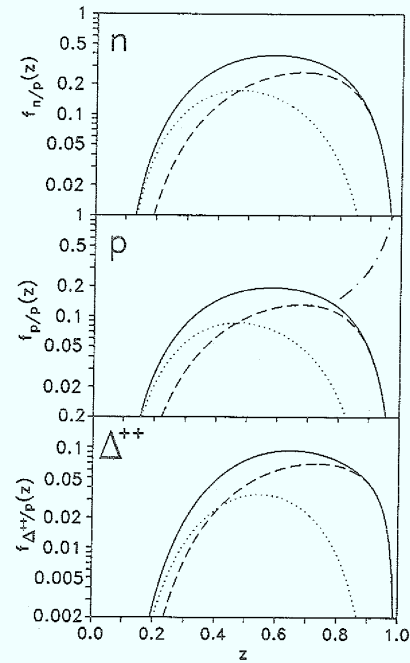


Fig. 1: Longitudinal momentum distribution of (a) neutrons, (b) protons and (c) Δ^{++} . The contributions of the π and ρ exchange mechanisms are shown by the dashed and dotted line, respectively. The contribution from the pomeron exchange mechanism to the $p \rightarrow p$ fragmentation is shown by the dash-dotted line.

Eq. (3) is a product of the universal flux factor which only depends on z and the structure function $F_2^p(x_\pi, Q^2)$ which is a function of $x_\pi = x/(1-z)$. This factorization property allows an important cross check of the model: Binning the semi-inclusive cross section data as a function of z at different fixed values of x_π one can verify that the shape of the flux factor as a function of z does not depend on x_π . Furthermore, this z -dependence must come out identical to the z -dependence of the inclusive spectra of neutrons from the hadronic pN interactions. Remarkably, the FNC of the ZEUS collaboration enables the latter cross check to be performed *in situ*, directly comparing the spectra of neutrons from inclusive beam-gas interactions and from deep inelastic ep scattering. Such a comparison of the two spectra will allow to verify that the background contribution to $z \sim 0.7-0.8$ from deep inelastic scattering off the baryonic core is as small as in hadronic reactions. Reversing the argument, one can determine the x_π dependence of the pion structure function changing x at fixed values of z and verify that this x_π dependence comes

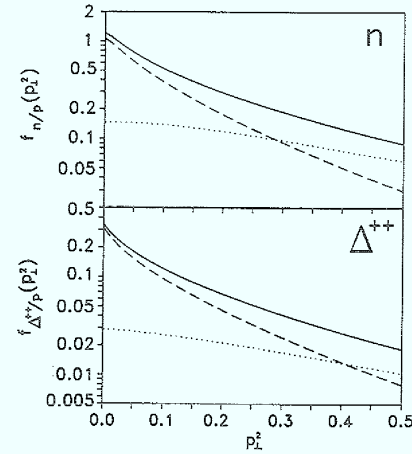


Fig. 2: Inclusive perpendicular momentum distribution of (a) neutrons and (b) Δ^{++} . The contributions corresponding to the π and ρ exchanges are shown by the dashed and dotted line, respectively.

out the same at all values of z . Our conclusion is that the separation of the pion exchange contribution and determination of the pion structure function is experimentally feasible, the possible backgrounds to the pion exchange are under reasonably good control, and important *in situ* cross checks of the pion-exchange mechanism are possible.

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Recently the CERN NA51 collaboration [1] presented the first direct measurement of the $\bar{u}/\bar{d}$ ratio

$$\frac{\bar{u}(x)}{\bar{d}(x)} = 0.51 \pm 0.04 \pm 0.05 \quad \text{at } x = 0.18. \quad (1)$$

We interpret the NA51 result as a strong piece of evidence for the pion-induced nonperturbative sea in the nucleon. The pion-nucleon interaction leads naturally to an admixture of a πN Fock component in the physical nucleon. Ever since the NMC data on the Gottfried Sum Rule (GSR) [2], there have been attempts to accommodate the observed GSR violation into the parton model analysis of structure functions invoking different parameterizations of the asymmetry. Here we wish to emphasize that as opposed to the experiment-driven parameterizations, the dynamical model of the pion-induced sea gives a unique prediction for the $\bar{u}-\bar{d}$ asymmetric part of the sea using pion-nucleon interaction parameters inferred from low and high energy soft hadronic interactions. Furthermore, the region of $x \sim 0.2$ is precisely the region where the nucleon sea is dominated by the pion-induced nonperturbative sea. The predictive power of the model rests on an important observation that the pion exchange contribution enters the high energy nucleon (delta) production in hadronic reactions and the lepton deep inelastic scattering virtual pions in similar kinematical conditions. Both processes are described by the same light cone πN ($\pi\Delta$, $K\Lambda$, ...) wave function, which allows one to constrain the principal nonperturbative parameters—the radii of the meson-baryon light cone wave functions—from an analysis of the $p \rightarrow n, \Delta, \Lambda$, etc. fragmentation spectra [3]. The meson-baryon couplings, too, are nonperturbative parameters of the model; in our analysis we use the standard low energy determinations of these couplings.

In the region of x of interest, the meson-induced sea is dominated by the valence quarks of the pion. Indeed, typically $f_{\pi/N}(y)$ has a maximum at $y \sim 0.3$, and $x = 0.18$ studied in the NA51 experiment corresponds to the rather large Bjorken variable in deep inelastic scattering by pions, $x_\pi \simeq x/y \simeq 0.6$. Thus the recent parameterization [4] based on the πN Drell-Yan data can be regarded as reliable, and the meson-induced sea can be calculated with good accuracy.

The primordial sea distribution of the nucleonic core is the least known ingredient of the model. As a first approximation, we adopt the following phenomenological procedure: We assume the primordial sea in the nucleonic core to be $\bar{u}-\bar{d}$ symmetric with the shape given by the symmetric MRS (S'_0) [5] parameterization and the magnitude adjusted so as to obtain rough agreement with the absolute normalization of the MRS(D'_0, D'_+) [5] $\bar{u}$ and $\bar{d}$ distributions, which are shown in Fig. 1 by dotted lines. Such a proce-

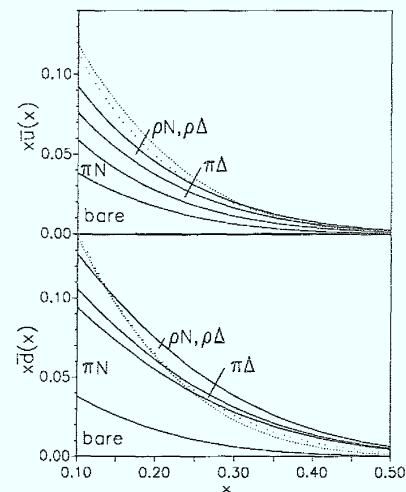
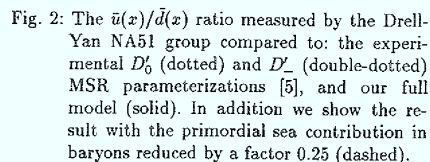


Fig. 1: The contributions of various Fock states to the $x\bar{u}$ (Fig. 1a) and $x\bar{d}$ (Fig. 1b) distributions. The dotted curves correspond to the asymmetric quark distributions D'_0 (dotted), D'_+ (double-dotted) from Ref.[5].

dures suggests that the primordial sea is about 50% of that given by the symmetric MRS (S'_0) parameterizations. This procedure is very crude, since the MRS(D'_0, D'_+) parameterization does not reproduce the NA51 result, but the effects of the primordial sea turn out to be small. The resulting decomposition of the $\bar{u}(x)$ and $\bar{d}(x)$ densities into contributions from different mechanisms is shown in more detail in Fig. 1.

Our results for the $\bar{u}/\bar{d}$ ratio are shown in Fig. 2. Because the primordial sea only is a correction to the dominant meson-induced sea, the predicted $\bar{u}/\bar{d}$ ratio changes little if the normalization of the primordial sea in nucleons is changed by a factor 2 from 50% to 25%. Reversing this argument, the ratio $\bar{u}/\bar{d}$ is stable to the factor ~ 2 changes of the absolute flux of pions. Owing to high accuracy of the hadronic data successfully described in the same framework a very conservative estimate of the normalization uncertainty is $\leq (20-30)\%$. For comparison we also show $\bar{u}/\bar{d}$ ratio obtained from the MSR (D'_0, D'_+) parameterizations [5] of the previous deep inelastic scattering and the Drell-Yan production data. These parameterizations were fine-tuned to the GSR violation, but placed the asymmetry at small values of x .



The model of the meson-induced sea gives a unique prediction that the $\bar{u}-\bar{d}$ asymmetry is predominantly a large- x phenomenon, in good agreement with the NA51 result.

References:

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Whether the QCD pomeron is described by the fixed or moving singularity in the complex j -plane, remains one of the topical issues. In this communication we prove that the generalized BFKL pomeron [1,2] is a moving cut and present the first direct calculation of the slope α'_{IP} for the pomeron trajectory. Being dimensionfull, the slope α'_{IP} is the nonperturbative quantity, and we relate it to the nonperturbative parameter of the generalized BFKL pomeron - the correlation radius for the perturbative gluons. The starting point of our analysis is the generalization of our BFKL equation [1,2] to the profile function of the dipole cross section $\Gamma(r, \vec{b})$. Defining the impact parameter $\vec{b}$ with respect to the center of the parent $q-\bar{q}$ dipole, and repeating the derivation [1,2], we obtain the homogeneous equation

$$\begin{aligned} \frac{\partial \Gamma(\xi, r, \vec{b})}{\partial \xi} = \mathcal{K} \otimes \Gamma(\xi, r, \vec{b}) = \frac{3}{8\pi^3} \int d^2 \vec{\rho}_1 \mu_G^2 \\ \left| g(R_1) K_1(\mu_G \rho_1) \frac{\vec{\rho}_1}{\rho_1} - g(R_2) K_1(\mu_G \rho_2) \frac{\vec{\rho}_2}{\rho_2} \right|^2 \\ \times \left[\Gamma(\xi, \rho_2, \vec{b} + \frac{1}{2} \vec{\rho}_1) + \Gamma(\xi, \rho_1, \vec{b} + \frac{1}{2} \vec{\rho}_2) - \Gamma(\xi, r, \vec{b}) \right]. \quad (1) \end{aligned}$$

in which $\vec{\rho}_2 = \vec{\rho}_1 - \vec{r}$, the arguments of the running QCD charge $g(r) = \sqrt{4\pi\alpha_S(r)}$ are $R_i = \min\{r, \rho_i\}$, $K_1(x)$ is the generalized Bessel function and $R_c = 1/\mu_G$ is the correlation radius for the perturbative gluons. Here we use the standard definition of the profile function when $A(s, t) = 2is \int d^2 \vec{b} \exp(-i\vec{q}\vec{b}) \Gamma(\vec{b})$ and the dipole cross section equals $\sigma(\xi, r) = 2 \int d^2 \vec{b} \Gamma(\xi, r, \vec{b})$. Hereafter we shall discuss the reduction of (2) to the equation for the diffraction slope $B(\xi, r) = \frac{1}{2} \langle \vec{b}^2 \rangle = \lambda(\xi, r)/\sigma(\xi, r)$, where $\lambda(\xi, r) = \int d^2 \vec{b} \vec{b}^2 \Gamma(\xi, r, \vec{b})$. Evidently, the diffraction slope for the dipole of size r contains the purely geometrical contribution $\frac{1}{8} r^2$ which comes from the elastic form factor of the dipole. Then, it is more convenient to consider $\eta(\xi, r) = \lambda(\xi, r) - \frac{1}{8} r^2 \sigma(\xi, r)$, for which we obtain the inhomogeneous equation

$$\begin{aligned} \frac{\partial \eta(\xi, r)}{\partial \xi} = \frac{3}{8\pi^3} \int d^2 \vec{\rho}_1 \mu_G^2 \\ \left| g(R_1) K_1(\mu_G \rho_1) \frac{\vec{\rho}_1}{\rho_1} - g(R_2) K_1(\mu_G \rho_2) \frac{\vec{\rho}_2}{\rho_2} \right|^2 \\ \times \{ \eta(\xi, \rho_1) + \eta(\xi, \rho_2) - \eta(\xi, r) + \\ \frac{1}{8} (\rho_1^2 + \rho_2^2 - r^2) [\sigma(\rho_2, \xi) + \sigma(\rho_1, \xi)] \} \\ = \mathcal{K} \otimes \eta(\xi, r) + \beta(\xi, r). \quad (2) \end{aligned}$$

Here the crucial point is that the homogeneous Eq. (2) is precisely our generalized BFKL equation

for the dipole cross section

$$\frac{d\sigma(\xi, r)}{d\xi} = \mathcal{K} \otimes \sigma(\xi, r), \quad (3)$$

which enables us to prove that on a very generic grounds that $\alpha'_{\text{IP}} = \frac{1}{2} dB(\xi, r)/d\xi \neq 0$ and, consequently, the generalized BFKL pomeron is described by the moving cut in the complex j -plane. The proof goes as follows: In [3,4] we have shown that the generalized BFKL operator $\mathcal{K}$ has the continuous spectrum, which corresponds to the cut in the j -plane. Let $-\infty < \nu < \infty$ be the "wavenumber" which labels eigenfunctions $E(\nu, r) \exp[\Delta(\nu)\xi]$ of Eq. (3) with eigenvalue $\Delta(\nu)$. For the guidance, in the scaling limit of $\alpha_S = \text{const}$ and $R_c \rightarrow \infty$ $E(\nu, r) = r \exp[i\nu \log(r^2)] = \sigma_{\text{IP}}(r) \exp[i\nu \log(r^2)]$ with the orthogonality condition [3-5]

$$\delta(\nu - \mu) = \frac{1}{2\pi} \int \frac{d \log(r^2)}{[\sigma_{\text{IP}}(r)]^2} E^*(\nu, r) E(\mu, r), \quad (4)$$

and ν is indeed the wavenumber of plane waves in the $\log(r^2)$ space.

The large- ξ behaviour of solutions of the inhomogeneous equation (2) is controlled by properties of the inhomogeneous term

$$\begin{aligned} \beta(\xi, r) = \mathcal{L} \otimes \sigma(\xi, r) = \frac{3}{64\pi^3} \int d^2 \vec{\rho}_1 \mu_G^2 \\ \left| g(R_1) K_1(\mu_G \rho_1) \frac{\vec{\rho}_1}{\rho_1} - g(R_2) K_1(\mu_G \rho_2) \frac{\vec{\rho}_2}{\rho_2} \right|^2 \\ (\rho_1^2 + \rho_2^2 - r^2) [\sigma(\rho_2, \xi) + \sigma(\rho_1, \xi)]. \quad (5) \end{aligned}$$

If $G(\nu, r) = \mathcal{L} \otimes E(\nu, r) = \int d\omega g(\omega, \nu) E(\omega, r)$, then the inhomogeneous term (5) can be written as

$$\begin{aligned} \beta(\xi, r) = \mathcal{L} \otimes \sigma(\xi, r) = \\ \int d\nu E(\nu, r) \int d\omega f(\omega) g(\omega, \nu) \exp[\Delta(\omega)\xi]. \quad (6) \end{aligned}$$

The singularity structure of $g(\omega, \nu)$ can be found considering the large- r behaviour of $G(\nu, r) = \mathcal{L} \otimes E(\nu, r)$. Making use of properties of solutions of the generalized BFKL solutions found in [2], it can be shown that $g(\omega, \nu)$ contains the two singular components $g_1(\omega, \nu) = 2G_1 \delta(\omega - \nu)$ and $g_2(\omega, \nu) = G_2(\nu) \delta(\omega)$ besides the smooth function of ω and ν . The $2G_1 \delta(\omega - \nu)$ component of $g(\omega, \nu)$ gives a contribution to $\eta(\xi, r)$ of the form

$$\begin{aligned} \eta_1(\xi, r) = \\ 2G_1 \int_0^\xi d\xi' \int d\nu f(\nu) E(\nu, r) \exp[\Delta(\nu)\xi] = \\ = 2G_1 \xi \sigma(\xi, r), \quad (7) \end{aligned}$$

which gives precisely the Regge growth of the diffraction slope $B(\xi, r)$ with $\alpha'_{\text{IP}} = G_1$. Furthermore, we find an explicit estimate for the slope of the

pomeron trajectory in terms of the nonperturbative correlation radius R_c ,

$$\alpha'_{\mathbb{P}} \sim \frac{3}{16\pi^2} \int d^2\vec{r} \alpha_S(r) \mu_G^2 r^2 K_1^2(\mu_G r) \propto \frac{3}{64\pi} R_c^2 \alpha_S(R_c). \quad (8)$$

The effect of $g_2(\omega, \nu) = G_2(\nu)\delta(\omega)$ can be evaluated making use of the explicit form of the pomeron cross section found in [2,3] [3]. It also contributes to the slope of the pomeron trajectory $\alpha'_{\mathbb{P}} \sim G_2(0) \sim G_1$. The smooth, nonsingular, part of $g(\omega, \nu)$ does not contribute to the slope of the pomeron trajectory.

In the numerical calculation of the slope $\alpha'_{\mathbb{P}}$ we start with the dipole-dipole cross section and the corresponding diffraction slope as described in [5,6]. We calculate the ξ dependence of the dipole cross section $\sigma(\xi, r)$ and of the diffraction slope $B(\xi, r)$ and verify that at $\xi \rightarrow \infty$ the effective intercept $\Delta_{eff}(\xi, r) = \partial \log \sigma(\xi, r) / \partial \xi$ and the effective slope $\alpha'_{eff}(\xi, r) = \partial B(\xi, r) / \partial \xi$ tend to the limiting values $\Delta_{\mathbb{P}}$ and $\alpha'_{\mathbb{P}}$, respectively, which are independent of the size of the projectile and target colour dipoles. We take the running coupling with the infrared freezing $\alpha_S(r) \leq \alpha_S^{(fr)} = 0.82$ [3,4]. The dependence of the slope $\alpha'_{\mathbb{P}}$ on $\mu_G = 1/R_c$ is shown in Fig.1.

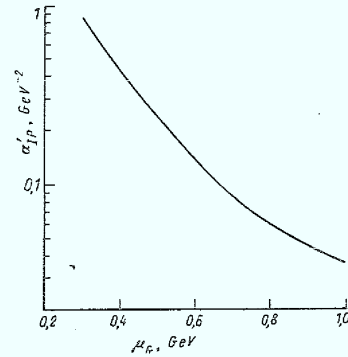


Fig. 1: The slope of the pomeron trajectory $\alpha'_{\mathbb{P}}$ as a function of the inverse correlation radius $\mu_G = 1/R_c$ for the perturbative gluons.

The simple estimate (8) is close to these numerical results.

To summarize, we have shown that the generalized BFKL pomeron [1-4] is the moving cut in the complex angular momentum plane. The slope $\alpha'_{\mathbb{P}}$ of the pomeron trajectory is the dimensionfull nonperturbative quantity and we related it to the nonperturbative parameter of the model - the gluon correlation radius.

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At moderately high energies, the salient feature of diffraction dissociation $a + p \rightarrow X + p'$ of ($a = h$) hadrons and ($a = \gamma$) real photons ($Q^2 = 0$) is the so called triple-pomeron regime

$$\frac{M^2}{\sigma_{tot}(ap)} \cdot \frac{d\sigma_D(a \rightarrow X)}{dt dM^2} \Big|_{t=0} \approx A_3 \mathbb{IP} \quad (1)$$

with the approximately energy- and projectile-independent dimensionfull triple-pomeron coupling $A_3 \mathbb{IP}$. Here t is the (p, p') momentum transfer squared and the mass M of the excited state satisfies $m_p^2 \ll M^2 \ll W^2$, where W is the total c.m.s. energy. The FNAL data on the diffractive real photoproduction give $A_3 \mathbb{IP}(Q^2 = 0) \approx 0.16 \text{ GeV}^{-2}$ [1]. In the Regge theory language, $A_3 \mathbb{IP}$ measures the projectile-pomeron cross section [2]

$$\sigma_{tot}(a\mathbb{IP}) = 16\pi A_3 \mathbb{IP} \frac{\sigma_{tot}(ap)}{\sigma_{tot}(pp)} \approx 0.08 \sigma_{tot}(ap). \quad (2)$$

Why $A_3 \mathbb{IP}$ is small and $\sigma_{tot}(a\mathbb{IP})$ is more than one order in magnitude smaller than $\sigma_{tot}(aN)$, is one of outstanding mysteries of the pomeron.

In diffractive DIS too, one can operationally define the triple-pomeron coupling

$$\frac{M^2 + Q^2}{\sigma_{tot}(\gamma^* p)} \cdot \frac{d\sigma_D(\gamma^* \rightarrow X)}{dt dM^2} \Big|_{t=0} \approx A_3 \mathbb{IP}(Q^2). \quad (3)$$

The dimensionfull $A_3 \mathbb{IP}(Q^2)$ is the nonperturbative quantity. The subject of the present paper [3] is an evaluation of the Born approximation for $A_3 \mathbb{IP}(Q^2)$ in terms of the nonperturbative parameter of the dipole cross section approach to the QCD pomeron [4,5]: the correlation radius R_c for the perturbative gluons. We rely upon the microscopic dipole pomeron description of diffractive DIS [4,5]. The $\propto 1/(M^2 + Q^2)$ component of the mass spectrum comes from the diffraction excitation of the $q\bar{q}g$ Fock state of the photon containing the soft gluon which carries the fraction $z_g \ll 1$ of photon's lightcone momentum and gives rise to $(M^2 + Q^2) \propto Q^2/z_g \gg Q^2$. Let $\vec{r}, \vec{\rho}_1$ and $\vec{\rho}_2 = \vec{\rho}_1 - \vec{r}$ be the $\bar{q}$ - q , g - q and g - $\bar{q}$ separations in the impact parameter (transverse size) plane. Then, in the triple-pomeron regime

$$(Q^2 + M^2) \frac{d\sigma_D}{dt dM^2} \Big|_{t=0} = \int dz d^2\vec{r} d^2\vec{\rho}_1 \{z_g |\Phi(\vec{r}, \vec{\rho}_1, \vec{\rho}_2, z, z_g)|^2\}_{z_g=0} \frac{\sigma_3^2}{16\pi} \quad (4)$$

where σ_3 stands for the 3-body cross section [5]

$$\sigma_3(r, \rho_1, \rho_2) = \frac{9}{8} [\sigma(\rho_1) + \sigma(\rho_2)] - \frac{1}{8} \sigma(r), \quad (5)$$

the square of the 3-parton wave function equals [5]

$$|\Phi(\vec{r}, \vec{\rho}_1, \vec{\rho}_2, z, z_g)|^2 = \frac{1}{z_g} \frac{1}{3\pi^3} |\Psi_{\gamma^*}(Q^2, z, r)|^2 \mu_G^2 \left| g(R_1) K_1(\mu_G \rho_1) \frac{\vec{\rho}_1}{\rho_1} - g(R_2) K_1(\mu_G \rho_2) \frac{\vec{\rho}_2}{\rho_2} \right|^2. \quad (6)$$

Here $|\Psi_{\gamma^*}(Q^2, z, r)|^2$ is the lightcone wave function of the quark-antiquark state derived in [6], $g(r)$ is the running color charge, $\alpha_S(r) = g^2(r)/4\pi$, the arguments of color charges are $R_i = \min\{r, \rho_i\}$ and $R_c = 1/\mu_G$ is the correlation (propagation) radius for perturbative gluons.

In the case of DIS, the typical size of the $q\bar{q}$ dipole $r^2 \lesssim R_{q\bar{q}}^2 \propto 1/Q^2$ and, to the standard Leading-Log Q^2 approximation, the dominant contribution to the diffraction cross section (4) comes from $r^2 \ll \rho_1^2 \approx \rho_2^2 \sim R_c^2$. In this region $\sigma_3(r, \rho_1, \rho_2) \approx \frac{9}{4} \sigma(\rho) \gg \sigma(r)$, where $\rho = \frac{1}{2}(\rho_1 + \rho_2)$, the virtual photon interacts as an effective octet-octet color dipole of size $\rho \sim R_c$, with the $q\bar{q}$ pair acting as an octet color charge, the 3-parton wave function factorizes and (4) takes on the factorized form

$$(Q^2 + M^2) \frac{d\sigma_D}{dt dM^2} \Big|_{t=0} = \int dz d^2\vec{r} |\Psi_{\gamma^*}(Q^2, z, r)|^2 \cdot \frac{16\pi^2}{27} \cdot \alpha_S(r) r^2 \times \frac{1}{2\pi^4} \cdot \left(\frac{9}{8}\right)^3 \cdot \int d\rho^2 \left[\frac{\sigma(\rho)}{\rho^2}\right]^2 \mathcal{F}(\mu_G \rho), \quad (7)$$

where $\mathcal{F}(z) = z^2 [K_1^2(z) + z K_1(z) K_0(z) + \frac{1}{2} z^2 K_0^2(z)]$. Because modulo to the factor $\sim \log[1/\alpha_S(Q^2)] \sim 1$,

$$\int dz d^2\vec{r} |\Psi_{\gamma^*}(Q^2, z, r)|^2 \cdot \frac{16\pi^2}{27} \cdot \alpha_S(r) r^2 \approx \sigma_{tot}(\gamma^* p) = \int dz d^2\vec{r} |\Psi_{\gamma^*}(Q^2, z, r)|^2 \sigma(r), \quad (8)$$

Consequently, (7) becomes equivalent to (3) with

$$A_3 \mathbb{IP}(Q^2) \sim A_3^* \mathbb{IP} = \frac{1}{2\pi^4} \cdot \left(\frac{9}{8}\right)^3 \cdot \int d\rho^2 \left[\frac{\sigma(\rho)}{\rho^2}\right]^2 F(\mu_G \rho). \quad (9)$$

The factorization (7) holds simultaneously for all the $q\bar{q}g$ states and we predict independence of $A_3 \mathbb{IP}(Q^2)$ on the flavour " i ". Eq. (9) shows very clearly that $A_3^* \mathbb{IP}$ is the nonperturbative quantity dominated by $\rho \sim R_c$.

For real photons, $Q^2 = 0$, in the dipole distribution (6) the typical $q\bar{q}$ size is large,

$$r^2 \sim R_{q\bar{q}}^2 \approx \frac{1}{m_q^2} \gg R_c^2. \quad (10)$$

Because of inequality (10), the dipole distribution $|\Phi|^2$ will be dominated by configurations with $\rho_1^2 \lesssim R_c^2 \ll \rho_2^2 \sim r^2$ and $\rho_2^2 \lesssim R_c^2 \ll \rho_1^2 \sim r^2$ and takes on the factorized form first considered in [4]:

$$|\Phi(\vec{r}, \vec{\rho}_1, \vec{\rho}_2, z, z_g)|^2 = \frac{1}{z_g} \frac{4}{3\pi^2} |\Psi_{\gamma^*}(Q^2, z, r)|^2 \mu_G^2 [\alpha_S(\rho_1) K_1^2(\mu_G \rho_1) + \alpha_S(\rho_2) K_1^2(\mu_G \rho_2)]. \quad (11)$$

We recover a sort of the constituent quark model with the $q\bar{q}$ and/or $q\bar{q}g$ clusters of size $\rho \lesssim R_c$, with

the square of the $q\bar{q}$ wave function of the "constituent" quark $\propto \frac{1}{z_i} \alpha_s(\rho) K_1^2(\mu_G \rho)$. Diffraction dissociation of the photon $\gamma^* \rightarrow q + \bar{q} + g$ proceeds via diffraction excitation of the "constituent" (anti)quark of the photon $q(\bar{q}) \rightarrow q(\bar{q}) + g$. For the size ordering (10), Eq. (5) leads to $\sigma_3(r, \rho_1, \rho_2) \approx \sigma(r) + \frac{9}{8} \sigma(\rho)$, the factorization $\sigma_3^2(r, \rho_1, \rho_2) - \sigma^2(r) \approx \frac{9}{4} \sigma(r) \sigma(\rho)$, where $\rho = \min\{\rho_i\}$, and the factorized form of the diffraction cross section (4):

$$M^2 \left. \frac{d\sigma_D}{ddM^2} \right|_{t=0} \approx \int dz d^2\vec{r} |\Psi_{\gamma^*}|^2 \sigma(r) \frac{3}{8\pi^2} \cdot \int d\rho^2 \left[\frac{\sigma(\rho)}{\rho^2} \right] f^2(\mu_G \rho) = \sigma_{tot}(\gamma p) A_{3\mathbb{P}}(0), \quad (12)$$

in which

$$A_{3\mathbb{P}}(0) \approx \frac{3}{8\pi^2} \int d\rho^2 \alpha_s(\rho) \left[\frac{\sigma(\rho)}{\rho^2} \right] f^2(\mu_G \rho) \quad (13)$$

and $f(z) = z K_1(z)$. Notice a close similarity between equations (9) and (13). A direct calculation gives $A_{3\mathbb{P}}(0) = 0.23 \text{ GeV}^{-2}$, in reasonable agreement with the experimental determination $A_{3\mathbb{P}}(0) \approx 0.16 \text{ GeV}^{-2}$ [1].

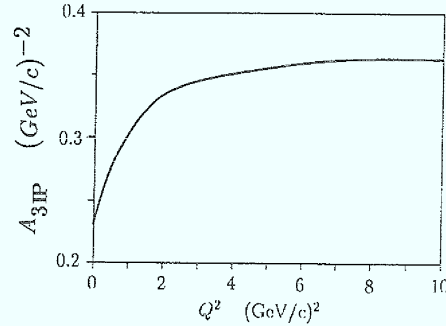


Fig. 1: Our prediction for the Q^2 dependence of $A_{3\mathbb{P}}(Q^2)$.

In Fig. 1 we present the Q^2 dependence of the triple-pomeron coupling $A_{3\mathbb{P}}(Q^2)$. The main feature of $A_{3\mathbb{P}}(Q^2)$ is its weak Q^2 dependence; still, we predict a slight, by a factor ~ 1.6 , growth of $A_{3\mathbb{P}}(Q^2)$ from $A_{3\mathbb{P}}(Q^2 = 0) = 0.23 \text{ GeV}^{-2}$ to the DIS value $A_{3\mathbb{P}}(Q^2 \gtrsim Q^2) = 0.36 \text{ GeV}^{-2}$. This rise of $A_{3\mathbb{P}}(Q^2)$ takes place at $Q^2 \lesssim Q^2 \approx (2-3) \text{ GeV}^2$ and describes the transition from the regime of diffraction dissociation of the "constituent" quark to the regime of diffraction dissociation of the octet-octet dipole state of the photon. The experimental test of the predicted Q^2 dependence of $A_{3\mathbb{P}}(Q^2)$ is possible at HERA and will shed light on our understanding of dipole cross section at $r \sim R_c$. To summarize, with the nonperturbative parameter of the model - the correlation (propagation) radius for the perturbative gluons $R_c \approx 0.3 \text{ fm}$ as used in

other successful phenomenological application of the model, we found good agreement with the experimental determination of $A_{3\mathbb{P}}(0)$. We related the small numerical value of $A_{3\mathbb{P}}$, which was a mystery, to the small value of R_c .

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Nuclear effects remain one of major open issues in the theory of heavy-flavor production in hadronic interactions. In leading order pQCD the cross section for heavy quark-antiquark production can be written as

$$\sigma_{Q\bar{Q}}(s) = \sum_{i,j} \int dx_1 dx_2 f_{h_1}^i(x_1, \mu^2) f_{h_2}^j(x_2, \mu^2) \cdot \sigma_{ij}(x_1 x_2 s, \mu^2). \quad (1)$$

Here $f_{h_1}^i(x_1, \mu^2)$ and $f_{h_2}^j(x_2, \mu^2)$ are the densities of partons "i" and "j" in the projectile h_1 and target h_2 respectively, taken at the factorization scale $\mu^2 \sim 4m_Q^2$, where m_Q is the mass of the heavy-quark. $\sigma_{ij}(x_1 x_2 s, \mu^2)$ is the cross section of the partonic subprocess $i + j \rightarrow Q + \bar{Q}$ at a squared center of mass energy $x_1 x_2 s > 4m_Q^2$.

In this work [1] we develop a theory of coherent nuclear effects at $x_2 \lesssim 0.1 \cdot A^{-1/3}$, in the domain of the gluon fusion subprocess $g + g \rightarrow Q + \bar{Q}$. In this region, heavy-flavor hadroproduction in the laboratory can be described in our formalism of light-cone wave functions (LCWF) [2,3] as diffractive dissociation of projectile gluons into $Q\bar{Q}$ pairs. The physical gluon in the LCWF formalism has a non trivial Fock decomposition which also contains a heavy $Q\bar{Q}$ state

$$|G\rangle = \sqrt{1 - n_{Q\bar{Q}}|g\rangle} + \sum_{z, \vec{r}} \Psi_G(z, \vec{r}) |Q\bar{Q}; z, \vec{r}\rangle. \quad (2)$$

Here $\Psi_G(z, \vec{r})$ is the projection of the physical gluon wave function onto a $Q\bar{Q}$ state in the $(z, \vec{r})$ -representation, where the light-cone variable z represents the fraction of the gluon momentum carried by the quark, $\vec{r}$ is the transverse separation of the Q - $\bar{Q}$ in the impact parameter space and

$$n_{Q\bar{Q}} = \int_0^1 dz \int d^2 \vec{r} |\Psi_G(z, \vec{r})|^2, \quad (3)$$

The $Q\bar{Q}$ wave function Ψ_G of the gluon can be obtained immediately from the $Q\bar{Q}$ wave function of a photon Ψ_γ , as derived in [2] After the interaction with the target nucleus at an impact parameter $\vec{b}$, the incident gluon state is transformed into

$$|G; \vec{b}\rangle \Rightarrow \hat{S}(\vec{b}) |G; \vec{b}\rangle. \quad (4)$$

$\hat{S}(\vec{b})$, the scattering matrix in the impact parameter representation, acts differently on the g and $Q\bar{Q}$ components of the physical gluon G , which leads to a production of the $Q\bar{Q}$ state. Unitarity of the $\hat{S}$ -matrix at a fixed impact parameter $\vec{b}$ simplifies the calculation of the inclusive $Q\bar{Q}$ production cross section. For every value of $\vec{b}$ it relates the probability $P_{Q\bar{Q}}$, to detect a $Q\bar{Q}$ pair in the final state with the probability P_G to find a physical gluon $P_G + P_{Q\bar{Q}} = 1$. After some algebra, making use of the explicit form

of the $\hat{S}$ -matrix, we find

$$P_{Q\bar{Q}} = 2 \int dz d^2 \vec{r} |\Psi_G(z, \vec{r})|^2 \langle (Q\bar{Q}g)_1 N | [1 - \hat{S}(\vec{b})] | (Q\bar{Q}g)_1 N \rangle. \quad (5)$$

Here, the crucial finding is that $P_{Q\bar{Q}}$ takes the form of a profile function for the elastic scattering of a $(Q\bar{Q}g)_1$ state on a nucleon, so that

$$\sigma(GN \rightarrow Q\bar{Q}X) = \int dz d^2 \vec{r} |\Psi_G(z, \vec{r})|^2 \sigma_3(\vec{r}_1, \vec{r}_2) = \langle \sigma_{3N} \rangle_{Q\bar{Q}g}, \quad (6)$$

where

$$\sigma_{3N} = \frac{9}{8} \left[\sigma(r_1) + \sigma(r_2) - \frac{1}{9} \sigma(r) \right]. \quad (7)$$

As indicated σ_{3N} is the total interaction cross section for the scattering of a color-singlet three-parton state $(Q\bar{Q}g)_1$ from a nucleon [3]. The g - Q and g - $\bar{Q}$ separation in the impact parameter space are labeled $\vec{r}_1$ and $\vec{r}_2$ respectively. Their difference $\vec{r} = \vec{r}_1 - \vec{r}_2$ is identical to the Q - $\bar{Q}$ separation. The distances $\vec{r}_{1,2}$ are related to the light-cone momentum fractions carried by the Q and $\bar{Q}$ respectively:

$$\vec{r}_2 = -z\vec{r}, \quad \vec{r}_1 = (1-z)\vec{r}. \quad (8)$$

The cross section $\sigma(r)$ is for the interaction of a $Q\bar{Q}$ color-dipole of size $r = |\vec{r}|$ with a nucleon. The color transparency property of the dipole cross section provides the infrared and ultraviolet finite heavy-quark production cross section in spite of the dissociating gluon being a colored state without any need for color cancellations by comoving quarks.

In order to calculate the heavy-quark production cross section for hadron-nucleon collisions $\sigma_{Q\bar{Q}}(h, N)$, we have to multiply $\sigma(GN \rightarrow Q\bar{Q}X)$ from Eq. (6) with the gluon density of the incoming hadron projectile and integrate over the virtuality k_1^2 of the projectile gluon. Making use of the relationship between the dipole cross section and the gluon density in the target nucleus [3], we recover the standard parton model formula (1).

The above formalism comes truly of age when applied to production off nuclear targets. Generalization of the above to nuclear targets is straightforward and reduces to a substitution of the Glauber cross section

$$\sigma_{3A} = 2 \int d^2 \vec{b} \left\{ 1 - \exp \left[-\frac{1}{2} \sigma_{3N} T(\vec{b}) \right] \right\}. \quad (9)$$

for σ_{3N} . Here $\vec{b}$ is the impact parameter of the $(Q\bar{Q}g)$ -nucleus scattering process, $T(\vec{b})$ stands for the optical thickness of the nucleus

$$T(\vec{b}) = \int_{-\infty}^{+\infty} dz n_A(\vec{b}, z). \quad (10)$$

Expanding the exponential in (9) in powers of

$$\tau_A = \sigma_{3N} T(\vec{b}) \propto \sigma_{3N} A^{1/3}. \quad (11)$$

one can identify terms $\propto \tau_A^k$ as resulting from the coherent interaction of the $(Q\bar{Q}g)$ state with ν nucleons inside the target nucleus. In the impulse approximation, nuclear transparency $T_A = 1$. The driving contribution to nuclear attenuation results from the coherent interaction of the three-parton $(Q\bar{Q}g)$ state with two nucleons inside the target:

$$T_A = 1 - \frac{1}{4} \frac{\langle \sigma_{3N}^2 \rangle_{Q\bar{Q}g}}{\langle \sigma_{3N} \rangle_{Q\bar{Q}g}} \int d^2\vec{b} T^2(\vec{b}). \quad (12)$$

At $k_1^2 \ll 4m_Q^2$ shadowing effects can easily be estimated, using the fact that $|\Psi_G(z, r)|^2 \propto \exp(-2m_Q r)$ and $\sigma(r) \approx C r^2$:

$$1 - T_A \approx \frac{15CA}{32\pi R_{ch}^2 m_Q^2}. \quad (13)$$

In the accessible region of small $x_2 \sim 0.01$ we find at $r \sim 1/m_c$ typically $C \approx 2.5$ [4], which leads to

$$1 - T_A \sim 6 \cdot 10^{-3} \cdot A^{1/3}. \quad (14)$$

in good agreement with the experimental data [5]. Following the analysis of refs.[2,6] we note, that the shadowing effect in Eq.(12) does not vanish at high virtualities of the projectile gluon $k_1^2 \gg 4m_Q^2$, and consequently is a leading twist effect. (The same is true for nuclear shadowing in DIS.) There is a close similarity between nuclear shadowing in heavy-flavor hadroproduction and heavy-flavor contributions to nuclear structure functions. It mainly involves an exchange of the three-parton cross section σ_{3N} in (12) with the dipole cross section $\sigma(r)$, which differs in the considered process by a factor ~ 2 . Hence the A -dependence of both processes cannot be described in terms of modified nuclear gluon distributions, which would be unique for all hard processes. More detailed analysis has shown that shadowing of $Q\bar{Q}$ contributions in DIS must be reinterpreted as a modification of the nuclear sea quark distributions [3]. This, however, is not possible for heavy-flavor hadroproduction.

Shadowing effects, which can be ascribed to nuclear modifications of gluon structure functions, emerge in DIS first through contributions from higher order $Q\bar{Q}g$ Fock states of the photon. The contributions of $Q\bar{Q}g$ components to shadowing can be related to diffractive dissociation of virtual photons into large mass states $M^2 \gg Q^2$, which is successfully described in terms of the so-called triple-pomeron coupling $A_3\mathbb{P}$. At $x_2 \lesssim x_A = 0.1 A^{1/3}$ one finds [2,3,6]

$$1 - T_A \approx \frac{3A}{R_{ch}^2} A_3\mathbb{P} \log\left(\frac{x_A}{x}\right) \sim 1.5 \cdot 10^{-2} A^{1/3} \log\left(\frac{x_A}{x}\right). \quad (15)$$

We used $A_3\mathbb{P} \approx 0.16 \text{ GeV}^{-2}$ as determined in the photoproduction experiment [7]. Here, the crucial point is that the triple pomeron coupling becomes approximately independent of Q^2 and quark masses

[2,3,8]. Therefore we may reinterpret the nuclear attenuation of $Q\bar{Q}g$ Fock states as nuclear modifications of target gluon distributions. These are then unique for all hard processes. However in DIS the non-factorizable contributions of light flavor $q\bar{q}$ pairs dominate nuclear shadowing down to $x \lesssim 10^{-3}$ [6]. For the reason that in the $Q\bar{Q}g$ admixture in the physical gluon, the size of the $Q\bar{Q}$ pair is much smaller than the separation of the gluon from the $Q\bar{Q}$ pair, eq. (15) shall be applicable also to the heavy-flavor hadroproduction.

Weak nuclear attenuation in inclusive hadroproduction of open-charm does not imply weak nuclear effects for exclusive charmonium channels. As far as nuclear effects are concerned, the difference between both should be quite similar to the well established difference between inclusive photoproduction of open-charm and exclusive photoproduction of charmonium bound states [8]. A detailed treatment of exclusive charmonium production will however be presented elsewhere.

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The spectrum and solutions of the generalized BFKL equation for total cross section

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The asymptotic behavior of high-energy scattering in perturbative QCD is usually discussed in terms of the Balitskii-Fadin-Kuraev-Lipatov (BFKL) equation for the differential gluon distribution function [1,2]. The work [1,2] focused on a somewhat academic case of the fixed strong coupling and infinite range of colour forces. On the other hand, the indispensable feature of the realistic QCD is the running QCD coupling and finite range of color forces. In this communication [3] we present a derivation of the generalized BFKL equation, which incorporates the running strong coupling. Also, we introduce a finite correlation radius for the perturbative gluons in a gauge invariant manner. We report on an analysis of solutions of this generalized BFKL equation and demonstrate that it has a continuous spectrum. Consequently, the generalized BFKL pomeron is described by the cut in the complex j -plane, and introduction of the running QCD coupling does not lead to a sequence of poles as it has been claimed by Lipatov [2].

The principal idea behind the derivation of the generalized BFKL equation is that the growth of the total cross section is related to a familiar growth with energy of the multiplicity of the Weizsäcker-Williams (WW) gluons in interacting hadrons [4]. The perturbative $q\bar{q}g$ Fock state generated radiatively from the parent colour singlet $q\bar{q}$ state of size $\vec{r}$ has the interaction cross section

$$\sigma_3(r, \rho_1, \rho_2) = \frac{9}{8}[\sigma_0(\rho_1) + \sigma_0(\rho_2)] - \frac{1}{8}\sigma_0(r), \quad (1)$$

where $\vec{\rho}_{1,2}$ are separations of the gluon from the quark and antiquark respectively, $\vec{\rho}_2 = \vec{\rho}_1 - \vec{r}$ and $\sigma_0(r)$ is the dipole cross section for the $q\bar{q}$ dipole $\vec{r}$. The increase of the cross section for the presence of gluons equals

$$\begin{aligned} \Delta\sigma_g(r, \rho_1, \rho_2) &= \sigma_3(r, \rho_1, \rho_2) - \sigma_0(r) \\ &= \frac{9}{8}[\sigma_0(\rho_1) + \sigma_0(\rho_2) - \sigma_0(r)]. \end{aligned} \quad (2)$$

The lightcone density of soft, $z_g \ll 1$, gluons in the $q\bar{q}g$ state derived in [4] equals

$$\begin{aligned} |\Phi_1(\vec{r}, \vec{\rho}_1, \vec{\rho}_2, z_g)|^2 &= \frac{1}{z_g} \frac{1}{3\pi^3} \mu_G^2 \\ &\times \left| g(R_1)K_1(\mu_G \rho_1) \frac{\vec{\rho}_1}{\rho_1} - g(R_2)K_1(\mu_G \rho_2) \frac{\vec{\rho}_2}{\rho_2} \right|^2. \end{aligned} \quad (3)$$

Here $g(r)$ is the effective colour charge, $\alpha_S(r) = g_S(r)^2/4\pi$ is the running strong coupling, $R_{1,2} = \min\{r, \rho_{1,2}\}$, $K_1(x)$ is the modified Bessel function, z_g is a fraction of the (lightcone) momentum of $q\bar{q}$ pair carried by the gluon, and $\int dz_g/z_g = \log(s/s_0) = \xi$ gives the usual logarithmic multiplicity of radiative gluons, which is at the heart of the LLXA. The wave function (3) counts only the physical, transverse, gluons. In (3), $\mu_G K_1(\mu_G \rho) \vec{\rho}/\rho$

emerges as the gradient $\partial K_0(\mu_G \rho)/\partial \vec{\rho}$, and $K_0(\mu_G \rho)$ is precisely the 2D Coulomb-Yukawa screened potential, which in a more generic case must be understood as the correlation function for the perturbative gluons. Eq. (3) describes the spatial density of WW gluons generated by two oppositely charge sources separated by distance $\vec{r}$. With allowance for the perturbative $q\bar{q}g$ Fock component in the beam dipole A , the dipole cross section takes the form

$$\begin{aligned} \sigma(\xi, r) &= \sigma_0(r) + \\ &\int dz_g \vec{\rho}_1 |\Phi_1(\vec{r}, \vec{\rho}_1, \vec{\rho}_2, z_g)|^2 \Delta\sigma_g(r, \rho_1, \rho_2) \\ &= \sigma_0(r) + \sigma_1(r)\xi = [1 + \xi \mathcal{K} \otimes] \sigma_0(r), \end{aligned} \quad (4)$$

where the kernel $\mathcal{K}$ is defined by

$$\begin{aligned} \sigma_1(r) &= \mathcal{K} \otimes \sigma_0(r) = \frac{3}{8\pi^3} \int d^2 \vec{\rho}_1 \mu_G^2 \\ &\times \left| g(R_1)K_1(\mu_G \rho_1) \frac{\vec{\rho}_1}{\rho_1} - g(R_2)K_1(\mu_G \rho_2) \frac{\vec{\rho}_2}{\rho_2} \right|^2 \\ &\times [\sigma_0(\rho_1) + \sigma_0(\rho_2) - \sigma_0(r)]. \end{aligned} \quad (5)$$

Eq. (4) shows that the effect of gluons can be reabsorbed into the generalized dipole cross section $\sigma(\xi, r)$. To higher orders in ξ , this generalized dipole cross section can be expanded as $\sigma(\xi, r) = \sum_{n=0}^{\infty} \frac{1}{n!} \sigma_n(r) \xi^n$. Then, Eq. (5) generalizes to $\sigma_{n+1} = \mathcal{K} \otimes \sigma_n$, and

$$\frac{\partial \sigma(\xi, r)}{\partial \xi} = \mathcal{K} \otimes \sigma(\xi, r) \quad (6)$$

is our generalized BFKL equation for the dipole cross section.

In the BFKL scaling limit of $r, \rho_1, \rho_2 \ll R_c$ and fixed α_S

$$\mu_G^2 \left| K_1(\mu_G \rho_1) \frac{\vec{\rho}_1}{\rho_1} - K_1(\mu_G \rho_2) \frac{\vec{\rho}_2}{\rho_2} \right|^2 = \frac{r^2}{\rho_1^2 \rho_2^2}, \quad (7)$$

the kernel $\mathcal{K}$ becomes independent of the gluon correlation radius R_c and with the fixed α_S it takes on the scale-invariant form (hereafter we refer to this limiting case as the BFKL approximation). In the BFKL approximation the eigenfunctions of Eq. (6) are $E(\omega, r) \exp[\Delta(\omega)\xi]$, where

$$E(\omega, r) = (r^2)^{\frac{1}{2} + i\omega} \quad (8)$$

with the eigenvalue (intercept)

$$\begin{aligned} \Delta(\omega) &= \\ &\frac{3\alpha_S}{\pi} [2\Psi(1) - \Psi(\frac{1}{2} - \omega) - \Psi(\frac{1}{2} + \omega)], \end{aligned} \quad (9)$$

where $\Psi(x)$ is the digamma function. The final result for $\Delta(\omega)$ coincides with eigenvalues of the BFKL equation [1,2].

Both the correlation radius R_c and the running/freezing coupling break the scale invariance of

the BFKL approximation. What is the impact of the broken scale invariance on the spectrum and the complex- j plane properties of the generalized BFKL pomeron?

We start with simple limiting cases and consider first the case of finite R_c and fixed coupling α_S . Evidently, in this case, on the infinite semiaxis $r \ll R_c$, our kernel is identical to the kernel of the BFKL approximation, the spectrum of eigenvalues will be continuous, the cut in the j -plane and the intercept $\Delta_{\mathbb{P}}$ will be the same as in the BFKL approximation of $R_c \rightarrow \infty$. The second limiting case is of the running/freezing coupling $\alpha_S(r)$ at $R_c \rightarrow \infty$. In this case, at $r > R_f$, we recover the BFKL approximation with $\alpha_S = \alpha_S^{(f)}$, the spectrum of eigenvalues will evidently be continuous, and the intercept $\Delta_{\mathbb{P}}$ and the cut in the complex j -plane will be the same as in the BFKL approximation. The realistic case of the interest is a combination of the two previous cases. A somewhat tedious analysis shows that in this case too we can exclude the node-free localized solutions with $\sigma(r)$ vanishing at large r and, consequently, the spectrum of our generalized BFKL equation remains the continuous one. This shows that, within the considered infrared regularization, our generalized BFKL equation produces a cut in the complex j -plane. The rightmost singularity in the j -plane will have the dipole cross section $\sigma_{\mathbb{P}}(r) = \text{const}$ vs. r at large r . Its intercept will be controlled by the behavior of $\sigma_{\mathbb{P}}(r)$ at $r \sim R_c$.

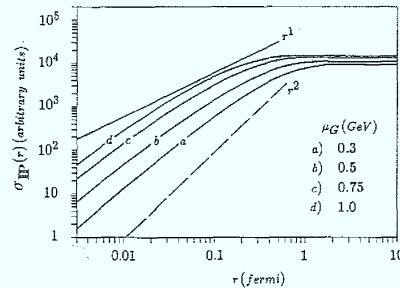


Fig. 1: The pomeron dipole cross section $\sigma_{\mathbb{P}}(r)$ for different values of μ_G . The straight lines show the r^1 and r^2 behavior.

In order to determine the intercept of the generalized BFKL pomeron, we study the large- ξ behavior of numerical solutions of Eq. (6) and verify that $\sigma(\xi, r)$ has the same asymptotic behavior $\sigma(\xi, r) \Rightarrow \sigma_{\mathbb{P}}(r) \exp(\Delta_{\mathbb{P}} \xi)$ irrespective of the boundary condition at $\xi = 0$. Namely, we compute $\Delta_{eff}(\xi, r) = \partial \log \sigma(\xi, r) / \partial \xi$ and check that at large ξ the effective intercept $\Delta_{eff}(\xi, r)$ tends to the same limiting value $\Delta_{\mathbb{P}}$ for all r . In Fig.1 we present the eigenfunction $\sigma_{\mathbb{P}}(r)$ for few values of μ_G (we keep $\alpha_S^{(f)} = 0.8$). As we have anticipated above, $\sigma_{\mathbb{P}}(r)$ flattens at large r and decreases towards small r faster than the BFKL

solution $\propto r^1$ of Eq. (8), but slower than $\propto r^2$.

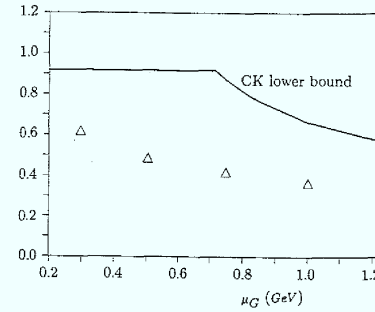


Fig. 2: The intercept $\Delta_{\mathbb{P}}$ for $\mu_G = 0.3, 0.5, 0.75, 1.0$ GeV. The solid curve shows the Collins-Kwiecinski lower bound.

The corresponding intercepts are shown in Fig.2. Our $\Delta_{\mathbb{P}}$ is significantly smaller than the BFKL result (9) and smaller than the Collins-Kwiecinski (CK) lower bound [5] $\Delta_{\mathbb{P}} > 3.6 \alpha_S^{(f)}/\pi$, which has been derived using the sharp infrared cutoff of $k^2 > k_0^2$. We notice, that the CK cutoff breaks the initial symmetry of the BFKL equation, does not give an identical cutoff of all gluon propagators and is suspect from the point of view of gauge invariance [6,7].

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Nucleon-nucleon bremsstrahlung is a process, where two strongly interacting nucleons emit a photon ($N_1 + N_2 \rightarrow N_1 + N_2 + \gamma$). Since the nucleon in the intermediate state is necessarily off its mass-shell, this reaction has been considered as an ideal tool for investigating the off-shell behaviour of the nucleon-nucleon (NN) interaction. With the recent developments, the theory of NN bremsstrahlung reaction has reached a considerable level of sophistication. Indeed, nowadays, most of the NN bremsstrahlung calculations include simultaneously various terms regarded as higher-order corrections: relativistic spin correction, rescattering or double-scattering, and Coulomb correction [1]. Also, quite recently, the influence of the Δ -isobar degree-of-freedom in NN bremsstrahlung has been investigated within a coupled-channel approach [2]. However, despite these developments, virtually all existing NN bremsstrahlung calculations take into account, beyond the soft-photon-approximation, only a part of the total electromagnetic transition current, the so-called one-body current $\vec{j}^1B$, while the gauge invariant constraints demand the existence of a second part, the two-body current $\vec{j}^2B$. In the neutron-proton bremsstrahlung reaction, this latter current (in the soft-photon-approximation) is known to dominate over the one-body current for energetic photon productions, while for the proton-proton bremsstrahlung, its contribution vanishes identically. To best investigate the off-shell behaviour of the NN interaction using the NN bremsstrahlung reaction, one requires kinematic conditions in which high energy photons are produced. Therefore, it may be important to take into account the two-body current contribution even for the proton-proton bremsstrahlung reaction. It is the purpose of the present work to investigate the influence to the two-body current on proton-proton bremsstrahlung observables.

The effective two-body current $\vec{j}^2B$ should obey the gauge condition, which is given (under certain restrictions) by

$$\vec{k} \cdot (\vec{j}^1B + \vec{j}^2B) = -[(K + V), \rho^1B], \quad (1)$$

where $\vec{k}$ is the photon momentum, K the kinetic energy operator of the two nucleons, V the Nucleon-Nucleon-Potential, and ρ^1B is the one-body charge density. The two-body current can be derived uniquely from (1) only in the soft-photon-approximation. We, therefore, follow the feynman diagram prescription to obtain a two-body current that fulfills the gauge condition (1) in all orders in photon energy. We consider a *complete set* of feynman diagrams that contributes to the bremsstrahlung amplitude and that is not already included in the one-body current. By complete we mean, that the two-body current, derived from this set, in addition to the one-body current, fulfills eq. (1). This choice of feynman diagrams is not unique, since one can always imagine diagrams, which are gauge invariant by themselves, and therefore can be added to the two-body current without violating eq. (1). An example for these kind of diagrams are the meson decay diagrams, which have been discussed in ref.

[3, 4].

To investigate the influence of the two-body current to the proton-proton bremsstrahlung reaction, we restrict ourselves to contributions of second order in the strong coupling constant and of order $\frac{p^2}{m^2}$, where $\vec{p}(m)$ is the nucleon momentum (mass). In this limit one can easily check the condition (1) explicitly.

For static potentials such as the OBEPQ [5], the set of pair diagrams (fig. 1) forms a complete set of diagrams in

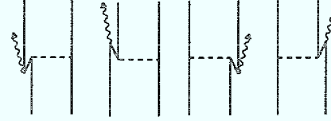


Fig. 1: Pair diagrams.

the sense mentioned above. For the non-static potentials such as the time ordered potential OBEPQ [5], in addition to the set shown in Fig. 1, we include the recoil diagrams (fig.2). These type of diagrams are a part of the one-body current that were neglected so far. It turns out, that they are of minor importance.

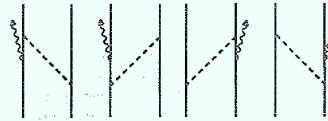


Fig. 2: Recoil diagrams.

Other diagrams, like contact diagrams for the ρ -meson or mesonic diagrams are not present for the $pp\gamma$ -reaction in the considered limit.

These new contributions, derived as described above, allow us to investigate effects of the two-body current on the $pp\gamma$ -reaction. At incident energies below pion threshold ($T_{Lab} = 280 MeV$) we compared the predictions for various bremsstrahlung observables after including the two-body current based on the OBEPQ and OBEPQ potentials. Since the parameters used to calculate these new contributions, such as coupling constants, meson masses, and cut-off parameters for the hadronic formfactors, are in general different for different potentials, effects of the two-body current on the bremsstrahlung observables depend on the potential used. Therefore not only the different off-shell behaviour of the potential, but also the two-body current must be considered while comparing bremsstrahlung observables for different NN-potentials. With increasing incident energy, the influence of the two-body current becomes more important. Fig. 3 shows the cross section for $T_{Lab} = 400 MeV$ and $\theta_1 = \theta_2 = 30^\circ$.

Here the two-body current clearly enhances the cross section at forward and backward photon angles, while around $\theta = 90^\circ$ it reduces the cross section. (at $T_{Lab} = 280 MeV$ the effect of the two-body current on the cross section is negligible). The momentum carried by the exchanged

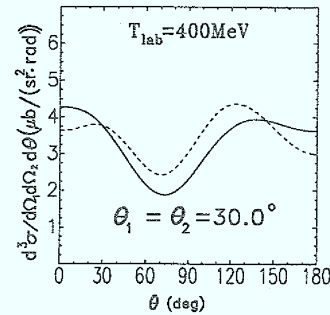


Fig. 3: Cross section at $T_{Lab} = 400 MeV$ at large nucleon scattering angles with (solid line) and without (dashed line) two-body current.

mesons $\vec{q} = \vec{p} - \vec{p}'$, where $\vec{p}(\vec{p}')$ are the relative momenta of the nucleons before (after) the scattering, is responsible for this rapid increase of the effect of the two-body current on the cross section rather than the increase of the photon momentum. Due to this, one also observes that the effect of the two-body current increases as the proton scattering angles increase.

In conclusion, the present qualitative study indicates that the influence of the two-body current on proton-proton bremsstrahlung observables should not be neglected. In particular, it is important for a valid comparison of the predictions of different nucleon-nucleon potentials, especially for spin observables. At larger incident energies, the effects of the two-body current on the cross section become generally important. Therefore for a proper description of the $pp\gamma$ -reactions above pion threshold, such as the experiment planned at the COSY facility, not only the contributions from subnucleonic degrees-of-freedom (Δ -isobar) is important, but also the two-body current seems to become of considerable importance.

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Quantum chromodynamics (QCD) is generally held to be the fundamental theory of strongly interacting particles. One of its implications seems to be that quarks and gluons are not observable by themselves but only in colorless compound states. Although theoretical studies of QCD provide much evidence to that effect, at present the precise dynamical mechanism of quark confinement from first principles is still missing [1]. In the absence of such a direct derivation from the QCD Lagrangian, a number of different model approaches have been considered in recent years to achieve confinement. Conceptually arguably one of the most elegant mechanisms to this end is based on the idea that the quark propagator should not have a pole, a mechanism which was proposed already twenty years ago [2]. (For recent implementations of this approach, see [3], and references therein.) We present here a model of quark confinement along similar lines, with a quark propagator without any singularities for real energies.

This is achieved by making the quark mass a dynamical quantity whose energy dependence is such that the quark cannot go on-shell; it is thus confined dynamically. In the usual Feynman-type covariant formulation, this immediately poses problems as to the existence of spinor solutions of the Dirac equation and hence as to the meaningful definition of color current matrix elements. In order to circumvent these difficulties and, furthermore, avoid possible problems with causality, we use here the time-ordered scattering formalism of ref. [4] which is manifestly covariant under an *off-shell* modification $\mathcal{M}$ of the usual Lorentz transformations $\mathcal{L}$. The transformation $(e', \mathbf{p}') = \mathcal{M}(e, \mathbf{p})$ of an arbitrary four-vector $(e, \mathbf{p})$ describing a cluster with three-momentum $\mathbf{p}$ and off-shell energy e is defined by relating the three-momenta $\mathbf{p}'$ and $\mathbf{p}$ in two different frames by Lorentz transformations involving only the respective *on-shell* energies ω' and ω , i.e., $(\omega', \mathbf{p}') = \mathcal{L}(\omega, \mathbf{p})$, where

$$\omega = \omega(m, \mathbf{p}) = \sqrt{m^2 + \mathbf{p}^2}, \quad (1)$$

etc., with m being the mass. The relation between off-shell energies e' and e is then defined by

$$e - \omega(m, \mathbf{p}) = e' - \omega(m, \mathbf{p}'), \quad (2)$$

similar to Galilei transformations in Euclidean space. The modified transformations $\mathcal{M}$, therefore, leave the difference between on- and off-shell energies invariant, and reduce to $\mathcal{L}$ when going on-shell. We note that while $\mathcal{M}$ appears as a nonlinear transformation of four-vectors $(e, \mathbf{p})$, it can be understood [4] as a covariant four-dimensional projection of a linear five-dimensional extension of the Lorentz transformation $\mathcal{L}$, i.e., Minkowski space is a hypersurface of an underlying five-dimensional manifold whose extra, fifth, dimension is interpreted as the off-shell

energy variable with transformation properties defined in (2), while the usual Minkowski-space energy is mapped onto the mass shell.

This off-shell extension is essential for ensuring the existence of spinor solutions of the Dirac equation even if the quark mass m is to be a dynamical quantity such that $e = \omega$ is never true. Letting m depend on the invariant quantity $e - \omega$, i.e., $m = m(e - \omega)$, the Dirac equation

$$(\omega \gamma^0 - \mathbf{p} \cdot \boldsymbol{\gamma} - m(e - \omega)) u = 0 \quad (3)$$

is immediately seen to admit the usual solutions, albeit with the mass depending on the off-shell parameter e . Moreover, in view of the transformation properties of $(\omega, \mathbf{p})$ and $e - \omega$ described above, eq. (4) remains covariant.

Defining now the dressed time-ordered quark propagator as having the same form as for an elementary spin-1/2 fermion, i.e.,

$$t(e + i0, \mathbf{p}) = \frac{\mathcal{P}}{e - \omega(m, \mathbf{p}) + i0}, \quad (4)$$

where $\mathcal{P} = \sum_{\lambda} u_{\lambda} \bar{u}_{\lambda}$ is the covariant on-shell spin-1/2 projector, the functional form of the quark mass can be extracted from the time-ordered version of the Dyson-Schwinger equation for the dressed propagator, viz.

$$t = t_b + t_b \Sigma[t] t, \quad (5)$$

with t_b being the bare propagator and the self-energy gluon loop $\Sigma[t]$ being a functional of t .

In the rest frame where $e - \omega \rightarrow E - m$, for a vanishing bare mass, the mass is simply given by $\Sigma[t]$ which, using the rules of ref. [4], is found to be [5]

$$m(E) = \frac{g^2}{3\pi^2} \int_0^\infty dq q \frac{m_q}{\omega_q} \frac{N_q}{E - q - \omega_q}, \quad (6)$$

where g is the quark-gluon coupling constant and $m_q = m(E - q - \omega_q + m_q)$ and $\omega_q = (m_q^2 + q^2)^{1/2}$ is a coupled system of implicit functions. This equation provides a highly nonlinear self-consistency condition for the dynamical quark mass $m(E)$ written here as an explicit function of the rest-frame energy E instead of as an implicit function of the invariant $E - m$. The spinor matrix element N_q is given by

$$N_q = \frac{1}{4} \text{Tr} \left((1 + \gamma^0) \gamma_{\mu} \mathcal{P}_q D^{\mu\nu} j_{\nu} \right) F, \quad (7)$$

where $\mathcal{P}_q$ is the on-shell spin-1/2 projector for the dressed quark with momentum q in the loop, γ_{μ} is the current at the bare quark-gluon vertex and j_{ν} the dressed current operator with F the corresponding dressed form factor. $D^{\mu\nu}$ is the gluon tensor which is taken in the Feynman gauge, that is $D^{\mu\nu} = -g^{\mu\nu}$.

In ref. [5], the dressed vertex in the quark-gluon loop was approximated by a bare vertex, i.e., $j_\nu = \gamma_\nu$ and $F = 1$, and the resulting spinor matrix element N_q was taken in the infrared limit, which provided $N_q = -1$. This was seen to be equivalent to taking into account only purely Coulombic gluons in the loop. The corresponding solution of eq. (6) is indeed found [5] to provide a functional form $m(E)$ for the quark mass without a fixed point, i.e., $E = m(E)$ is never realized and the quark is thus confined dynamically. On the real axis, the functional behavior of $m(E)$ is approximated quite well by [5]

$$m_{\text{approx}}(E) = \frac{E}{2} + \sqrt{\frac{E^2}{4} + m_0^2}, \quad (8)$$

where m_0 provides the mass scale; it can be chosen as unity. The dynamical mass is thus seen to rise monotonically from zero at $-\infty$ through unity at $E = 0$ to infinity at $+\infty$, always staying above the line corresponding to the pole condition $m = E$.

Equation (6) corresponds to a nonlinear eigenvalue condition for the coupling constant. For the approximate case $j_\nu = \gamma_\nu$, $F = 1$, and $N_q = -1$, its value is numerically found to be $g^2/4\pi = 4.712$ which can actually be shown [5] to correspond to

$$\frac{g^2}{4\pi} = \frac{3\pi}{2} \quad (9)$$

if one assumes the vertex to be dominated by soft gluons.

As far as current conservation is concerned, one sees immediately that the bare γ^μ current does not admit a Ward-Takahashi identity (WTI) since the dynamical quark masses before and after the gluon absorption (or emission) in general are different and hence their difference does not vanish. For the spinor matrix element of the dressed current operator j^μ we may now impose the requirement that its divergence provide a Ward-Takahashi identity to actually deduce its functional form (or at least constrain it to a large extent). Making a general ansatz for the current $j^\mu(p, q)$ in terms of invariants, one finds that

$$j^\mu(p, q) = \left(i\sigma^{\mu\nu} + \beta \frac{p^\mu q^\nu - q^\mu p^\nu}{m_p m_q} \right) \frac{(p - q)_\nu}{m_p + m_q}, \quad (10)$$

where the masses m_p and m_q pertain to the quarks with momenta p and q , respectively. The open parameter β cannot be determined by the requirement of the WTI (see below). Let us emphasize here that for the derivation of eq. (10), the existence of spinor solutions for eq. (3) is crucial. Without an appropriate off-shell formalism, this would not have been possible.

Associated with the current (10) one furthermore finds a dressed vertex F given by

$$F_\lambda(\eta^2) = \frac{\lambda^2}{\eta^2}, \quad (11)$$

where λ provides some scale and η^2 is an invariant [4],

$$\eta^2 = \frac{(p \cdot q)^2 - m_p^2 m_q^2}{m_p^2 m_q^2}, \quad (12)$$

which essentially corresponds to the relative velocity between the quarks before and after the absorption (emission) of the gluon. Although this form factor appears to be infrared divergent, it actually is not since the scale λ itself turns out to be η dependent, viz.

$$\lambda = \Lambda \frac{\eta^2}{\lambda_0^2 + \eta^2}, \quad (13)$$

where Λ and λ_0 are parameters. [This form for λ is a convenient parametrization; a more general form would be $\lambda = \Lambda \eta^2 f$ with an open function f . It is essential, however, that λ vanishes like η^2 .] There exists now a very delicate interplay between the coupling "constant" g , the form factor F , and the spin matrix element N_q of eq. (7). One finds that the coupling scales with λ , i.e.,

$$\frac{g^2}{4\pi} = \frac{3\pi}{\beta \lambda^2} = \frac{1}{\lambda^2} \frac{g_0^2}{4\pi}; \quad (14)$$

in other words, the λ dependence cancels altogether between (11) and (14), and eq. (6) does not depend on the scale λ at all. Furthermore, the $1/\eta^2$ singularity from (11) is cancelled by a zero of the remaining factor in the spin matrix element (7), and everything remains finite and well defined. This allows one now to fix the parameter β by requiring that the rescaled color charge of the dressed vertex for vanishing relative velocity between the quarks be equal to g_0 [cf. eq. (14)]. One then has $\beta = 5/2$, and the current (10) is now completely determined.

Although somewhat different in detail, the resulting solution of eq. (6) again exhibits the generic confining behavior found in the previous approximate case. Preliminary results from calculations of meson masses generated by one-gluon exchange between fully dressed quark-gluon vertices described by the current (10) and the form factor (11) are very encouraging. In particular, it is found that the pion mass comes out to be zero, in accordance with chiral symmetry.

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It is still an unsolved problem, to derive hadronic properties at intermediate energies directly from QCD. Therefore suitable models are required, not only containing the features of QCD considered to be important, but also allowing to perform practical calculations. It is especially of high interest, whether it is possible to describe both baryons and mesons with a common set of parameters.

Recently relativistic Faddeev quark-models have been developed for the baryon octet [1, 2, 3]. We extended the Faddeev approach to the baryon decuplet as well [5]. This was possible via the inclusion of the axial vector diquark in addition to the scalar one as it was already performed by Ishii et al. [4] for the nucleon.

Quark models of the nucleon have to respect chiral symmetry. The Nambu–Jona-Lasinio (NJL) model [6] provides a particularly simple implementation of chiral symmetry as it is reviewed in refs. [7]. As most relativistic Faddeev models we therefore start from a NJL-type Lagrangian:

$$\begin{aligned} \mathcal{L} = & \bar{q}(i\partial - m_0)q \\ & - G_1(\bar{q}\tilde{A}_1^{\mu k}q)(\bar{q}\tilde{A}_{\mu 1}^kq) \\ & + G_2(\bar{q}\tilde{A}_2^{\mu k}q)(\bar{q}\tilde{A}_{\mu 2}^kq), \end{aligned} \quad (1)$$

where $m_0 = \text{diag}(m_0^u, m_0^d, m_0^s)$, $\tilde{A}_1^\mu = \gamma^\mu \otimes I_f \otimes \lambda_c^k$ and $\tilde{A}_2^\mu = \gamma^\mu \gamma_5 \otimes I_f \otimes \lambda_c^k$, λ_c^k being the Gell-Mann matrices.

Using the Fierz transformation the interaction can be expressed in terms of a particle–antiparticle interaction with suitable colour channels

$$\begin{aligned} \mathcal{L}_f^M = & \frac{g_1^2}{2}(\bar{q}\Lambda^{[1]}q)(\bar{q}\Lambda^{[1]}q) \\ & + (\text{colour} - \text{octet} - \text{terms}), \end{aligned} \quad (2)$$

where $\Lambda^{[1]} = \Gamma_c^a \otimes \lambda_f^k \otimes \Gamma^a$ and a particle–particle interaction

$$\begin{aligned} \mathcal{L}_f^D = & -\frac{g_2^2}{2}(\bar{q}\Lambda^{[3]}C\bar{q}^T)(q^TC^{-1}\Lambda^{[3]}q) \\ & + (\text{colour} - \text{sextet} - \text{terms}), \end{aligned} \quad (4)$$

where $\Lambda^{[3]} = t_f^i \otimes \lambda_f^k \otimes \Gamma^a$, with $\Gamma_s^a \in \{I, i\gamma_5\}$ and $\Gamma_v^a \in \{i\gamma^\mu, i\gamma^\mu\gamma_5\}$. The color antitriplet is generated by $t_{\alpha\beta}^i = i\sqrt{\frac{3}{2}}\epsilon_{\alpha\beta\gamma}$. The generators of the flavor group $U(3)_F$ are denoted λ_A^i by $A \in \{0, \dots, 8\}$. The coupling constants are $g_1^2 = \frac{4}{9}(G_1 + G_2)$, $g_2^2 = \frac{1}{2}g_1^2$ and $g_1^2 = \frac{1}{2}g_2^2$.

We take into account only the two low lying diquarks namely the scalar one ($\Gamma_s = i\gamma_5$) and the axial vector one ($\Gamma_v^\mu = i\gamma^\mu$).

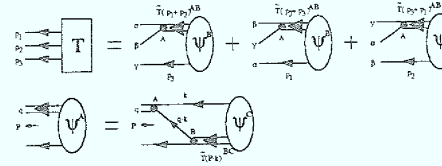


Figure 1: The Faddeev equation. Greek letters denote the internal degrees of freedom of a single quark, A, B stand for the type of the diquark and $\bar{T}$ stands for the diquark propagator.

A detailed derivation of the relativistic Faddeev equations is given in ref. [1], the Feynman graphs are shown in figure 1.

One has to use the same kind of approximation as has already been used in the meson case. The essential one is to neglect the three-body irreducible graphs. This reduces the three-body problem to an effective two-body Bethe–Salpeter equation with an energy dependent interaction. To solve the equations we used the so called static approximation suggested in ref. [2] and justified in refs. [3, 4] via the comparison of the approximate results with the full calculation.

The model contains 5 parameters: in addition to the 2 current masses ($m_0^u = m_0^d, m_0^s$) and the 2 coupling constants (g_1^2, g_2^2) we need a cut off (Λ) to get finite results. On the other hand we have 14 observables:

- 4 masses from the baryon decuplet ($M_\Delta, M_{\Sigma^*}, M_{\Xi^*}, M_{\Omega^*}$),
- 4 masses from the baryon octet ($M_N, M_\Sigma, M_\Lambda, M_\Xi$),
- 3 masses from the nonet of the vector mesons ($m_\rho = m_\omega, m_{K^*}, m_\phi$), where the ϕ is treated as pure $s\bar{s}$ state,
- 2 masses from the nonet of the pseudoscalar mesons (m_π, m_K) (To calculate the mass of the η' -meson one would have to include the t'Hooft interaction) and
- the pion decay constant f_π using the Gell-Mann–Oaks–Renner relation.

For a given Cut-off we can therefore fix all 4 parameters using only mesonic data in contrast to previous work, where there was at least one purely baryonic parameter. We investigate the behavior of the spectrum in dependence of the constituent mass of the up quark. The cut off is determined by the NJL gap eqn. [6]. Note that we have the constraint of $m_u \geq 411$ MeV in order to bind the $\Delta_{33}(1232)$.

The deviation of the masses of decuplet and octet baryons calculated in our model from the experimental values is shown in figure 2 as a function of the quark mass. The masses of both nucleon and

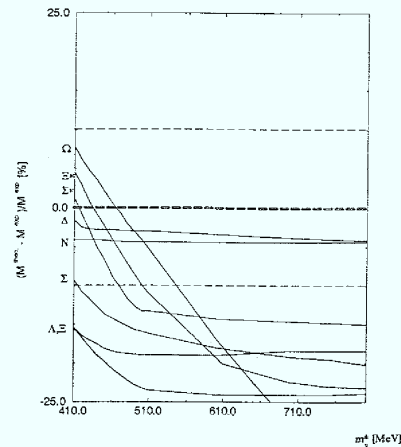


Figure 2: The deviation of the numerical results from the experimental data in percent.

$\Delta_{33}(1232)$ are reproduced within 5 % for quark masses between 411 MeV and 800 MeV. With increasing quark masses the coupling constants are adjusted so that the meson masses remain constant. This suffices to produce both a nucleon and a delta mass which show only weak dependence on the constituent quark mass. For a constituent quark mass of 411 MeV, the masses of the strange baryons deviate from the experimental ones by approximately 15 % and one can not improve the agreement with the data by increasing the quark masses. We checked that treating the scaling factor $\alpha = \frac{2}{3}$ as a free parameter reduces the maximal deviation from the experimental masses to 7 %. The fact that the strange baryons deviate more strongly from the data than the proton and the delta may be taken as a hint that the 't'Hooft interaction should not be neglected. The 't'Hooft interaction affects the scalar diquarks only and therefore only the octet baryons.

To summarize, we have shown that the Faddeev approach to the baryonic sector of the Nambu-Jona-Lasinio model is able to provide a consistent description of both octet and decuplet baryons.

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Quasielastic ($e, e'p$) scattering on nuclei at large missing momentum p_m is a relevant tool for investigation of short-distance nucleon-nucleon interaction (initial state two-nucleon correlations (ISC)) in the nuclear medium and constitutes an important part of the CEBAF experimental program [1]. The measured p_m distribution is distorted by final state interaction (FSI) of the struck proton in the target nucleus debris. In the CEBAF range of energies, the radius b_o of the FSI is very close to the radius r_c of ISC. This makes FSI an important background in the production of protons with large p_m , which may obscure the relationship between the observed p_m distribution and the ISC.

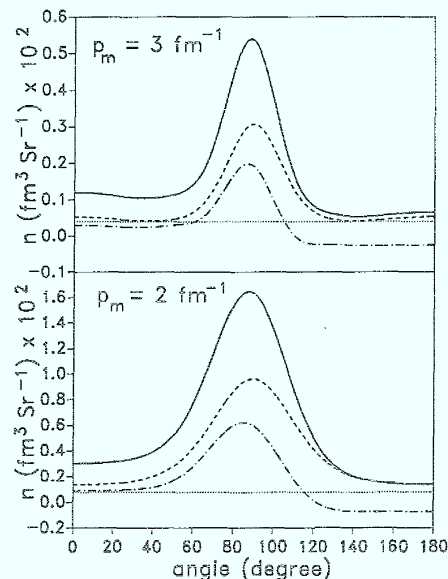


Fig. 1: The angular dependence of the three most important components of $w(\vec{p}_m)$ at $p_m = 2 \text{ fm}^{-1}$ (lower panel) and at $p_m = 3 \text{ fm}^{-1}$ (upper panel). The dotted line is the undistorted distribution $n_F(p_m)$ taken from [4], the dash-dotted line represents the ISC-FSI interference component $w(CT' + C'T; \vec{p}_m)$, the dashed line shows the elastic rescattering component $w(\Gamma'T; \vec{p}_m)$, and the solid line is the sum of the three contributions.

We have presented simple estimates of the effect of FSI in $^4\text{He}(e, e'p)$ scattering[2]. We show how FSI leads to an anisotropic p_m distribution, with

the dominance of FSI effects in transverse kinematics and significant FSI corrections to ISC effects in parallel kinematics. This last point can be important for the y -scaling analysis. We find a novel effect of quantum-mechanical interference of ISC and FSI which gives a large correction to the elastic rescattering of the struck proton and which must be included when discussing more sophisticated features of FSI such as color transparency effects. We concentrate on the region of 4-momentum transfer squared $Q^2 \gtrsim (1-2) \text{ GeV}^2$, which is relevant to the planned CEBAF experiments [1]. Furthermore, in this range of Q^2 , FSI can be described by Glauber's multiple scattering theory [3], which simplifies the evaluation of FSI effects.

In Fig. 1 we show the angular distribution $w(\vec{p}_m)$ for two different values of the missing momentum p_m and its most important components at large p_m . The large- p_m tail of the undistorted distribution $n_F(p_m)$ is borrowed from ref. [4]. For the elastic rescattering and the ISC-FSI interference we take our results. We find a prominent signal of the elastic rescattering and ISC-FSI interference in parallel kinematics. Notice a significant forward-backward asymmetry of the ISC-FSI interference component, which is generated by the real part of the pN scattering amplitude, and also the unexpectedly large contribution from the elastic rescattering term at $p_\perp = 0$, which comes from the θ -function generated tail of the p_z distributions.

To illustrate this general FSI effect, in Fig. 2 we show what happens if one substitutes $\frac{1}{4}$ for the product of $\theta(z_i - z_4)$ and $\theta(z_i - z_4')$. The distortion effect becomes stronger at larger p_m , leading to the forward and backward peaks alongside the FSI peak in the transverse kinematics. This could influence the y -scaling analysis. The sensitivity of these forward and backward peaks to modifications of the conventional Glauber eikonal Green function to account for the finite size of nucleons are discussed elsewhere [5]. Summarizing the main results, we have presented a simple classification of FSI effects compared to the ISC effects. The anisotropic angular dependence of the elastic rescattering and ISC-FSI interference effects compared to the isotropic momentum distribution in the PWIA is a striking signature of FSI effects.

An important finding is that apart from the peak at 90° , distortions by FSI lead also to forward and backward peaks, which slowly build up with increasing missing momentum p_m . We found a large contribution to the p_m distribution in transverse kinematics from the novel effect of the quantum mechanical interference of FSI and ISC effects. The considered aspects of ISC and FSI were not discussed in previous work on FSI and ISC effects in inclusive $A(e, e'p)$ scattering.

Our results suggest that the experimental separa-

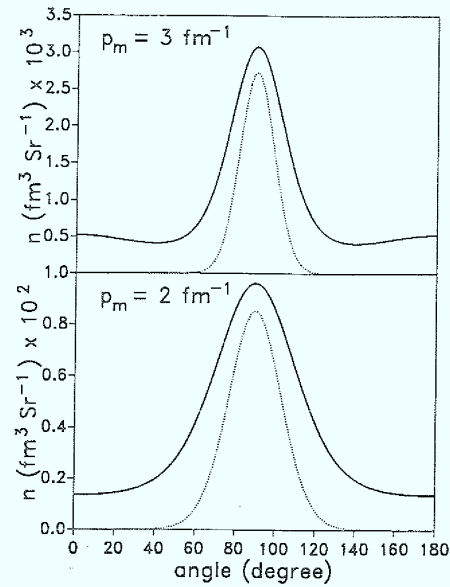


Fig. 2: The angular dependence of the elastic rescattering term $w(\Gamma' T; \vec{p}_m)$ including the product $\theta(z_i - z_4)$ and $\theta(z'_i - z_4)$ (solid line) and replacing the product of the θ -functions by $\frac{1}{4}$ (dotted line) for $p_m = 2 fm^{-1}$ (lower panel) and $p_m = 3 fm^{-1}$ (upper panel).

tion of the FSI and ISC effects is more difficult than thought before. The expansion parameter of the problem $\sim r_c/R_0$ is not very small, and more detailed numerical analysis with allowance for the tensor correlation function and the D-wave in 4He are called upon to establish the sensitivity to models for the correlation function.

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There is a basic necessity of the knowledge of the nucleon momentum distribution $n(\vec{p})$ in nuclei, in particular at large momenta, for it provides unique information on short-distance nucleon-nucleon interaction (initial-state N - N correlations ISC) in the nuclear medium. A quantitative understanding of the final-state interaction (FSI) effects is crucial for disentangling $n(\vec{p})$ from the missing momentum distribution $W(\vec{p}_m)$ as measured in inclusive $A(e, e'p)$ scattering in quasielastic kinematics. In a paper [1] on $^4\text{He}(e, e'p)$ scattering we found that at large missing momenta $\vec{p}_m$, which are usually believed to come from the ISC effects, the FSI effects are large and overwhelm the ISC effects. In [1] the FSI was shown to produce a strong angular anisotropy of $W(\vec{p}_m)$, which at large p_m develops a sharp peak in transverse kinematics, at $\theta \sim 90^\circ$, two minor peaks at $\theta = 0^\circ$ and 180° in parallel kinematics and a sizable backward-forward asymmetry (here θ is the angle between the missing momentum and the (e, e') momentum transfer $\vec{q}$).

To generalize these results and get a deeper insight into them, a systematic study of the $D(e, e'p)n$ unpolarized scattering was performed [2]. Prediction of FSI effects in this process is of great interest on its own sake, as experiments with the deuterium target constitute an important part of the experimental program at CEBAF [3]. Other motivations are that (i) in the region of dominance of ISC effects, the momentum distribution for heavy nuclei are expected to closely resemble that in the deuteron [4]; (ii) the realistic models of the deuteron [5] allow an accurate evaluation of FSI effects in the longitudinal missing momentum distribution which is important for the y -scaling analysis [6]; (iii) one can assess differences between FSI effects for the S and D waves in the initial N - N pair. Despite the deuteron being a dilute target, we find large FSI effects at $p_m \gtrsim 1.5 \text{ fm}^{-1}$. At the large $Q^2 \gtrsim (1.2) \text{ GeV}^2$ of the interest in the CEBAF experiments, the kinetic energy of the struck proton $T_{kin} \approx Q^2/2m_p$ is high and FSI can be described by Glauber theory [7]. In our analysis of $D(e, e'p)$ scattering in the few-GeV region we have found that FSI effects completely dominate the observed missing momentum distribution $W(\vec{p}_m)$ at high p_m (over $\sim 1.5 \text{ fm}^{-1}$) (see Fig.1). They produce a marked angular anisotropy in $W(\vec{p}_m)$, consisting in a dip at $\theta \sim 90^\circ$ (in transverse kinematics) for $p_m \sim (1.5) \text{ fm}^{-1}$, and a prominent peak at the same angle for larger p_m (see Fig.2). This anisotropy is mostly due to the FSI distortions of the S -wave contribution. The anisotropy effects in the D -wave contribution are also sizable, but much smaller than in the S -wave contribution. The D -wave contribution remains the dominant component for $p_m \gtrsim 1.5 \text{ fm}^{-1}$ in parallel kinematics. One of the points we addressed in more detail are forward and backward peaks in

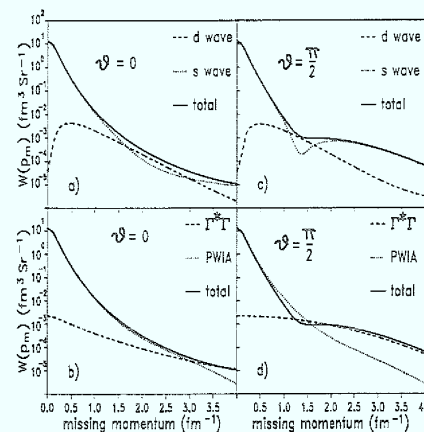


Fig. 1: The missing momentum distribution $W(\vec{p}_m)$ vs. p_m and its decomposition into (a,c) the S and D -wave contributions and (b,d) into the PWIA and the rescattering contributions $\propto \Gamma^* \Gamma$ at (a,b) $\theta = 0^\circ$ and (c,d) $\theta = 90^\circ$. In panels (b,d) the full distribution $W(\vec{p}_m)$ (the solid curve) includes also the contribution $\propto (\Gamma^* + \Gamma)$ from the interference of the PWIA and FSI amplitudes, which is not shown separately.

$W(\vec{p}_m)$, which have their origin in the idealized step-function factor $\theta(z)$ in the eikonal Green function derived [7] under an implicit assumption of pointlike nucleons (z is the longitudinal separation of nucleons). We analyzed the sensitivity of the forward and backward peaks to the smearing of the θ -function to account for the finite size of nucleons and conclude that predictions at $p_m \lesssim 3 \text{ fm}^{-1}$ are free of uncertainties with the smearing. Another interesting quantity is the $p_\perp$ -integrated, longitudinal missing momentum distribution $F(p_z) \equiv \int d^2 \vec{p}_\perp W(\vec{p}_\perp, p_z)$. In the PWIA, $F_{PWIA}(p_z) = \int d^2 \vec{p}_\perp n(\vec{p}_\perp, p_z)$ is related to the asymptotic y -scaling function [6]. In Fig. 3 we show the ratio $R(p_z) = F(p_z)/F_{PWIA}(p_z)$. At larger $p_z \gtrsim 1.2 \text{ fm}^{-1}$, the contribution from FSI enhances $F(p_z)$ by $\sim (30-40)\%$ as compared to the PWIA distribution $F_{PWIA}(p_z)$. At large $Q^2 \gtrsim 2 \text{ GeV}^2$ and $T_{kin} = Q^2/2m_p \gtrsim 1 \text{ GeV}$, the proton-neutron total cross section is approximately constant. Therefore, this large FSI effect in $F(p_z)$ should persist at large $Q^2 \gg 1 \text{ GeV}^2$ and should be taken into account in the asymptotic y -scaling regime, too. The departure from the PWIA further increases at $p_z \gtrsim 3 \text{ fm}^{-1}$, but

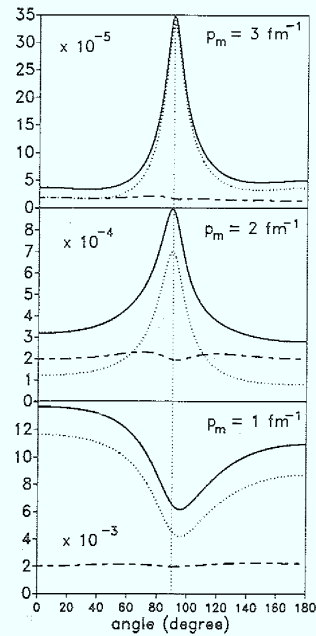


Fig. 2: The angular dependence of (solid curve) the missing momentum distribution $W(\vec{p}_m)$ and of its (dotted curve) S -wave and (dashed curve) D -wave components at different values of the missing momentum p_m .

the quantitative description of FSI in this region of p_m requires improving upon the Glauber approach, which goes beyond the scope of our investigation.

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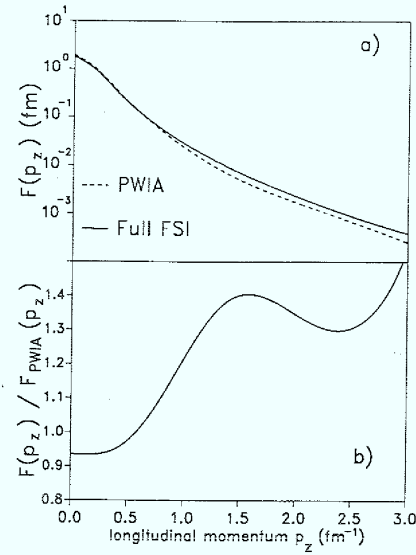


Fig. 3: (a) The longitudinal missing momentum distribution $F(p_z)$ calculated with full FSI (solid curve) and in the PWIA (dashed curve). (b) The ratio of the full $F(p_z)$ to the PWIA distribution $F_{PWIA}(p_z)$.

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Recently, the proton magnetic form factor has been determined in the time-like region by measuring the exclusive cross section for proton-antiproton annihilation into the electron-positron final state. The new data has been obtained for momentum transfers in the region $8.9 \text{ GeV}^2 \leq q^2 \leq 13 \text{ GeV}^2$ (Fermilab E-760) [1] and in the vicinity of the proton-antiproton threshold (LEAR PS170) [2]. The new Fermilab E-760 data confirm the predictions of perturbative QCD for the slope of magnetic form factor in the time-like region. The vector meson dominance (VMD) model [3, 4], on the other hand, predicts a slope of the magnetic form factor of the proton at the proton-antiproton threshold which is about five times smaller than the the new Low Energy Antiproton Ring (LEAR) data. Given the apparent failure of the vector meson dominance model and the success of pQCD, the data have been used to derive a running coupling constant in the time-like region [1].

We want to point out that the vector dominance model works quite well, if one includes the effects of excited vector meson resonances on the electromagnetic form factors. We employ a hybrid vector meson dominance formalism which is a generalization of the model developed by Gari and Krumpelmann [5]. The model provides a smooth transition from the low $-q^2$ behaviour predicted by vector meson dominance to the high $-q^2$ scaling behaviour predicted by perturbative QCD. The essential contribution of this work is to incorporate both time-like and space-like nucleon form factor data to determine the role of excited vector resonances. Details of the model can be found in [6].

In Fig.1, we display the result for the magnetic form factor of the proton. The parameterization gives a good account of the available data. The electric form factor of the proton and the magnetic form factor of the neutron are reproduced in a comparable quality [6]. We find that the behaviour of the data just above the $p\bar{p}$ threshold is due the $\omega'(1600)$ and $\rho'(1700)$ resonances. Close to $q^2 = 5 \text{ GeV}^2$, the form factor levels off. We assign this unusual behaviour to the effect of the ω'' and ρ'' mesons which are not yet established experimentally. Masses of 1.85 GeV and 2.15 GeV were therefore assumed for these mesons. Since the model is able to describe both electric and magnetic form factors of the proton and the magnetic form factor of the neutron, we investigated whether the new data help to constrain the electric form factor of the neutron. We found that the excited vector mesons play a negligible role in the space like region. The electric neutron form factor depends significantly on the coupling strength of the phi-meson. A detailed discussion of the effect of the phi-meson on the electric form factor of the neutron can be found in [6].

In conclusion, we have shown that the excited vector mesons are required in order to account for the puz-

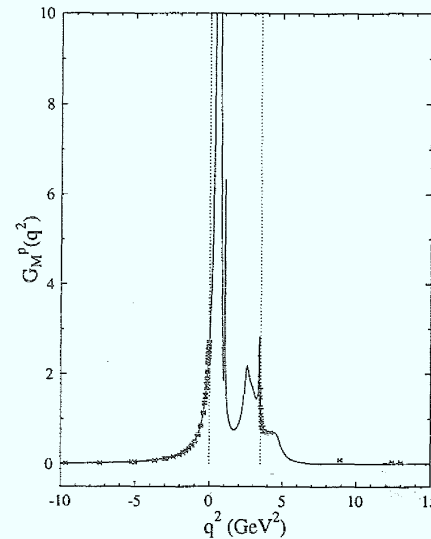


Fig. 1: Proton magnetic form factor

zling threshold behaviour of the form-factors in the time-like region.

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Previously we have analyzed data of pion- and photon induced reactions off ^{12}C in the Δ resonance region [1, 2, 3]. In this report we present a comparison of our results for elastic pion scattering, coherent pion photoproduction and coherent pion production from a charge exchange reaction $^{12}\text{C}(^3\text{He}, t\pi^+)^{12}\text{C}(\text{g.s.})$.

In the charge exchange reactions the target is exposed to the virtual pion- and rho-meson fields produced by the projectile-ejectile system. The pion-like interaction excites the spin-longitudinal (LO) response function of the target while the rho-meson-like interaction excites the spin-transverse (TR) response function. Due to the kinematics the virtual meson fields obey the energy-momentum relation $\omega < q$ and thus explore the LO and TR response functions in an $(\omega, \vec{q})$ region which is inaccessible to real pion and photon scattering. The two responses can be experimentally separated by measuring the exclusive $^{12}\text{C}(^3\text{He}, t\pi^+)^{12}\text{C}(\text{g.s.})$ reaction. This reaction triggers selectively on the LO channel. In Fig. 1 (a) we show the measured data of Hennino *et al.* [4] and compare them to previously [3] calculated angular distributions. The solid curve is the result from both LO and TR excitation. It turns out, however, that the contribution of the LO excitation to the pion production cross section is dominant while that of TR excitation is suppressed because of the unfavorable TR-LO channel mixing (see Fig. 2 (a)). The shape of the angular distribution is strongly forward peaked. This is an effect of the spin structure of the excitation ($\vec{S}^{\dagger} \cdot \vec{q}$) and de-excitation ($\vec{S} \cdot \vec{q}'$) operators involved in the LO channel. The product of both operators is proportional to $qq' \cos \theta_{\pi}$ and peaks at $\theta_{\pi} = 0^{\circ}$. In Fig. 1 (b) we show the corresponding elastic pion cross section at an appropriately chosen incident energy. One observes the direct proportionality between the coherent pion production cross section of Fig. 1 (a) and the elastic pion cross section of Fig. 1 (b). We remark that the coherent pion production process in the $^{12}\text{C}(^3\text{He}, t\pi^+)^{12}\text{C}(\text{g.s.})$ reaction can be viewed as 'virtual' pion scattering off ^{12}C where the initially off-mass shell pion is converted into an on-mass shell pion when the pion is emitted from the nucleus. The momentum needed to lift the pion on its mass shell is provided by the recoil nucleus.

In Fig. 2 we show angular distributions for the spin-transversely excited coherent pion production in charge exchange reactions (a) and for the pion photoproduction (b). Because of the $(\vec{S}^{\dagger} \times \vec{q})$ structure of the excitation operators the calculated angular distributions (solid curves) are now proportional to $(qq')^2 \sin^2 \theta_{\pi}$ and vanish at $\theta_{\pi} = 0^{\circ}$. This property is directly confirmed by the pion photoproduction data [6]. It can be noticed that the calculated TR angular distribution in Fig. 2 (a) peaks at a slightly smaller angle than the angular distribution of the pion photo-

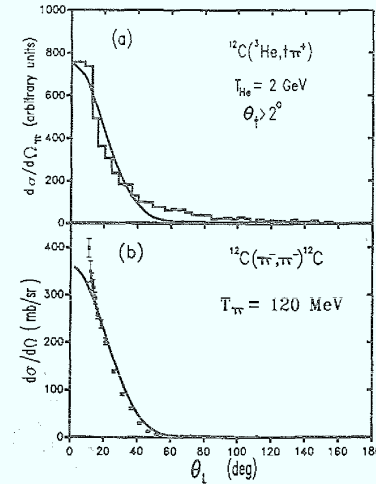


Fig. 1: Angular distributions of coherent pions for (a) the $^{12}\text{C}(^3\text{He}, t\pi^+)^{12}\text{C}(\text{g.s.})$ reaction at $T_{\text{He}} = 2$ GeV and for (b) the elastic pion scattering off ^{12}C at $T_{\pi} = 120$ MeV. (a) The solid curve represents the result of the complete calculation. The data are taken from Ref. [4]. (b) The solid curve represents the calculated elastic pion cross section. The data are taken from Ref. [5].

production in Fig. (b). The reason for this is twofold: on the one hand the charge exchange reaction involves a larger momentum transfer than the pion photoproduction. Therefore the nuclear form factor suppresses the cross section at higher pion angles in the charge exchange reaction. On the other hand the $(^3\text{He}, t)$ formfactor also cuts off the cross section at high momentum transfers, whereas the $\gamma N \Delta$ coupling is pointlike.

In Fig. 3 we compare the $^{12}\text{C}(^3\text{He}, t\pi^+)^{12}\text{C}(\text{g.s.})$ coincidence cross section with the measured data of Hennino *et al.* [4]. The solid and the dashed curve represent calculations with Landau-Migdal parameters of $g'_{\Delta\Delta} = 0.33$ and $g'_{\Delta\Delta} = 0.4$, respectively. We find that both the magnitude and the peak position of the calculated cross section are very sensitive to the choice of the $g'_{\Delta\Delta}$ parameter. While the solid curve reproduces the magnitude of the measured coherent pion production cross section very well [3] the dashed curve was multiplied by a factor of ~ 1.5 to match the data. The energy shift ΔE between both curves is directly proportional to the change in $g'_{\Delta\Delta}$.

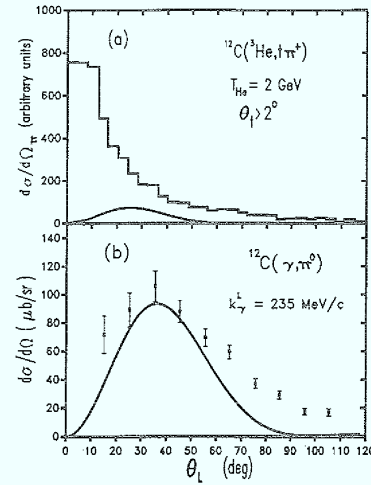


Fig. 2: Angular distributions of coherent pions for (a) the $^{12}\text{C}(^3\text{He}, t\pi^+)^{12}\text{C}(\text{g.s.})$ at $T_{\text{He}} = 2$ GeV and for (b) the coherent pion photoproduction on ^{12}C . (a) The solid curve represents the contribution to the pion production cross section resulting from the TR excitation of the nucleus. (b) The data for coherent pion photoproduction are taken from Ref. [6]. The solid curve represents the result of our calculation.

namely $\Delta E \approx \Delta g'_{\Delta\Delta} (\hbar c f_{\pi N\Delta}^2 / m_\pi^2) \rho_0$ with $\rho_0 = 0.17 \text{ fm}^{-3}$ being the central nuclear density. The proportionality of ΔE to ρ_0 is due to the fact that the coherent pions come from the nuclear interior. The $(^3\text{He}, t)$ kinematics forces the pions to travel into the direction of the momentum transfer $\vec{q}$. Therefore most of the coherent pions are detected in this direction. By changing the triton emission angle and the triton energy loss one can vary $\vec{q}$ and ω and can map out the energy and momentum dependence of the ΔN^{-1} residual interaction.

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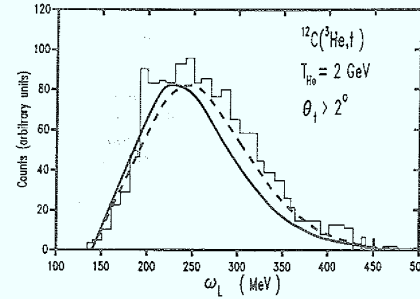


Fig. 3: Pion coincidence spectrum for the $^{12}\text{C}(^3\text{He}, t)$ reaction at $T_{\text{He}} = 2$ GeV and triton scattering angle $\theta_t \geq 2^\circ$. The data are taken from Ref. [4]. The solid and dashed curves represent cross section calculations with Landau-Migdal parameters of $g'_{\Delta\Delta} = 0.33$ and $g'_{\Delta\Delta} = 0.4$, respectively. The dashed curve has been multiplied with a factor of 1.5 in order to compare the peak positions.

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The $^{12}\text{C}(\pi^+, pp)^{10}\text{B}$ reaction

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In an other contribution to this annual report we have shown that the energy shift observed in the cross section of inclusive charge exchange reactions is due to a coherent medium effect on the spin-longitudinal response function. The medium effect produces a coherent pion wave in the nucleus. The pion wave is damped by the coupling to various reaction channels one of which is the two proton decay channel. In the $^{12}\text{C}(^3\text{He}, t2p)^{10}\text{B}$ reaction the coupling of the pion wave to the $2p$ decay channel was represented by a spin-isospin dependent zero range force. By comparison of the calculations to the $2p$ decay cross section data an interaction strength parameter of $V_0 = 220 \text{ MeVfm}^3$ was extracted. It is an important question whether this parameter is consistent with data of the $^{12}\text{C}(\pi^+, pp)^{10}\text{B}$ reaction. Therefore we have analysed this reaction within the Δ -isobar model at a pion kinetic energy of $T_\pi = 165 \text{ MeV}$ [1].

In Fig. 1 we compare our calculations for the $^{12}\text{C}(\pi^+, pp)^{10}\text{B}$ reaction with the data [2]. The threefold differential cross section is plotted versus the angle of the second proton, where the first proton was detected at $\theta_{p1} = 130^\circ$. The dashed curve shows the result of the calculations with $V_{\Delta N, \Delta N} = 0$ while the solid curve shows the result with $V_{\Delta N, \Delta N} \neq 0$. The dotted curve shows the result with additional inclusion of the distortion effects on the protons.

The inclusion of the residual interaction results in a reduction of the cross section by a factor of ~ 4 .

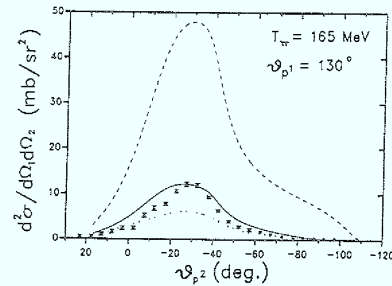


Fig. 1: Twofold differential cross section for the $^{12}\text{C}(\pi^+, pp)^{10}\text{B}$ reaction as function of the angle of the second proton, where the first proton was measured at $\theta_{p1} = 130^\circ$. The data are taken from ref. [2]. The dashed curve shows the result of our calculation with $V_{\Delta N, \Delta N} = 0$, the solid curve the result with $V_{\Delta N, \Delta N} \neq 0$. The dotted curve represents the result with $V_{\Delta N, \Delta N} \neq 0$ and inclusion of distortion effects on the outgoing protons.

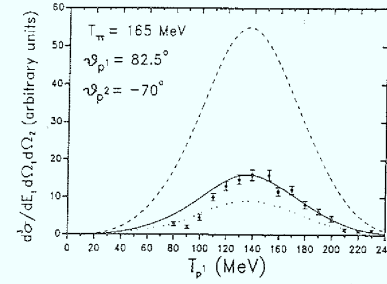


Fig. 2: Threefold differential cross section for the $^{12}\text{C}(\pi^+, pp)^{10}\text{B}$ reaction as function of the kinetic energy of the proton, which was measured at $\theta_{p1} = 82.5^\circ$. The angle of the second proton was fixed at $\theta_{p2} = -70^\circ$. The dashed curve shows the result of our calculation with $V_{\Delta N, \Delta N} = 0$, the solid curve the result with $V_{\Delta N, \Delta N} \neq 0$. The dotted curve represents the result with $V_{\Delta N, \Delta N} \neq 0$ and inclusion of distortion effects on the outgoing protons. Our calculation with inclusion of $V_{\Delta N, \Delta N}$ (solid curve) was normalized to the counting rate from ref. [2].

The distortion of the two outgoing protons leads to a further reduction of the cross section by a factor of ~ 2 . The calculation with inclusion of all distortion effects corresponds to a lower limit of the $2p$ decay cross section.

In Fig. 2 we show again results for the $^{12}\text{C}(\pi^+, pp)^{10}\text{B}$ reaction, where the experiment was performed in coplanar geometry with the angle of proton 1 fixed at $\theta_{p1} = 82.5^\circ$ and that of proton 2 at $\theta_{p2} = -70^\circ$. Since the data from ref. [2] are only counting rates, our calculation with inclusion of $V_{\Delta N, \Delta N}$ (solid curve) was normalized to these. Therefore we can not compare the absolute magnitude, but we describe the variation with the kinetic energy of proton 1 rather well. The good agreement between our calculation and the data shows that the coupling strength $V_0 = 220 \text{ MeVfm}^3$ is a good choice and is in agreement with the result from the $^{12}\text{C}(^3\text{He}, t2p)^{10}\text{B}$ experiment.

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The $^{12}\text{C}(\gamma, p\pi^-)^{11}\text{C}$ reaction

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Pion and photon induced nuclear reactions in the $\Delta(1232)$ resonance energy region contain different information on the behaviour of the Δ in the nucleus [1]. While the photon cross section ($\sigma_{\gamma A}/A$) is found to be practically independent of the nuclear mass number A , the pion cross section does show a strong A -dependence. In particular, the peak position of the pion-nucleus total cross section is systematically shifted downward in energy with mass number A relative to that of the free πN system. This is a result of the different spin structure of the $\pi N\Delta$ and $\gamma N\Delta$ couplings. The former is spin-longitudinal and proportional to $\vec{S} \cdot \vec{q} \vec{T}$ while the latter is spin-transverse and proportional to $\vec{S} \times \vec{q} \vec{T}$ (where $\vec{S}$ and $\vec{T}$ are the spin and isospin transition operators, respectively; $\vec{q}$ is the pion or photon momentum). Therefore in pion scattering the probe strongly couples to the virtual pion field inside the nucleus, which in our formalism is described by the ΔN^{-1} residual interaction. This strong coupling leads to pion multiple scattering inside the nucleus producing the energy shift and broadening in the elastic pion-nucleus cross section. In photon-induced reactions, however, the multiple pion scattering is strongly suppressed since the $\gamma N\Delta$ and $\pi N\Delta$ couplings are orthogonal to each other. Therefore the coherent π propagation is greatly inhibited in $\vec{q}$ direction. This makes the photon a good probe to selectively study the propagation of the Δ in the medium, i.e. to investigate the Δ -nucleus mean field without the effects due to pion multiple scattering.

In this report we use the $^{12}\text{C}(\gamma, p\pi^-)^{11}\text{C}$ reaction to investigate the Δ propagation in the nuclear mean field. In the analysis of the data we have to keep in mind, however, that the reaction mechanism of the $(\gamma, p\pi^-)$ reaction involves nucleon Born terms and Δ terms [2, 3]. The nucleon Born terms are large, actually often larger than the Δ terms. Therefore the information on the Δ propagation might be covered by the cross section contribution due to the Born terms. We remark that most calculations for the $(\gamma, p\pi^-)$ reaction assume a free Δ propagator. In the present report we only calculate the contributions of the Δ terms to the cross section including the Δ mean field effect and compare them to results of ref. [3] with a free Δ propagator.

The nuclear reaction formalism and the nuclear structure aspects of the problem are described in refs. [4] and [2, 3].

In Fig. 1 we show the results for the $^{12}\text{C}(\gamma, p\pi^-)^{11}\text{C}$ reaction at an incident photon energy of $E_\gamma = 380$ MeV. The threefold differential cross section is plotted versus the kinetic energy of the outgoing proton. In the coincidence measurement the decay proton and decay pion were measured in coplanar geometry keeping the proton and pion angles fixed at $\theta_p = 20^\circ$ and $\theta_\pi = -120^\circ$, respectively. The upper part of the

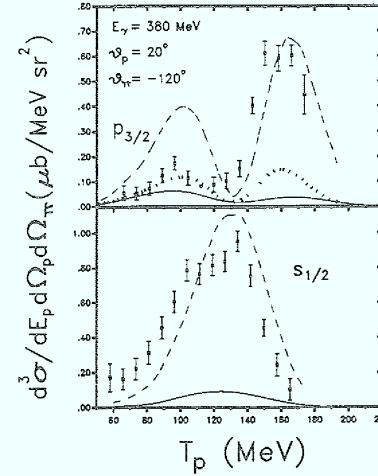


Fig. 1: Threefold differential cross section for the $^{12}\text{C}(\gamma, p\pi^-)^{11}\text{C}$ reaction as function of the kinetic energy of the outgoing proton. The data are taken from ref. [5]. The upper part of the figure shows the photon absorption on a $p_{3/2}$ neutron, the lower part the absorption on a $s_{1/2}$ neutron. The solid curves show the results of our calculations, the dashed curves the results of ref. [3] including Δ contributions and Born terms. The dotted curve represents the result of [3] with Δ contribution alone.

figure shows the cross section for the photon absorption on a $p_{3/2}$ neutron and the lower part shows the corresponding result for the photon absorption on a $s_{1/2}$ neutron. The distinction between the s shell and p shell photon absorption is experimentally possible since energy-momentum conservation

$$\begin{aligned} E_\gamma + m_A &= E_p + E_\pi + E_{A-1} - E_S \\ \vec{k}_\gamma &= \vec{q}_p + \vec{q}_\pi + \vec{q}_{miss} \end{aligned} \quad (1)$$

allows for the determination of the separation energy E_S of the neutron in A . In Eq. (1) $E_i = \sqrt{M_i^2 + \vec{q}_i^2}$ is the total energy of particle i ($i = A-1, p, \pi$), M_i is its rest mass, and $\vec{q}_i$ is its momentum. In impulse approximation the unobserved momentum $\vec{q}_{miss}$ is equal to the negative momentum of the bound nucleon before the photon absorption, i.e. $\vec{q}_{miss} = -\vec{q}_i$. In this approximation we may write

$$\vec{k}_\gamma + \vec{q}_i = \vec{q}_\Delta = \vec{q}_p + \vec{q}_\pi \quad (2)$$

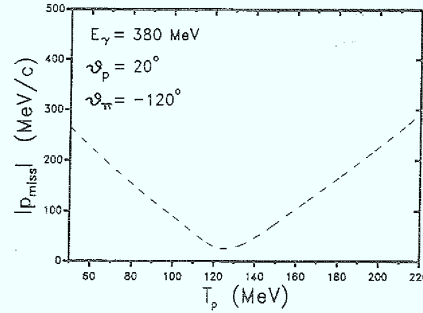


Fig. 2: $|\vec{q}_{miss}|$ as function of the kinetic energy of the outgoing proton for the $^{12}\text{C}(\gamma, p\pi^-)^{11}\text{C}$ reaction. $\vec{q}_{miss}$ is calculated with help of Eq. (1).

where $\vec{q}_\Delta$ is the momentum of the Δ . For a given kinetic energy T_p of the decay proton we find $E_\pi = E_\gamma - T_p + E_S$ and $p_\pi = \sqrt{E_\pi^2 - m_\pi^2}$. By variation of T_p for fixed pion and proton angle we can vary $\vec{p}_i$. Therefore the threefold differential cross section of Fig. 1 measures the momentum distribution of the neutron wave function via the $(\gamma, p\pi^-)$ reaction mechanism. Insofar the reaction mechanism is dominated by the Δ terms the cross section also includes information on the momentum distribution of the Δ wave function.

From Fig. 1 we recognize that the cross section for photon absorption on the $p_{3/2}$ neutron shows two maxima while that for the $s_{1/2}$ neutron has only one maximum. This is an effect of the kinematics of the reaction. In Fig. 2 we show $|\vec{q}_{miss}|$ as function of T_p . It can be recognized that $|\vec{q}_{miss}|$ is large for small proton energies, that it has a minimum for $T_p = 125$ MeV, and that it is growing again at higher proton energies. We find that the minimum of $|\vec{q}_{miss}|$ corresponds to the maximum in the $s_{1/2}$ neutron cross section. This is immediately understandable from the momentum distribution of the $s_{1/2}$ hole wave function which has its maximum value at $\vec{q}_i = -\vec{q}_{miss} = 0$. The situation is a little more complicated for the γ absorption on the $p_{3/2}$ neutron. The $p_{3/2}$ hole wave function has its maximum near $|\vec{q}_i| = |\vec{q}_{miss}| = 110$ MeV/c. From the kinematic conditions of Eq. 1 we find that the value of $|\vec{q}_{miss}| = 110$ MeV/c can be met by proton energies of $T_p = 90$ MeV and $T_p = 180$ MeV, respectively. In the first case $\vec{k}_\gamma$ and the momentum of the struck nucleon, $\vec{p}_i$, point in opposite direction while in the second case they point in the same direction. Therefore a Δ^0 created in the scattering process receives a relatively small momentum $\vec{p}_\Delta$ for $T_p = 90$ MeV and a large momentum for $T_p = 180$ MeV. The cross section ratio between both kinematical situations should contain information on the Δ propagation in the nucleus.

In Fig. 1 we compare our cross section calculations for the $(\gamma, p\pi^-)$ reaction via intermediate Δ^0 excitation (solid curve) with a corresponding calculation of Johansson and Sherif (dotted curve) [3]. The difference between both calculations consists in the choice of the Δ propagator. We consider the Δ propagation in the nuclear potential while Johansson and Sherif use a free Δ propagator. We observe a strong reduction in cross section with our propagator. The dashed curve is again a result of Johansson and Sherif including both the nucleon Born terms and the Δ terms in the reaction mechanism. From the relative enhancement between the dashed and the dotted curves one can estimate the importance of the nucleon Born terms for the $(\gamma, p\pi^-)$ reaction. All calculations shown in Fig. 1 consider the nuclear distortion effect on the final proton and pion wave functions. In conclusion our calculations within the Δ -hole model show that by consideration of the Δ propagation in the medium the quasifree decay of Δ 's with small momentum ($T_p = 90$ MeV) are favoured relative to the quasifree decay of Δ 's with large 3-momentum ($T_p = 180$ MeV).

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The $^{12}\text{C}(\pi^+, p\pi^+)^{11}\text{B}$ reaction

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In the preceding report we found that in the $^{12}\text{C}(\gamma, p\pi^-)^{11}\text{C}$ reaction the reaction mechanism is dominated by the nucleon Born terms and that the Δ terms are relatively small. In the $^{12}\text{C}(\pi^+, p\pi^+)^{11}\text{B}$ reaction, however, the situation is known to be just the opposite in the sense that the reaction mechanism is now dominated by the Δ terms and the nucleon Born terms are suppressed. As a new feature, however, in pion-induced reactions the multiple scattering of the pion inside the nucleus becomes very important. Since we describe the pion propagation in the Δ -isobar model on a quantitative level we obtain a good description of the $^{12}\text{C}(\pi^+, p\pi^+)^{11}\text{B}$ data [1].

In Fig. 1 we compare our calculations for the $^{12}\text{C}(\pi^+, p\pi^+)^{11}\text{B}$ reaction with data at $T_\pi = 245$ MeV [2]. The angles of the outgoing proton and pion were taken in coplanar geometry and fixed at $\theta_p = 17.5^\circ$ and $\theta_\pi = -130^\circ$. In the experimental data no distinction between the pion absorption on the $s_{1/2}$ and $p_{3/2}$ protons was made. In Fig. 1 the threefold differential cross section is plotted versus the kinetic energy of the outgoing proton. The dashed curve shows the result of the calculations with $V_{\Delta N, \Delta N} = 0$ while the solid curve shows the result with $V_{\Delta N, \Delta N} \neq 0$. The dotted curve shows the result with additional inclusion of the distortion effect on the proton and pion wave functions.

By the comparison of the calculations with and without inclusion of $V_{\Delta N, \Delta N}$ we recognize that there

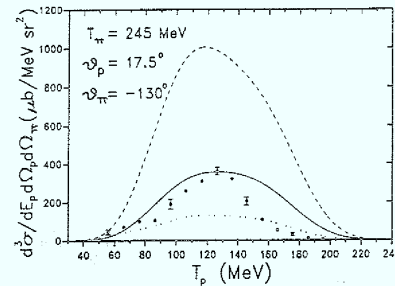


Fig. 1: Threefold differential cross section for the $^{12}\text{C}(\pi^+, p\pi^+)^{11}\text{B}$ reaction as function of the kinetic energy of the outgoing proton. The data are taken from ref. [2]. The dashed curve shows the result of our calculation with $V_{\Delta N, \Delta N} = 0$, the solid curve the result with $V_{\Delta N, \Delta N} \neq 0$. The dotted curve represents the result with $V_{\Delta N, \Delta N} \neq 0$ and inclusion of distortion effects on the outgoing proton and pion.

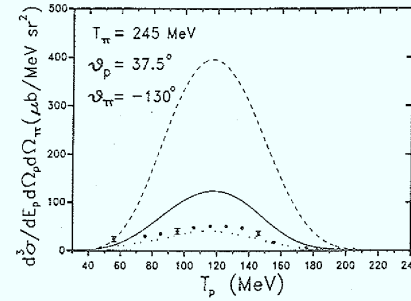


Fig. 2: Same as in Fig. 1, but the angles of the outgoing proton and pion were fixed at $\theta_p = 37.5^\circ$ and $\theta_\pi = -130^\circ$.

is a strong effect of the residual interaction on the quasifree decay cross section. In the calculation with inclusion of $V_{\Delta N, \Delta N}$ the cross section is by about a factor of ~ 2.5 smaller than in the calculation without inclusion of $V_{\Delta N, \Delta N}$. Obviously we obtain a good description of the data in the former case. The strong reduction in cross section is a result of the multiple pion scattering in the nucleus. The reduction factor of ~ 2.5 is not too surprising since it is comparable to the reduction factor of ~ 4 observed in the total pion-nucleus cross section. The inclusion of the distortion effects on the outgoing proton and pion leads to a further reduction of the cross section by a factor of ~ 3 . We used the optical potentials, as derived from elastic proton-nucleus and elastic pion-nucleus scattering. In this way we probably overestimate the absorption and thus obtain a lower limit for the $^{12}\text{C}(\pi^+, p\pi^+)^{11}\text{B}$ cross section.

In Fig. 2 we show the threefold differential cross section for the reaction $^{12}\text{C}(\pi^+, p\pi^+)^{11}\text{B}$ at $T_\pi = 245$ MeV. The angle of the outgoing proton corresponds here to $\theta_p = 37.5^\circ$ and $\theta_\pi = -130^\circ$. Again we observe a reduction of the cross section by a factor of ~ 4 due to the residual interaction and by a further factor of ~ 2 due to the distortion. We obtain a good description of the data even in this case where we chose an unfavoured combination of proton and pion angles.

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The $^{12}\text{C}(^3\text{He},\text{tp}\pi^+)$ and $^{12}\text{C}(^3\text{He},\text{tpp})$ coincidence reactions

P. Oltmanns, F. Osterfeld, and T. Udagawa*

Recently, new information on the Δ propagation in nuclei has been obtained from intermediate energy ($^3\text{He},\text{t}$) charge exchange reactions[1, 2]. Similar to the π -nucleus total cross section data, the inclusive cross section data of the charge exchange reactions show also a substantial downward energy shift of the Δ resonance in nuclei. The analysis of the data shows[3, 4, 5, 6] that a large part of this shift is due to a nuclear medium effect on the spin-longitudinal response function. The medium effect is caused by the strongly attractive energy (ω)-dependent π -exchange interaction $V_{\Delta N,\Delta N}$ between Δ -particle nucleon-hole (ΔN^{-1}) states in the nucleus. This interaction produces a coherent pion wave in the medium.

The damping of the coherent pion wave can be studied by measuring the various decay channels. In the following we discuss the $^{12}\text{C}(^3\text{He},\text{tp}\pi^+)$ and $^{12}\text{C}(^3\text{He},\text{tpp})$ coincidence reactions. In Fig. 1 (a) and 1 (b) we show the corresponding coincidence spectra measured by Hennino *et al.* at Saclay in Paris[2]. In Fig. 1 (a) we compare the results of our microscopic $^{12}\text{C}(^3\text{He},\text{tp}\pi^+)$ decay calculation with the data. The calculations are performed by projecting the many-body Δ -nucleus wave function onto the $\pi\pi^+$ decay channel using a microscopic model for the decay matrix elements. The model includes the distortions of the outgoing pion and proton. The decay interaction is fixed from the free decay width of the Δ resonance. Without any further adjustment we produce the solid curve in Fig. 1 (a). We remark that we see no effects of the $V_{\Delta N,\Delta N}$ residual interaction in the $\pi\pi^+$ decay cross section. This indicates that the $\pi\pi^+$ events come from the nuclear surface and carry no information of the coherent pion propagation in the nucleus.

In Fig. 1 (b) we compare our microscopically calculated $^{12}\text{C}(^3\text{He},\text{tpp})$ coincidence spectrum with the data of Hennino *et al.*[2]. The microscopic model for the $2p$ decay includes the distortions of the outgoing protons and assumes that the $\Delta+N \rightarrow N+N$ decay interaction can be approximated by a pointlike interaction with a strength of $V_0 = 220 \text{ MeV fm}^3$. With this single parameter we can reproduce the experimental data in Fig. 1 (b). Two calculations are compared with the data: The solid curve shows the result of our calculation with inclusion of $V_{\Delta N,\Delta N}$, while the dashed curve shows the result without inclusion of the residual interaction, respectively. The residual interaction shifts the $2p$ decay spectrum down in energy by 20 MeV relative to the result with $V_{\Delta N,\Delta N} = 0$. This is so since the two protons carry in average 100 MeV kinetic energy and thus suffer from little distortion. Therefore they can come out from the nuclear interior and carry information on the coherent pion propagation.

In conclusion we have shown that the existence of the coherent pion wave in the nucleus, as measured in

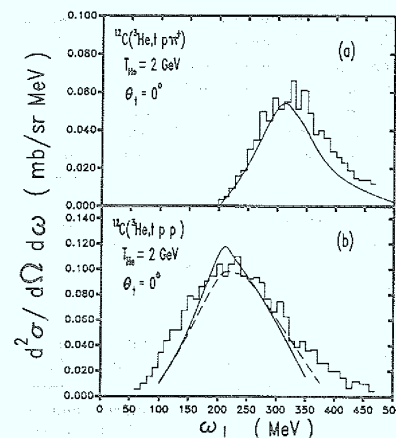


Fig. 1: Coincidence spectra for the $^{12}\text{C}(^3\text{He},\text{t})$ reaction at $T_{\text{He}} = 2 \text{ GeV}$ and triton scattering angle $\theta_t = 0^\circ$. (a) The $^{12}\text{C}(^3\text{He},\text{tp}\pi^+)$ data[2] in comparison with the result of our calculation. (b) The $^{12}\text{C}(^3\text{He},\text{tpp})$ reaction data[2]. The solid and dashed curve show the results with and without inclusion of the residual interaction, respectively.

intermediate energy charge exchange reactions, has also an appreciable effect on the $2p$ decay spectrum. The $\pi\pi^+$ spectrum is unaffected.

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Extremely strong electromagnetic fields of extremely short duration accompany the passage of nuclei colliding at LHC or RHIC energies. The production of dileptons, especially e^+e^- pairs („QED electrons“) is most conspicuous.

Electromagnetic pair production with e^- capture is calculated in Ref. 1. This is also useful for the consideration of luminosity decrease at the LHC (Pb-Pb)-collider. The Weizsäcker-Williams method is not generally applicable for the case of e^+e^- production. This is essentially due to the low mass of the electrons, $m_e < \frac{1}{R}$, where R is the nuclear radius. In the external field approximation, differential probabilities for e^+e^- pair production are calculated [2] for $b=0$, using the method of Feynman integrals.

For very heavy systems, lowest order perturbation theory is not justified, and multiple pair production must be taken into account [3]. E.g. for Pb-Pb collisions at LHC energies, about 3-4 e^+e^- pairs (mainly soft ones) are produced on the average in one collision, at impact parameters ranging from 0 up to about the Compton wave length of the electron ($\lambda_e=386$ fm). The impact parameter dependence is calculated in Ref. 4, using Bessel transformation and Feynman integral methods. Similar results were recently also obtained by Güçlü et al. [5] using Monte Carlo techniques.

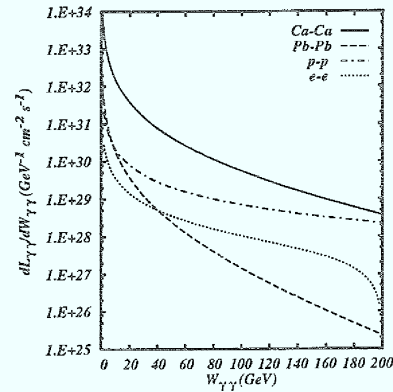


Fig. 1: The $\gamma\gamma$ -luminosity function as a function of the invariant mass $W_{\gamma\gamma}$. We take the following collider parameters (L is the luminosity in units of $\text{cm}^{-2}\text{s}^{-1}$, E_b the beam energy). LEP200; $L=10^{32}$, $E_b=100$ GeV, p-p: $L=10^{33}$, $E_b=7$ TeV. Ca-Ca: $L=5\cdot 10^{30}$, $E_b=3,5$ TeV/A, Pb-Pb: $L=5\cdot 10^{26}$, $E_b=2.76$ TeV/A.

In Fig. 1 $\gamma\gamma$ luminosity functions for various colliders are shown as a function of the invariant mass $W_{\gamma\gamma}$ of the $\gamma\gamma$ -system. It can be seen that Ca-Ca collisions at the planned LHC would be of great interest for photon-photon physics. At the high energy end ($W_{\gamma\gamma}$ around 200 GeV) the p-p collisions become competitive. In this case, also substantial inelastic contributions, due to the incoherent radiation from the quarks is expected [6]. This will not be the case for the heavy ion collisions [7,8].

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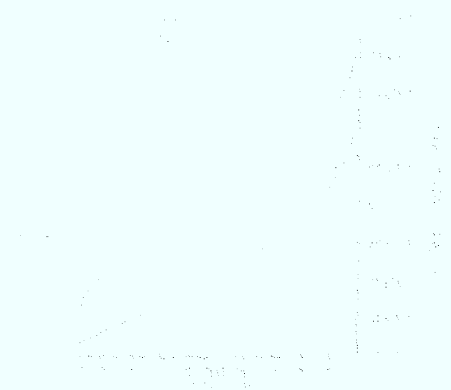
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4. NUCLEAR STRUCTURE AND REACTION MECHANISM

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In peripheral heavy-ion collisions at relativistic bombarding energies $\gtrsim 1$ GeV/nucleon the double-giant-dipole resonance (*DGDR*) is strongly excited [1]. The physical reason is a strong longitudinal focusing of the electromagnetic field in the target rest frame. This greatly enhances the field intensity in the vicinity of the target nucleus thus increasing the probability for two-photon absorption from the ground state.

Most simply, the *DGDR* is understood as a two-phonon excitation (with $J^\pi = 0^+, 2^+$ in a spherical nucleus) which, in the harmonic limit, is located at twice the giant dipole resonance (*GDR*) energy. This expectation is largely confirmed by experiment [1]. On the other hand, the single-dipole state

$$|D\rangle \equiv \sum_{i=1}^Z \tilde{r}_i \tau_i^z |0\rangle, \quad (1)$$

as well as the double-dipole state

$$|DD\rangle \equiv \left(\sum_{i=1}^Z \tilde{r}_i \tau_i^z \right) \left(\sum_{j=1}^Z \tilde{r}_j \tau_j^z \right) |0\rangle, \quad (2)$$

are not eigenstates of the nuclear Hamiltonian and hence acquire a width. Here the experimental situation is less clear and values for Γ_{DGDR} are observed which are bracketed by $\sqrt{2}\Gamma_{GDR}$ and $2\Gamma_{GDR}$ [1]. The lower value is expected from a folding of two Gaussians while the larger width results from the folding of two Lorentzians (recall that the Lorentzian gives a good fit to the total photoabsorption cross section).

To reach a quantitative understanding of the *DGDR* characteristics we have performed calculations of the single- and double-dipole response function by diagonalising a given model Hamiltonian, $\hat{H}$, in the space of *1p1h*- and *2p2h* excitations avoiding diagrammatic selections and fully respecting the Pauli principle. The theory has been used previously to provide a realistic description of linear response functions over a wide range of excitation energies [3].

We start by expanding the wave packets $|D\rangle$ and $|DD\rangle$ (eqs. (1) and (2)) in exact eigenstates $|N\rangle$ of the nuclear Hamiltonian $\hat{H}$ as

$$\begin{aligned} |D\rangle &= \sum_{N \neq 0} \langle N | \hat{D} | 0 \rangle |N\rangle \\ |DD\rangle &= \sum_{N \neq 0} \langle N | \hat{D} \hat{D} | 0 \rangle |N\rangle \end{aligned} \quad (3)$$

where $\hat{D} = \sum_{i=1}^Z \tilde{r}_i \tau_i^z$. When writing $\hat{H}$ as a mean field part and a residual interaction

$$\hat{H} = \sum_i \epsilon_i a_i^\dagger a_i + \frac{1}{4} \sum_{ij,kl} v_{ij,kl} a_i^\dagger a_j^\dagger a_l a_k, \quad (4)$$

($v_{ij,kl}$ denote antisymmetrized two-body matrix elements) the exact eigenstates $|N\rangle$ are obtained by diagonalization in the space of *npnh* excitations. A minimal

truncation of this, prohibitively large, space which can accommodate both single- and double dipole excitations must clearly encompass the *1p1h* and *2p2h* sector, so that

$$|N\rangle = \sum_1 c_1^N |1\rangle + \sum_2 c_2^N |2\rangle. \quad (5)$$

Due to the one-body nature of the dipole operator, the overlap matrix elements $\langle N | \hat{D} | 0 \rangle$ then project onto the *1p1h* subspace, $|1\rangle$, with

$$\langle N | \hat{D} | 0 \rangle = \sum_1 c_1^{N*} \langle 1 | \hat{D} | 0 \rangle \quad (6)$$

while $\langle N | \hat{D} \hat{D} | 0 \rangle$ project onto the *2p2h* space, $|2\rangle$, with

$$\langle N | \hat{D} \hat{D} | 0 \rangle = \sum_2 c_2^{N*} \langle 2 | \hat{D} \hat{D} | 0 \rangle. \quad (7)$$

Knowing the expansion coefficients, $c_{1,2}^N$ and the eigenvalues E_N of the Hamiltonian matrix in these spaces one then readily computes the 'strength functions'

$$S_D(E) = \sum_N |\langle N | \hat{D} | 0 \rangle|^2 \delta(E - E_N)$$

$$S_{DD}(E) = \sum_N |\langle N | \hat{D} \hat{D} | 0 \rangle|^2 \delta(E - E_N). \quad (8)$$

The angular-momentum-coupled expression for S_{DD} is somewhat involved and will be given elsewhere [6]. For comparison with experiment it is useful to define the centroids and variances of the strength distributions (8) in terms of their energy moments

$$m_{D(DD)}^k = \int dE E^k S_{D(DD)}(E) \quad (9)$$

as

$$\begin{aligned} E_{GDR} &= m_D^1 / m_D^0 \\ \sigma_{GDR} &= \sqrt{m_D^2 - (m_D^1)^2 / m_D^0} \end{aligned} \quad (10)$$

and similarly for the double-dipole excitation.

Although not measured, so far, we have selected the ^{40}Ca nucleus since the dimensions of the model space can be kept under reasonable numerical control. For the mean field we have chosen the same Woods-Saxon potential as in ref. [7]. Continuum excitations have been treated by discretization through an expansion of the positive-energy single-particle states in a harmonic oscillator basis. The density-dependent zero-range interaction of ref. [7] has also been adopted. Its form ensures the Pauli principle. To get a quantitative description of the *GDR* we have increased the isovector part, to 300 MeV fm³ for, both, the interior and exterior part of V_{01} . When truncating to *1p1h* and *2p2h* configurations with excitation energies ≤ 30 MeV (which yields 375 states) one obtains the strength distribution displayed in the upper part of Fig. 1. The centroid energy, E_{GDR} is at 19.7 MeV as compared to 20.3 MeV, deduced the global A-dependence

$E_{GDR} = 31.2A^{-1/3} + 20.6A^{-1/6}$ [9]. Upon further increase of the model space, both the total transition strength and the centroid energy remain unchanged. For further comparison, the lower part of Fig. 1 shows the normalized photo absorption cross section [8]

$$\sigma_{\gamma}(E) = \sum_N |\langle N | \hat{D} | 0 \rangle|^2 E_N \delta(E - E_N) / m_D^1 \quad (11)$$

together with data from Compton scattering [10] (for the theoretical distribution a Lorentzian smoothing of 1 MeV has been used).

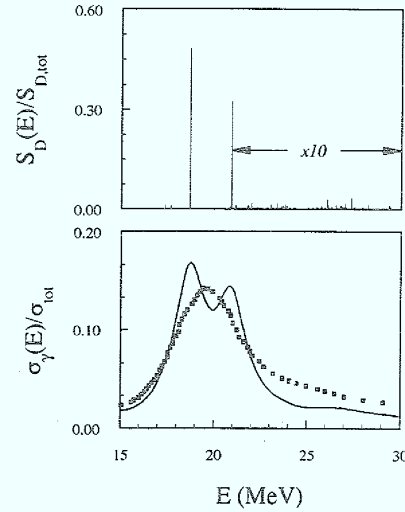


Fig. 1. upper part: the calculated dipole transition strength distribution in ^{40}Ca (eq. (8)). The high-energy tail has been multiplied by a factor of 10; lower part: the photoabsorption cross section, normalized to the TRK sum rule. The data, indicated by the filled squares, are inferred from the Compton-scattering analysis of ref. [9].

While not reproduced in all detail the overall agreement with experiment is quite satisfactory.

In the harmonic limit, the double-dipole state is expected at $2 \times E_{GDR}$ and a truncation at 30 MeV excitation energy is clearly insufficient. On the other hand the density of $2p2h$ states increases rapidly with excitation energy. A reasonable compromise is to truncate at 45 MeV (where the number of states for the 0^+ -component of the DGDR is 1067 while for the 2^+ -component one obtains 4144 states). At 45 MeV the total transition strength as well as the centroid energy are saturated. Any further increase in truncation energy will thus only influence the variance, σ_{DGDR} , i.e. the width. To test harmonicity one can compare the DGDR centroid energies $E_{DGDR}^{0^+} = 38.67$ MeV and $E_{DGDR}^{2^+} = 38.38$ MeV to $2 \times E_{GDR} = 39.34$ MeV, indicating that the anharmonicity is quite small. This is

corroborated by the fact that the 0^+ - and 2^+ components are nearly degenerate. The normalized DGDR strength distributions, (Fig. 2), indicate some differences between the 0^+ - and 2^+ components, however, in fine structure as well as in the overall width (the 0^+ -component is not strongly excited in experiment [11]).

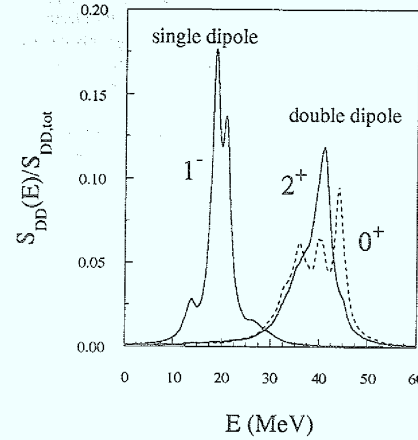


Fig. 2. The transition-strength distribution for the 0^+ -component (dashed line) and 2^+ -component (full line) the DGDR in ^{40}Ca . For comparison the single-dipole strength function is also given (short-dashed line). In all cases the discrete spectra have been folded with a 1 MeV-width Lorentzian (the experimental resolution is typically 1-2 MeV [1]).

To decide whether $\Gamma_{DGDR}/\Gamma_{GDR}$ is closer to $\sqrt{2}$ or 2 we have determined the ratio of variances $\sigma_{GDR}/\sigma_{DGDR}$ yielding 1.61 for the 0^+ -component and 1.24 for 2^+ component (σ_{GDR} was evaluated in a window from 10-30 MeV while σ_{DGDR} was obtained in a window from 30-50 MeV). It would be premature to conclude that a ratio of $\sqrt{2}$ is favored, however, since the model space might still be too small to get a reliable answer of the width. We estimate that this limitation can be overcome. Before embarking on large-scale calculations, however, one should evaluate the Coulomb excitation cross section for direct comparisons with experiment [12].

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One of the most interesting results of the previous study [1] is that the isoscalar state remains much more localized than the isovector state even though they evolve in the same phase-space. Is it a manifestation of stronger collectivity of the isoscalar state or, perhaps, is it that the specific properties of the coupling matrix elements block certain regions of the phase space and ergodization only occurs in the unblocked regions? A quantity which appears helpful in resolving this question originates from the concept of entropy. The information entropy of the state $|F\rangle$ can be defined, in a given basis $|k\rangle$, as

$$K = - \sum_k p_k \ln p_k, \quad (1)$$

where $p_k = |\langle k|F\rangle|^2$. It provides a quantitative measure of the complexity of the state $|F\rangle$ and its localization length in the basis $|k\rangle$ [2]. The so-defined K is, in principle, basis dependent, but the physically preferred basis is determined by the mean field. As the most smooth component of the nuclear Hamiltonian [3] it provides a natural reference for quantifying local GOE-type fluctuations.

In our case the mean field basis corresponds to the unperturbed basis of states $|1\rangle$ and $|2\rangle$. Calculating $K(t)$ along the 'trajectory' $|F(t)\rangle$ for the isovector and isoscalar states, we obtain asymptotically values of 7.10 and 6.53, while the corresponding initial values are 2.56 and 2.40, respectively. This is to be compared to $K^{GOE} = 7.29$ ($K^{GOE} = \psi(N/2 + 1) - \psi(3/2)$ [2], where ψ is the digamma function and N is the number of basis vectors). A comparison of these numbers indicates non-uniformities in the p_k distribution, especially for the isoscalar excitations. Actually, even the GOE-type fluctuations result in a gaussian distribution which is non-uniform (a uniform distribution maximizes the entropy and corresponds to $p_k = 1/N$ which, for $N=3040$, gives $K \approx 8.02$).

In view of the above mentioned non-uniform phase-space exploration, we find it instructive to calculate the spectrum of q -moments for $\{p_k\}$ and to introduce a generalized entropy [4]

$$K_q = \frac{1}{1-q} \ln \sum_{k=1}^N p_k^q. \quad (2)$$

From this definition it follows that $K_{q_1} \leq K_{q_2}$ if $q_2 < q_1$ (provided $\sum_k p_k = 1$). Equality holds for the uniform distribution. For $q \rightarrow 1$ eq. (2) yields the information entropy (eq. (1)). The most important property of K_q is that with increasing q a higher weight is given to the largest components in the set $\{p_k\}$. For $q \rightarrow 0$, on the other hand, K_q just counts the number of sites (here the basis vectors $|k\rangle$) visited, irrespective of how frequently they are sampled. For this reason eq. (2) also constitutes a basis for

defining the multifractal dimensions of non-uniform fractal sets [5].

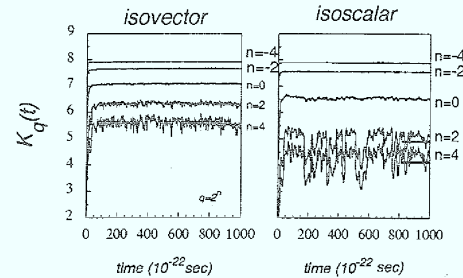


Fig. 1. The time evolution of the generalized entropies defined by eq. (2) for the isovector (l.h.s) and isoscalar (r.h.s) giant quadrupole resonances in ^{40}Ca coupled to the background characterized by GOE fluctuations. The horizontal marks on the left hand side of each panel denote the asymptotic values of K_q for $q=4$ and $q=16$.

For selected q -values Fig. 1 compares the time evolution of $K_q(t)$ for the isovector and isoscalar excitations when the background states have GOE fluctuations (case (c)). As one can see from the large- q behavior of $K_q(t)$, which are systematically smaller for the isoscalar excitation, the large components of these remain much more localized (larger) than those of the isovector excitation. Since, by probability conservation, the number of significant components is smaller in the former case, the amplitude of oscillations is larger in the corresponding $K_q(t)$. On the other hand, the dynamics start to look similar in both cases as q decreases and, for $q \rightarrow 0$, K_q approaches a value of 8. This signals that, on the level of small probabilities, the whole space spanned by 3040 states is visited. This aspect of the dynamics is consistent with the scaling properties of the transition strength distribution for the isovector and isoscalar states discussed in ref. [6]. On the level of small components they both scale.

For the purpose of addressing specific questions concerning the quantum-mechanical phase space exploration of collective modes, the full set of "generalized entropies" appears to be a useful theoretical tool.

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The study of chaos in nuclear physics has been mostly based – so far – on level statistics. In practical terms this is rather restrictive since a reliable statistical analysis requires very precise energy resolution. It also does not provide firm means for investigating the role of collectivity and mechanisms of its coexistence with chaos. In this respect the study of temporal correlations between an initially prepared non-stationary state and a state to which it evolves seems to be much more appropriate. Nuclear giant resonances are of central interest in this connection because they are located in a region of high level density which is expected to be dominated by chaotic dynamics. To make the theoretical studies meaningful one needs a scheme which incorporates the relevant elements, such as the possibility of defining a physical collective state, a realistic modeling of the background states whose complexity is consistent with the Gaussian orthogonal ensemble (GOE) of random matrices and, finally, the realistic coupling between the two.

The recently developed model [1], based on a diagonalization of the full nuclear Hamiltonian consisting of a mean field part and a residual interaction

$$\hat{H} = \sum_i \epsilon_i a_i^\dagger a_i + \frac{1}{2} \sum_{ij,kl} v_{ij,kl} a_i^\dagger a_j^\dagger a_l a_k, \quad (1)$$

in the basis of 1p-1h and 2p-2h states

$$|1\rangle \equiv a_p^\dagger a_h |0\rangle; \quad |2\rangle \equiv a_{p_1}^\dagger a_{p_2}^\dagger a_{h_1} a_{h_2} |0\rangle \quad (2)$$

fulfills these requirements and proves numerically manageable [2]. A prediagonalization of the two-body interaction v in the 1p-1h and 2p-2h subspaces defines $|\bar{1}\rangle = \sum_1 C_1^1 |1\rangle$ and $|\bar{2}\rangle = \sum_2 C_2^2 |2\rangle$, and the coupling is mediated by the off-diagonal elements $\langle \bar{1} | v | \bar{2} \rangle$ and their complex conjugate. An initially excited state, in response to an external one-body field $\hat{F} = \sum_{ij} F_{ij} a_i^\dagger a_j$ can be represented as

$$|F\rangle = \hat{F}|0\rangle = \sum_i F_i |\bar{1}\rangle. \quad (3)$$

As soon as the coupling between the subspaces $|\bar{1}\rangle$ and $|\bar{2}\rangle$ is taken into account, the state originally localized in the 1p-1h subspace, as defined by eq. (3), starts leaking into the 2p-2h space. The degree of mixing depends not only on the magnitude of v but also on the nature of the energy fluctuations in the 2p-2h space. The most natural quantity for describing the leakage is the survival probability, defined as

$$P(t) = |\langle F(0) | F(t) \rangle|^2. \quad (4)$$

The physical significance of $P(t)$ can be identified from its relation to the spectral autocorrelation function $G(E)$ [3]

$$P(t) = \int dE e^{-iEt/\hbar} G_F(E), \quad (5)$$

where

$$G_F(E) = \int dE' S_F(E') S_F(E' + E) \quad (6)$$

and $S_F(E)$ is the transition strength distribution $S_F(E) = \sum_n |F_n|^2 \delta(E - E_n)$. Thus $P(t)$ can be obtained from experiments which measure $S_F(E)$.

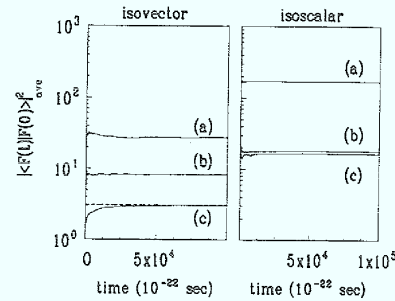


Fig. 1. The time-averaged survival probability $\bar{P}(t)$ in units of $1/N$ (N is the number of vectors spanning the space, 3040 in the present case) for the isovector (l.h.s.) and isoscalar (r.h.s.) quadrupole resonances: (a) no residual interaction in 2p-2h subspace; (b) including only particle-particle and hole-hole matrix elements in the diagonalization of the 2p-2h subspace; (c) diagonalization of the full residual interaction in the 2p-2h subspace

The calculations presented below for quadrupole excitations in ^{40}Ca are performed in the same basis as in ref. [2], i.e. including all 1p-1h and 2p-2h states up to 50 MeV and using the same residual interaction. Since the present study concentrates on the phase-space exploration and the role played by chaotic dynamics, we distinguish three cases corresponding to different classes of the spectral fluctuations in the 2p-2h subspace. As established in ref. [1] one finds: (a) with no residual interaction in the 2p-2h subspace there are many degeneracies in $|\bar{2}\rangle$ ($= |\bar{2}\rangle$) and the nearest-neighbor spacing is strongly peaked near zero; (b) inclusion of particle-particle and hole-hole matrix elements in $\langle \bar{2} | v | \bar{2} \rangle$ removes all degeneracies and leads to a Poissonian distribution of the nearest-neighbor spacings, characteristic of generic integrable systems; and (c) use of the full residual interaction yields GOE fluctuations, characteristic of chaotic dynamics.

The initial state $|F(0)\rangle = \hat{F}|0\rangle$ (eq. (3)) is already non-stationary in the 1p-1h subspace and therefore

An additional requirement for chaoticity is the initial dephasing of $P(t)$ below its asymptotic value [3]. Such a dephasing, indeed, takes place for the isovector case, as can also be seen from Fig. 1. As a consequence of the high density of states the corresponding 'correlation hole' extends over a time interval, four orders of magnitude longer than the characteristic 'excitation time' of $\sim 10^{-22}$ sec. The isoscalar excitation (right panel of Fig. 1) only shows a trace of such a behavior and the asymptotic values of $P(t)$ are systematically larger even though the initial state is coupled to the same background.

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A MICROSCOPIC THEORY OF GIANT RESONANCES IN NON-MAGIC NUCLEI

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The modern status of experimental giant resonance physics is determined by the following main directions:

1. Study of giant resonances in excited states, especially in hot nuclei. A huge amount of information on dipole resonances including the behaviour of the widths has already been obtained.
2. Study of giant resonances in unstable nuclei which can be produced in the radioactive beams facilities. There are no experimental data here at present but this direction offers a wide perspective and is very promising.
3. The quick improvement of the experimental resolution in the regions of binding energies and giant resonances, e.g. about 40 keV FWHM [1]. Very many experimental data about gross and fine structure have been accumulated and more will be obtained in the nearest future.

The first and the second directions are also of great interest for astrophysics.

These directions together with the coincidence experiments, which recently became one of the main methods in the investigations of the decay properties of resonances, confirm very strongly the necessity of an improvement of the RPA approach in magic nuclei or the QRPA in non-magic nuclei. These improvements are: taking into account the single particle continuum and, what is much more complicated and necessary in medium weight and heavy nuclei, the configurations which are more complex than those accounted for by the RPA or QRPA. Now it is clear that only a theory with such complex configurations can yield a description of the widths and fine structure of the resonances. As far as the calculations of unstable nuclei are concerned it is evident that such calculations need a theory the parameters of which are already known for the (unstable) nuclei under consideration.

Quite recently several variants of the microscopic theory which take into account these improvements microscopically and use non-separable forces have been developed and, what is probably the most difficult for the theory with non-separable forces, numerically realized [2-4] for heavy nuclei. In fact, non-separable forces is a minimal requirement for predictions of properties of unstable nuclei. However these variants have been developed and realized only for 3-8 magic nuclei including unstable ones [4]. The questions arise about hundreds of non-magic nuclei which are being studied in most of experimental works and are of great interest in astrophysics and other fields. There are no such giant resonance calculations in intermediate mass and heavy nuclei.

We have developed an approach for non-magic nuclei [5] which consistently uses the Green function method and generalizes the approach [3] for these nuclei but does it in a simpler form, i.e. in g^2 approx-

imation for the propagators of the integral equations for the transition density (g is the creation amplitude of low-lying collective phonons). Like in the case of magic nuclei there is a justification, at least in the first approximation, to use the known non-separable phenomenological particle-hole interaction of Landau-Migdal type, the parameters of which are already known. Pairing is described by the usual gap equation with the known particle-particle interaction. All the necessary results have been obtained for the case of most practical interest when the dynamic effects of the gap change in an external field can be neglected, but the general case was also treated shortly.

It turned out that the formulae with pairing are not so complicated when compared with the case of magic nuclei. The theory has been also formulated in the coordinate representation. This is important for giant resonances for two reasons. Firstly, it allows to account for the single particle continuum, i.e. the resonance escape width, exactly within the Green function method at least on the RPA level. Secondly, this decreases considerably the numerical difficulties of the problem especially in the case of treatment of the ground state correlations caused by our complex $1p1h$ -phonon configurations. The latter effect was rather noticeable and experimentally observable for magic nuclei especially for ^{40}Ca . Thus, the numerical realization of the approach is quite realistic and it is in progress now.

The approach developed makes it possible to calculate consistently the distribution of giant resonances' strength in non-magic nuclei with taking into account single particle continuum, complex configurations and ground state correlations caused by these configurations. It also allows to account for these effects in the calculations of static properties of nuclei, e.g. the isotopic variations of nuclear radii.

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The discovery of chaos has provided a new understanding of classical dynamical systems. At present it is not clear, however, how chaos manifests itself in quantum mechanics. Since in quantum mechanics trajectories are replaced by energy eigenvalues and eigenstates, it is conjectured that 'quantum chaos' shows up in the fluctuation properties of the spectrum which are characteristic of random matrix theory. In certain simple cases this conjecture has been shown to hold. It is therefore interesting to see whether 'real quantum systems' such as atoms and nuclei show similar signatures. In the case of nuclei, giant resonance excitations are particularly well suited for such studies, because they occur in a region of high level density.

Ignoring particle emission by placing the nucleus in large spherical cavity (closed system) yields a discrete eigenvalue spectrum, E_i . To study the energy distribution let $N(E)$ denote the cumulative level density of the Hamiltonian. It splits up into a smooth part $\bar{N}$ and fluctuations δN ,

$$N(E) = \bar{N}(E) + \delta N(E). \quad (1)$$

To isolate the fluctuating part, $\bar{N}$ is eliminated by a remapping of the energy spectrum,

$$E_i \mapsto e_i \equiv \bar{N}(E_i). \quad (2)$$

which sets the level density everywhere to 1. Define now the nearest neighbour spacing s_i and its corresponding distribution $P(s)$ by

$$s_i \equiv e_{i+1} - e_i, \quad (3)$$

$$P(s) ds \equiv \text{probability for } s \leq s_i \leq s + ds. \quad (4)$$

For a Hamiltonian which corresponds to a classical integrable system $P(s)$ can be shown to be poissonian

$$P(s) = e^{-s}. \quad (5)$$

However, if the Hamiltonian matrix is chosen from the Gaussian Orthogonal Ensemble (GOE) of random matrices, $P(s)$ becomes a Wigner function

$$P(s) = \frac{\pi}{2} s e^{-\frac{\pi}{4}s^2} \quad (6)$$

We have the following conjecture:

For time-reversible closed strongly chaotic systems the matrix of the Hamiltonian operator can be described by the GOE.

Since the Wigner function vanishes for small s , we expect to see linear level repulsion for a chaotic system.

To check these predictions, we start from a nuclear Hamiltonian

$$\hat{H} = \sum_i \epsilon_i a_i^\dagger a_i + \sum_{ijkl} v_{ijkl} a_i^\dagger a_j^\dagger a_l a_k, \quad (7)$$

where the first part denotes a Woods-Saxon mean-field potential and the second a Landau-Migdal zero-range residual interaction. Since a diagonalization of $\hat{H}$ in the full Hilbert space is numerically prohibitive, it is performed in the subspace of 1p-1h and 2p-2h excitations. The nearest-neighbor distribution of the resulting energy spectrum of ^{40}Ca for a fixed angular momentum and parity $J^\pi = 2^+$ fits nicely to a Wigner function (Fig. 1).

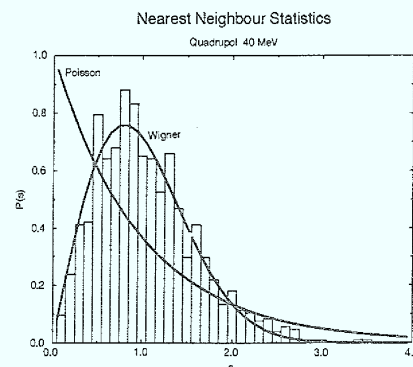


Fig. 1. Nearest-neighbour statistics for $J^\pi = 2^+$ excitation in ^{40}Ca .

A measure for the rigidity of the spectrum is the Δ_3 -statistics of Dyson/Mehta. It measures the least-square deviation of $N(E)$ from its best straight fit in an interval of length L :

$$\Delta_3(L) = \min_{A,B} \frac{1}{L} \int_L d\epsilon [N(E+\epsilon) - A - B\epsilon]^2 \quad (8)$$

Here for an integrable system the prediction is

$$\Delta_3(L) = \frac{L}{15}, \quad (9)$$

and for a matrix from the GOE

$$\Delta_3(L) = \frac{1}{\pi^2} \log L + \text{const.} \quad (10)$$

Applying the Δ_3 -statistics to the spectrum of ^{40}Ca , we again see very good agreement with the GOE prediction (Fig. 2).

Therefore, according to the conjecture the excited states of the ^{40}Ca nucleus are chaotic.

In reality giant resonance excitations lie above the particle emission threshold and the continuum has to be taken into account (open system). To treat these effects we add an antihermitian operator $i\gamma\hat{C}$ to the original Hamiltonian $\hat{H}_0$,

$$\hat{H} = H_0 + i\gamma\hat{C} \quad (11)$$

where γ is a real coupling parameter. In the eigenbasis of H_0 the operator $\hat{C}$ reads

$$C_{kl} = \sum_{c=1}^M V_k^c V_l^c. \quad (12)$$

Each of its components V_k^c is given by an independent Gaussian random number with zero mean and the variance given by

$$\langle (V_k^c)^2 \rangle = \frac{1}{N}, \quad (13)$$

where N denotes the dimension of the truncated Hilbert space of 1p-1h and 2p-2h excitations. Such matrices C_{kl} are known to have $(N-M)$ vanishing eigenvalues. Thus $\hat{H}$ describes a nuclear system with M open channels. Since $\hat{H}$ is non hermitian, its eigenvalues are distributed in the complex plane, rather than on the real axis (Fig. 3).

To examine the nearest-neighbor distribution, the distances to the next neighbor d_i have to be unfolded with the mean level density per area ρ_i

$$s_i = \frac{d_i}{\rho_i}, \quad (14)$$

and renormalized to a mean value equal to 1. For a classically integrable system the level-spacing distribution in the complex plane is given by a poissonian

$$P(s) = \frac{1}{2}\pi s \exp(-\frac{1}{4}\pi s^2), \quad (15)$$

while the classically chaotic case corresponds to the Ginibre ensemble

$$P(s) = \frac{81}{128}\pi^2 s^3 \exp(-\frac{9}{16}\pi s^2). \quad (16)$$

Comparing the spacing distribution of the non-hermitian Hamiltonian with these predictions, we find good agreement with the Ginibre ensemble (Fig. 4).

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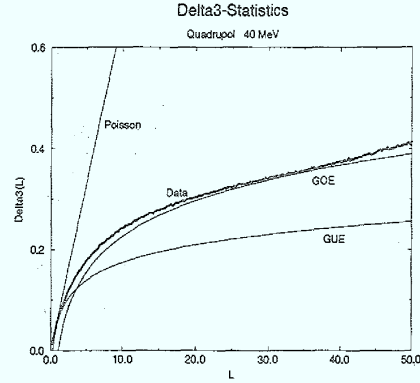


Fig. 2. Δ_3 -statistics for $J^\pi = 2^+$ excitation in ^{40}Ca .

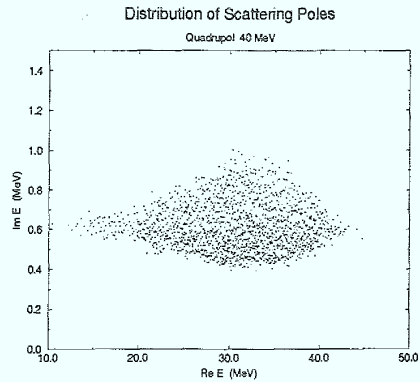


Fig. 3. Distribution of the scattering poles in the complex plane for $J^\pi = 2^+$ and $\gamma = 1.0$.

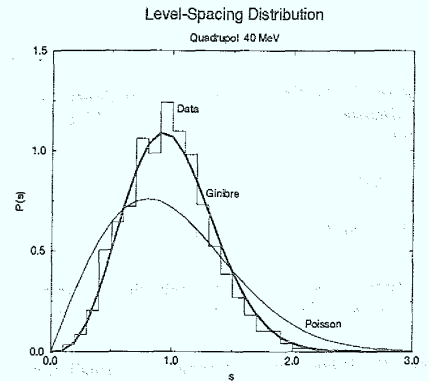


Fig. 4. Nearest neighbour statistics in the complex plane for $J^\pi = 2^+$ and $\gamma = 1.0$.

Double closed shell nuclei form, together with their near neighbors, a few islands in the chart of nuclides where nuclear structure calculations can be expected to be comparatively simple. In particular, RPA calculations within a Woods-Saxon or Hartree-Fock basis usually give reasonable results for most of the low-lying states of the even-particle nuclei in that area. The depletion of the single-particle states below the Fermi surface (and the corresponding partial occupation of those above) can be taken into account approximately by the Landau-Migdal renormalization [1].

To deal with less ideal systems in the spirit of this latter theory, one has to treat the interplay between the even-even and the adjacent even-odd nuclei more explicitly, i. e. one has, in a first step, to calculate explicitly the admixture of collective (two-particle one-hole) components with the single-particle orbitals. It has been shown [2], that this program can, in principle, be carried out by a consistent extension of the Landau-Migdal theory, whereby the appearance of three-body interaction terms is avoided.

Numerical calculations have been carried out [3] for the nuclei ^{88}Sr , ^{90}Zr and their neighbors, where fairly good subshell closure can be assumed.

Recently, new interest arose for the Nickel isotopes [4, 5, 6] with the even mass numbers between 56 and 64.

In this report we consider a few preliminary results for ^{56}Ni and ^{60}Ni . While the former can be considered magic, the latter may possibly have, in a first approximation, a closed neutron $2p_{3/2}$ subshell. On the other hand, ^{56}Ni is certainly far from stability. In fact, a standard RPA calculation with a simple schematic interaction gives already a significant improvement for the energy levels compared to the position of the dominant particle-hole components (Table 1).

Table 1: Experimental and RPA energies (in MeV , calculated in a Woods-Saxon basis with the density-dependent δ -interaction of Ref. [3]) for ^{56}Ni .

	Exp.	RPA
2^+	2.701	2.161
2^+	5.352	6.046
4^+	3.924	3.969
6^+	5.316	6.659

The coupling of the single-particle orbitals to the low-lying particle-hole (ph) phonons, however, does not, at present, lead to stable solutions for energies and single-particle strengths.

In ^{60}Ni , on the other hand, a Woods-Saxon calculation results in a filled neutron $2p_{3/2}$ subshell and its empty spin-orbit partner separated by only 2.1 MeV , compared to an empirical shell spacing of 3.57 MeV . Particle-phonon coupling gives stable solutions for energy shifts and single-particle strengths (Table 2). Remarkably, the shell gaps are widened, in the neutron case nearly to the empirical value.

Table 2: Shifted energies E , energy shifts ΔE and single-particle strengths z of neutron single-particle levels around the Fermi surface of ^{60}Ni .

	$E [\text{MeV}]$	$\Delta E [\text{MeV}]$	z
$1f_{5/2}$	-6.80	1.14	0.67
$2p_{1/2}$	-6.87	1.65	0.57
$1f_{7/2}$	-10.57	0.05	0.81
$1f_{7/2}$	-15.57	-0.53	0.60

For full consistency, particle-particle as well as odd-odd phonons will have to be included additionally. In the ^{56}Ni case this may also show specific effects of the large neutron-proton difference of the chemical potentials in this nucleus.

For ^{60}Ni it is an open question at the moment whether its structure can be described by starting from a closed-subshell picture. The fact that ^{61}Ni has a $3/2^-$ ground state may indicate that we are already in a region transient to superfluidity.

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Coulomb dissociation is by now a well established tool for nuclear structure physics and nuclear astrophysics. Recent reviews are given from a theoretical [1] as well as experimental point of view [2,3].

Apart from nuclear interaction effects (see e.g. Ref. [2]) the study of higher order electromagnetic effects seems most important. "Post-acceleration effects" have been experimentally observed in ^{11}Li dissociation at MSU [4] and in ^{11}Be dissociation at RIKEN [5]. Theoretical studies with a low ξ -expansion [6,7] of the S-matrix (ξ is the adiabaticity parameter) are not in detailed agreement with experiment, but they are consistent with other theoretical approaches, based on coupled channels or direct numerical integration techniques. We conclude that the "post-acceleration" effects are negligible in $^8\text{B} \rightarrow ^7\text{B} + p$ Coulomb dissociation studies at incident energies $E/A \lesssim 20$ MeV. So the ^8B Coulomb dissociation [8,9] can be used to measure the astrophysical S-factor for the $^7\text{Be}(p,\gamma)^8\text{B}$ reaction, which is vital for the solar neutrino problem.

The study of the $^{12}\text{C}(\alpha,\gamma)^{16}\text{O}$ capture reaction - very important for nuclear astrophysics - by means of the ^{16}O Coulomb dissociation poses severe problems [2,3]. Despite of this, there are strategies to overcome these problems, which are mainly related to the high excitation energy [10,11].

As an example from nuclear structure physics, we discuss higher order effects in the $\frac{1}{2}^+ \rightarrow \frac{1}{2}^-$ bound state transition in ^{11}Be . The relevant parameter

$$z = \frac{2Z_X Z_{\text{eff}}^{(1)} e^2}{\hbar v b (\eta_i + \eta_f)} \quad (1)$$

describes the strength of the interaction in terms of the target charge number Z_X , the impact parameter b , the effective dipole charge number $Z_{\text{eff}}^{(1)}$, and velocity v of the projectile, and the "looseness" of the system by means of the constants $\eta_{i(f)} = \frac{\sqrt{2\mu E_{i(f)}}}{\hbar}$ depending on the binding energy $E_{i(f)}$ of the initial (final) bound state. For further details see Ref. [12]. In Fig. 1 the influence of higher order electromagnetic effects as a function of the characteristic parameter z is shown.

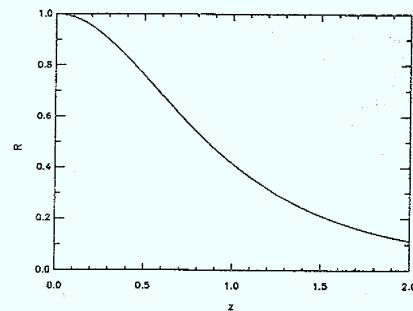


Fig. 1: The ratio $R = |\alpha_{\text{sudden}}|^2 / |\alpha^1|^2$ as a function of z .

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The cross-sections for the photodissociation of a ground state of angular momentum l_f and spin $I_f = l_f + 1/2$ given in [1] are dominated by the radial integrals

$$R_{l_f, l_i, f} \simeq \int_0^\infty dr r^2 g_{l_f}(r) g_{l_i}(r) \quad (1)$$

The radial eigenfunctions g_{l_f}, g_{l_i} have to be computed numerically for realistic nuclear models. Analytic relations are available in some cases (e.g. square well potential). Though the cross-sections are not so handsome, we need therefore easier but accurate approximations. We try to derive them by extending the external integral to the interaction region provide the integrand in eq.(1) does not diverge and the exponential factor does not vary much there. For transitions from s- wave bound states, we use the effective range approach [2] to find the photodissociation cross-section in the electric dipole approximation

$$\sigma^{(\gamma, n)} \simeq \frac{32\pi}{3} \alpha Z_{eff}^2 \frac{\beta}{1 - \beta\tau_0} \left(\frac{k}{k^2 + \beta^2} \right)^3 \quad (2)$$

where $\beta = \sqrt{(2\mu B/\hbar^2)}$, B is the binding energy, $Z_{eff} = \mu(Z_1/M_1 - Z_2/M_2)$ is the effective charge of the clusters, μ is their effective mass and τ_0 the effective range of the interaction.

Eq. (2) fits quite well data from recent Coulomb dissociation experiments of ^{11}Be on ^{208}Pb at 72 A MeV [3]. Fig.1 shows the dipole strength distribution

$$\frac{dB}{dE_x} = \frac{9\hbar c}{16\pi^3} \frac{\sigma^{(\gamma, n)}(E_x)}{E_x} \quad (3)$$

versus the transition energy. We have also applied the

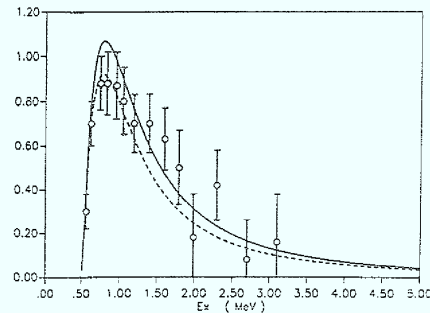


Fig. 1: Dipole strength distribution of ^{11}Be in units of $e^2\text{fm}$ from an effective range approximation (solid line) and a square well model (dashed line), the experimental data are taken from ref.[3]

The results are in good agreement with other theoretical models [4].

This method can be extended to bound states of higher angular momenta, however it will require an empirical normalisation of the cross-sections, for p -waves see e.g. [5]. In conclusion we find that the effective range approximation yields the qualitative behaviour and in many cases the magnitude of the electromagnetic cross-sections at low energy. It is a good starting point in making theoretical predictions and in the analysis of nuclear reaction data at low energies. It allows us to study higher order effects of the interaction in the Coulomb dissociation. Corrections should be found to improve the accuracy.

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III. Accelerator Division

5. COOLER SYNCHROTRON COSY

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The machine development of the COSY cooler synchrotron in the year 1994 was dominated by three topics:

- increase the knowledge of machine parameters
- acceleration of protons to the design momentum 3.3 GeV/c
- extraction of the beam for external experiments with sufficient intensity ($> 10^8$ protons / spill)

Study of machine parameters

Prior to the effort of either approaching the design momentum or increasing the extraction efficiency the overall machine parameters were studied carefully.

Controlling the tune

COSY is operated at a working point of $Q_x \approx 3.6$ and $Q_z \approx 3.6$. The tunes are set by changing the strength of the focussing and defocussing quadrupole families in the bending sections of the ring. The response of the machine to separately changing the focussing or defocussing strength was recorded and is depicted in Fig. 1.

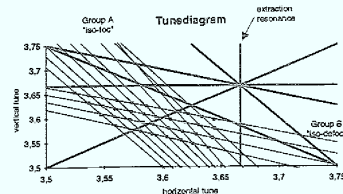


Fig. 1: Measured working points (dots) in COSY at 1.1 GeV/c. On the "iso-defoc"-lines the horizontally defocussing quads are kept constant, and the focussing strength of the focussing ones is increasing from left to right. On the "iso-foc"-lines the horizontally focussing quads are kept constant, and the focussing strength of the defocussing ones is increasing from right to left.

Thus the working point can be driven to any reasonable tune. A tool to predict the necessary quadrupole settings was developed and installed.

Controlling the orbit

Being injected into the ring the protons enter their orbit, which is significantly different from the ideal orbit. With the 57 beam position monitors (BPMs) the protons' orbit was measured at different momenta. Guided by MICADO-calculations kickstrengths for the correctormagnets (steers) were generated and applied. Best results are achieved after fineadjustment (scaling) of all magnetic fields in the order of 10^{-4} . It was verified that the corrections found are valid for momenta up to 1.7 GeV/c. All corrections needed refer to the horizontal plane. In the vertical plane the beam deviates less than 3 mm from the ideal orbit. So no corrections are applied. A typical result is shown in Fig. 2.

Closed Orbit in COSY with respect to the ideal orbit / mm

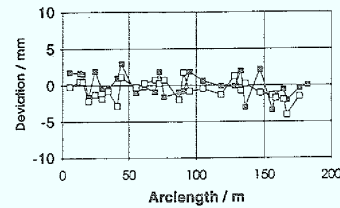
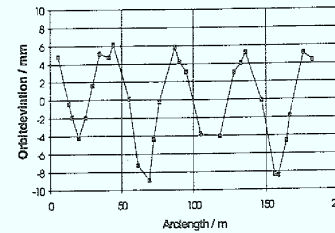


Fig. 2: Orbitdeviation around the ring. Filled squares mark the horizontal Orbit, open squares mark the vertical orbit.

Measurement of dispersion

Orbit for magnetic fields raised by 0.05%



Orbit for magnetic fields lowered by 0.05 %

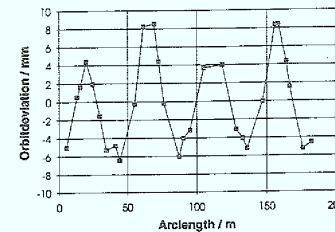


Fig. 3: Dispersionorbits for $\Delta p = 0.5 \%$ and $\Delta p = -0.5 \%$. Both orbits show deviations with opposite sign.

The success of the orbitcorrection was also verified by measuring the orbit with slightly detuned magnetic components. The main dipole and the quadrupoles were scaled to 0.9995 and afterwards to 1.0005 times their respective strength. The measured orbits are displayed in Fig. 3. As expected they show symmetric deviations from the reference orbit, the reference orbit being the same as in Fig. 2.

Acceleration to 3.3 GeV/c (2.5 GeV)

By optimizing the machine parameters step by step we succeeded in accelerating protons up to 1.85 GeV/c. Acceleration beyond this momentum was difficult as we approached the γ_{tr} -regime. By dynamically changing the optic the γ_{tr} -limit could be shifted to momenta higher than the design momentum. This is advantageous especially for the internal experiments. Now data can be taken during the whole ramp cycle. It turned out, that at about the same momentum corrections to the acceleration phase were necessary. On the basis of beam position measurements these corrections were evaluated. They were applied to the first time derivative of the accelerating frequency. The correction used (preset values) is shown in Fig. 4 together with the measurement of the actually applied correction.

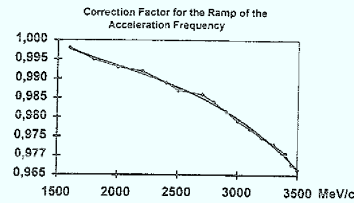


Fig. 4: Modification of the RF-ramp during acceleration. Preset values (smooth line) and measured values (data points) versus momentum

With this correction of the RF-ramp acceleration up to 2.2 GeV/c was possible. Further increase of proton momentum was achieved by compensating for the hysteresis effects of the dipoles magnets. Fig. 5 shows the evaluated and measured corrections to the current ramp of the dipoles. The small "deviation" at about 2.1 GeV/c lies within the measurement accuracy.

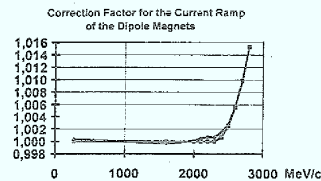


Fig. 5: Modification of the dipole-ramp during acceleration. Preset values (squares) and measured values (triangles) versus momentum

All these efforts resulted in accelerating $2 \cdot 10^{10}$ protons up to 3.3 GeV/c, which is the design momentum. For 3.2 GeV/c and 3.3 GeV/c the intensity of the accelerated beam together with the current in the dipoles is shown in Fig. 6.

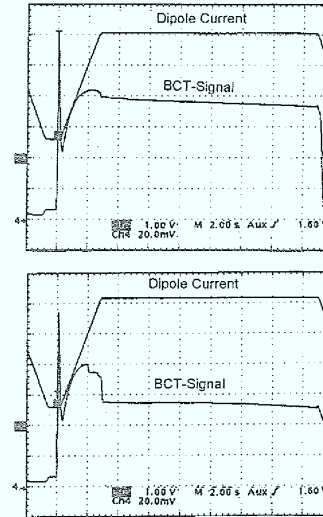


Fig. 6: Acceleration to 3.2 GeV/c (upper figure) and to 3.3 GeV/c (lower figure). The horizontal axis represents the time, full range = 20 seconds. The upper trace shows the current in the bending dipoles: The lower level corresponds to the injection momentum (285 A), the higher level corresponds to the maximum momentum (4200 A). The lower trace shows the current-proportional output voltage of the Beam Current Transformer (BCT).

Increasing extraction efficiency

In 1994 the resonant extraction at COSY could be established, but the extraction efficiency was only in the order of 0.001. Spill times were in the order of 2 seconds. In preparing to increase the extraction efficiency tailored diagnostic tools were installed. For instance the first MultiWire Proportional Chamber (MWPC) in the extraction beam line was equipped with an add on, which provides time dependent observation of beam intensity. As before the horizontal and vertical beamprofiles can be observed at predefinable (trigger) times. Ion optical settings for the first section of the extraction beamline were established to generate an image of the middle of dipole elements (most important: the magnetic septum) on the first or second MWPC.

Inside the ring detectors were installed to localize beam losses. Adiabatic injection into the ring (slow raising of RF-Voltage) increased the number of protons at injection time up to $1.5 \cdot 10^{10}$. Thus good conditions were given to increase the number of extracted protons. For efficient extraction the electrostatic septum has to be the narrowest location inside the ring. Therefore a local orbit bump of appr. 35 mm height was added onto the closed orbit using four steering magnets. The introduced net dipole component which is additionally seen by the beam is corrected for by fine adjusting the strength of the main dipole field in the sub-permille regime. Fig. 7 shows an example of such an orbit bump.

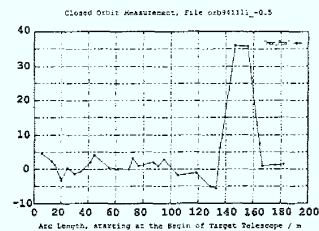


Fig. 7: Measured closed orbit with 35 mm bump at the electrostatic septum

Using the quadrupoles studied before (cf. Fig. 1) and using the optic with the shifted γ_e the quadrupoles are programmed to ramp to a working point with $q_H = 3.64$ and $q_V < q_H$. Subsequently the quadrupoles in the arcs are used for a tune jump towards the $2/3$ - resonance. Next the third order resonance is stimulated by ramping the sextupoles. Thus protons can be fed continuously into the extraction beam line. Up to now this process is so far controlled that more than 10^8 protons are extracted per spill. At 800 MeV/c spill times of about 0.5 seconds could be realized, whereas at 1100 MeV/c 5 seconds were achieved. Extraction at higher momenta is still on the agenda. Components for the ultra-low extraction were installed and their effect on the beam could be verified.

Machine uptime and usage

In Fig. 8 the usage of the synchrotron is summarized. From the scheduled 9532 hours the machine uptime was 8181 hours. 2373 hours were used by experiments. For every month the black column denotes the scheduled machine hours, the white column denotes the hours in which the machine was running, whereas the shaded column shows the machine hours used by experiments.

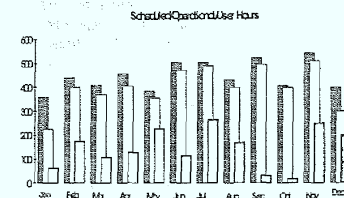


Fig. 8: Summary of machine usage in 1994. Black = scheduled hours, white = operational hours, shaded = hours for experimental use.

1. Introduction

The statistical current fluctuations of coasting beams in storage rings are known as the Schottky noise signals. Schottky spectra are observed via a pick-up electrode in absence of external excitation. They are relatively easy to interpret if one works far enough from the threshold of instabilities. Coherent instabilities are excited by "wake fields": beam particles interact via direct space charge fields and through the electromagnetic fields induced in the walls of the vacuum chamber and the environment surrounding the beam. When a small coherent beam disturbance occurs longitudinal (and transverse) electromagnetic fields are induced, which will perturb the particle trajectory through the Lorentz forces. These fields trail behind the emitting particles and this explains the name "wake field" [1, 2]. The fields, and in turn the beam perturbation, can "self amplify" each other such that instability occurs resulting into beam degradation and loss. The beam environment may be characterised by a longitudinal and transverse "impedance", which is an important quantity for beam stability.

Beam transfer functions (BTF), giving the frequency domain response of the beam to an external excitation, are widely used to measure the impedance of the beam environment and to analyse stability conditions.

Both Schottky and BTF signals have been monitored in COSY. Results obtained from longitudinal measurements will be discussed here.

2. Theory

The longitudinal BTF $r_{\parallel}^0(\omega)$ neglecting the longitudinal coupling impedance $Z_{\parallel}$ is proportional to the following function of the excitation frequency ω and the revolution frequency ω_r

$$r_{\parallel}^0(\omega) \propto j \int \frac{d\psi}{d\omega_r} d\omega_r, \quad (1)$$

where $\psi(\omega_r)$ is the particle's revolution frequency distribution and n the harmonic number [3, 4]. The integral is called the dispersion integral. It can be split into its principal value (PV) and the residue at the pole:

$$\int \frac{d\psi}{d\omega_r} d\omega_r = PV \int \frac{d\psi}{d\omega_r} d\omega_r + j \frac{\pi}{n} \frac{\partial \psi}{\partial \omega_r} \bigg|_{\omega_r = \omega/n} \quad (2)$$

Taking into account collective effects $r_{\parallel}^0(\omega)$ is modified by the shielding function $\varepsilon_{\parallel} = 1 + Z_{\parallel} r_{\parallel}^0$ to [5]

$$r_{\parallel}^c(\omega) = \frac{r_{\parallel}^0(\omega)}{\varepsilon_{\parallel}(\omega)} \quad (3)$$

This influence of the coupling impedance can be seen in the longitudinal stability diagram $1/r_{\parallel}^0(\omega)$ - a plot of the BTF in the complex plane - as a shift by the impedance vector. The instability threshold is reached, when the inverted response curve touches the origin in the stability diagram. If the impedance vector $Z_{\parallel}$ is known (for

example by estimating the shift of the stability diagram), the true momentum distribution $\psi(\omega_r)$ can be evaluated

from the measured $r_{\parallel}^c(\omega)$.

A possibility to check the exactness of the estimated $Z_{\parallel}$ is given by the additional measurement of the deformed longitudinal Schottky power spectrum $P^c(\omega)$. It is related to the non-distorted spectrum $P^0(\omega)$, which reflects the Schottky power spectrum in absence of any coupling effects via

$$P^c(\omega) = \frac{P^0(\omega)}{\|\varepsilon_{\parallel}(\omega)\|^2} \quad \text{with} \quad P^0(\omega) \propto \psi(\omega/n) \quad (4)$$

From the deformed Schottky spectrum one can calculate the distribution $\psi(\omega_r)$ for an estimated vector $Z_{\parallel}$. For the correctly chosen $Z_{\parallel}$ the distributions obtained from the two computations do fit [6, 7, 8].

3. Results

3.1 Longitudinal Schottky signals

When the resonance frequency of the COSY RF system is stepped over the revolution frequency of the circulating particles, the measured total power of the first longitudinal band changes drastically (figure 1). This shows that a detuning of the cavity resonance or a shortening of the cavity gap via a gap relay - both in preparation - will be necessary in order to observe pure beam fluctuation signals, which then will reflect the coasting beam distribution correctly. The total power per Schottky band should be constant and the width of the band should be proportional to the harmonic number. Figure 2 shows the measured total power with an open and (manually) shortened cavity gap at different harmonics [8].

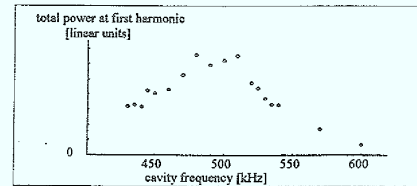


Figure 1: The total power of the beam signal at the first harmonic, measured for different cavity frequencies f_{RF} (revolution frequency $f_0 = 451 \text{ kHz}$).

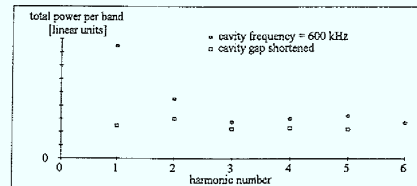


Figure 2: The total power of the Schottky bands at lower harmonic numbers ($f_{RF} = 600 \text{ kHz}$, $f_0 = 451 \text{ kHz}$).

3.2 Longitudinal BTF measurements

A narrow band noise signal has been used for external beam excitation. The frequency response has been measured using a Two Channel Vector Spectrum Analyser. The time gating procedure [9] has been applied to improve the quality of the results.

Due to the low phase space density of the uncooled COSY beam the influence of the coupling impedance was small compared to the amplitude of the inverted response. This is in agreement with equation 2: a Gaussian distribution

$\psi \propto e^{-(\omega - \omega_0)^2 / \Delta\omega^2}$ yields $d\psi/d\omega \propto \Delta\omega^{-2}$ (ω_0 being the average revolution frequency and $\Delta\omega$ the width of the beam distribution). The amplitude of the inverted response of this beam would increase with the square of the beam width.

A calibration of the stability diagram is possible, if the beam intensity is known. A longitudinal stability diagram is shown together with the reconstructed momentum distribution in figure 3 and 4. The measured imaginary part of the coupling impedance of around -1 k Ω per harmonic is in good agreement with the theoretically value expected from the space charge value. The loops in the stability trace are caused by a coupling resonance between the horizontal and vertical machine tune. This problem has been investigated by measuring the transverse BTF [8].

When switching off the driving amplifiers of the COSY RF system, a noise signal is still observable at the cavity gap. Longitudinal BTF measurements have also been performed using this noise signal as excitation (and reference) for the measurement. The measured correlation of the monitored beam signal to the noise at the cavity gap shows that the gap noise has obviously a very strong effect on the coasting beam.

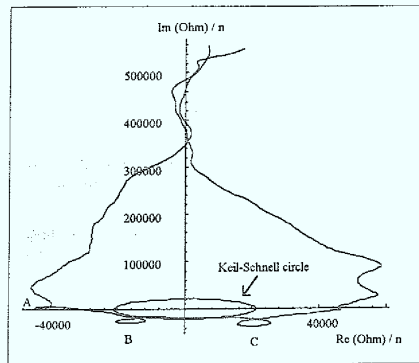


Figure 3: Longitudinal stability diagram of a beam with $\approx 2 \cdot 10^{10}$ protons at 269 MeV/c at the 3rd harmonic, with $(dp/p)_{\text{FWHM}} = 9 \cdot 10^{-4}$. From the stability diagram follows that the coupling impedance is below the theoretical Keil-Schnell limit (ellipse). The loops (marked by "A", "B" and "C") come from tune resonances, as shown by transverse BTF measurements [8].

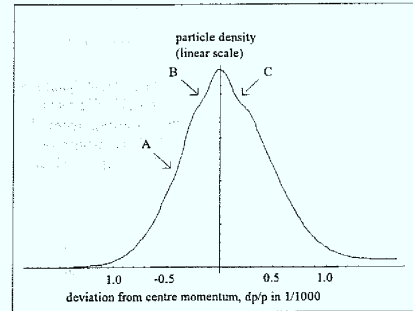


Figure 4: Momentum distribution reconstructed from the stability diagram of figure 3 by integrating the real part of the BTF.

4. Conclusion

In the low frequency region beam signals have been analysed without any external perturbation of the coasting beam (Schottky signals) and in the presence of external excitation signals (giving the BTF). From longitudinal BTF results follows a space charge impedance below the Keil-Schnell limit, as expected in absence of beam cooling.

5. References

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Normally for emittance measurements in particle beams a description by Twiss Parameters is chosen [1]. We used an equivalent but slightly different method starting from the simple law of imaging by a quadrupole lens, which often can be approximated by a thin lens. By this method one gets additional and more detailed information on the intensity distribution in normal and angular space and the possibility to cross check the results.

The envelope of the beam in one plane can be described by three parameters. We chose the distance of the minimum beam radius (waist) from the middle of the magnetic extraction septum of COSY, the beam extension (radius) in the waist x_w , and the beam extension in angle x'_w . These parameters can in principle easily be calculated from the beam radius at three positions z along the central beam. Instead of measuring at three positions we have varied the focal length of a calibrated quadrupole, in this way imaging intensity cuts through the beam at different object distances on one fixed multiwire proportional chamber. From the extension of the image of the beam on this position sensitive device the beam extension at the waist can be calculated.

Near the waist at $z = z_w$ the beam radius x can be described

$$x^2 = x_w^2 + (z - z_w)^2 \cdot x'^2_w \quad [2].$$

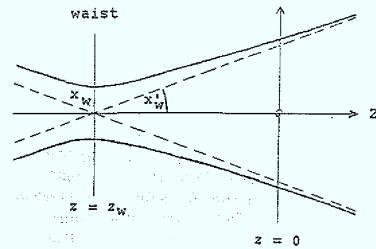


Fig. 1: Beam envelope near the waist. Illustration for the definitions of z_w , x_w , x'_w

The calibration of the quadrupole can easily be obtained with good accuracy by imaging a steerer (in our example SH10) on the position monitor MW12, (if one does not know it from previous experiments anyway). The imaging condition for the middle of the steerer (pivot point) is fulfilled, if the beam profile (especially its center) is constant, when changing the deflection angle in the steerer. Then:

$$1/g + 1/b = c \cdot I$$

where g = object distance = 2.07 m, b = image distance = 3.47 m, c = calibration constant, I = current in the quadrupole Q12. In the imaging case $I = (41 \pm 0.4)\%$ of 185 A at $p = 1.09$ GeV/c. This results in $c = 0.0188$, if I is measured in %. Then the magnification M (corresponding to r_{11} in TRANSPORT notation) for imaging on the position monitor

$$M = 1 - c \cdot b \cdot I = 1 - 0.065 \cdot I[\%]$$

and the object distance from the location where the beam is observed is

$$g = -b/M = -3.47/M.$$

These equations are valid for the horizontal plane. For the vertical plane the sign of I has to be inverted (TRANSPORT). Horizontal and vertical planes can in our case be observed independently as the beam focused in one plane is not defocused in the other plane so much that parts of the beam touch the beam tube.

The results of such measurements of the beam extension are demonstrated in Fig. 2.

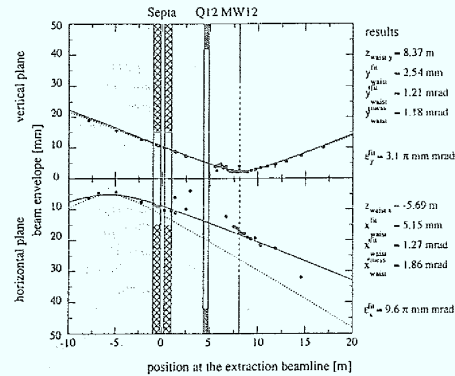


Fig. 2: Measurement of beam extensions in the extraction region of the beam from COSY. Superscript "fit" refers to results using the equation given in the text; superscript "mess" denote results based on angle imaging. These data were used for the ϵ -values given.

The resulting emittance values are

$$\epsilon_x = \pi x_w \cdot x'_w = 9.6 \pi \text{ mm mrad}$$

$$\epsilon_y = \pi y_w \cdot y'_w = 3.1 \pi \text{ mm mrad}.$$

For the 95% width the emittances are

$$\epsilon_x = 28 \pi \text{ mm mrad}$$

$$\epsilon_y = 10 \pi \text{ mm mrad}.$$

These values are still considerably higher than the design value of $\epsilon = 2.5 \pi \text{ mm mrad}$ (x and y) for the beamline to the experimental areas TOF and BIG KARL.

The positions of the waists in three measurements at different runs has not changed much and are mostly about 6 m before the middle of the magnetic septum in the horizontal (x) plane and 9 m after it for the vertical plane (y).

References

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- [2] TRANSPORT Manual (1980), p. 176.

H. Labus, J. Dietrich, N. Dolfus

The principle instrumentation to measure beam profiles in the COSY extraction beamline has already been described in the annual report 1994. We use multi wire proportional chambers (MWPC's), developed at the GSI Darmstadt and commercially available from Physikalische Nachweisgeräte Dr. H. Stelzer. Here we present experiences with these devices, mainly their usability for intensity measurements, and some modifications to improve their availability.

During the design period of the COSY extraction most attention was focussed on a good profile measurement, and no particular effort has been made to get absolute intensity information. Our experience is that we must introduce also precise intensity measurements, which we realize by the following way:

There are two major informations concerning beam intensity, its absolute value and its time structure. Unfortunately the intensity lays in a range between 10^6 and 10^9 /s i.e. 0.1 pA and 0.1 nA. Thus single particle counting is no more applicable, which else would allow absolute and fast intensity measurements simultaneously. Only ionisation chambers with high sensitive current amplifiers are suitable within this range, but have too low cutoff frequencies of some ten Hz to watch large intensity oscillations between DC and to some kHz at the lower intensities. On the other hand they have the great advantage that no intensity calibration is necessary. Absolute intensity may immediately be calculated to a few percent precision with the Bethe Bloch formular if particle energy and amplifier sensitivity are known. Fortunately our MWPC's are possibly the best choice in this situation as they behave like ionisation chambers with intrinsic amplification, which allows to use less sensitive but fast current amplifiers. Profile and fast intensity measurements are therefore possible with one and the same detector system, with the only disadvantage that an intensity calibration must be made at the higher gas amplifications.

The cutoff frequency of our MWPC amplifiers is about 1 kHz and our data acquisition software is able to measure one sample profile every 10 ms, whose area represents the beam intensity. Successive readout of this areas gives a rough information about the time structure of the intensity. In one of our MWPC's we introduced an analogue sum amplifier which adds all 64 currents of one grid for continuously monitoring the intensity. Fig.1 shows the result, with two different time resolutions. From the applied high voltage of 1500 V we can roughly estimate that the maximum intensity is between 10^7 and 10^8 /s, as we have no calibration of the intrinsic gas amplification until now. The intensity shows large fluctuations up to the cutoff frequency, so even larger fluctuation may be present at higher frequencies. In principle the cutoff frequency of the amplifiers could be shifted up to about some 100 kHz, but programmable filters must then be introduced to smooth the statistical fluctuations at the lower end of the intensity region.

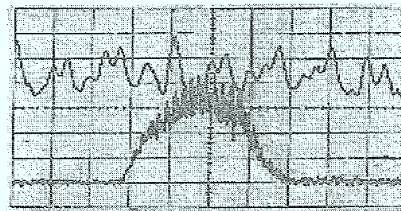


Fig.1: Lower curve: Intensity of one complete spill with 100 ms/unit and max. intensity of about 10^8 /s
Upper curve: Lower curve zoomed with 5ms/unit

We think that it is sufficient to have only one MWPC with continuous analog intensity monitoring at the begin of the extraction beamline. We now prepare the calibration of one MWPC with 40 MeV protons whose intensity will be measured with single event counting photomultipliers up to some 10^9 /s and with faraday cups for the great intensities up to 10^9 /s. The calibration will be reliable if both methods show the same result between 10^6 /s and 10^7 /s, as we try to get the absolute intensity with a precision of only $\pm 30\%$.

To get an online monitoring of the proper function of each MWPC we will attach a weak Sr-90 source in one corner of each MWPC frame which produces a small test profile only with the largest amplification at 2800 V. The area of this test profile indicates the long term stability of all MWPC components. Fig. 2 shows a test profile from a preliminary source at maximum voltage.

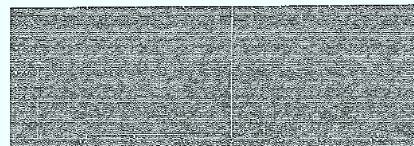


Fig. 2: Test profile from Sr-90 source at 2800 V.
Horizontal: x-profile [0-64 mm] and y-profile [0-64 mm]. Vertical: Intensity approximately 10^7 /s.

The Sr-90 source is situated outside the region of the extracted beam and does not influence measurements above rates $>10^6$ /s.

Each MWPC is supplied by a separate valve to control its gas flow. To reduce gas consumption we now try to operate up to three MWPC's in series with one valve.

First Stochastic Cooling Tank assembled in the COSY Ring

P. Brittner, D. Prasuhn, G. Schug, H. Singer, W. Spieß and R. Stassen

Cryogenic test of tanks

The first horizontal pick-up tank [1-3] has already passed several tests in the vacuum test field. The cryogenically cooled tank reached a temperature of about 30K at the electrode bars and a vacuum pressure of 10^{-9} mbar.

The dilatations caused by thermal shrinks led to a contact between a tuning screw and an inner conductor near a flexible stripline. The location was found using time-domain network analysis. The tank had been partly dismantled to repair this fault. Further similar problems at cooling down did not occur up to now because critical parts of different expansion coefficients have been elastically supported. Especially, the combiner/splitter boards have been embedded in contact springs. The springs also short gaps between the board metallizations and the electrode bars. Densities of around 1 contact per mm have been fabricated using Cu-Be foil of 20µm thickness. The spring surfaces have been prepared for UHV use either by Ag plating of the springs or by hardening of blank pieces under vacuum. The measured outgasing rate was sufficiently low.

Assembly of signal-processing system of the tanks

The earlier described RF components [2,3] have been assembled onto water-cooled alum. plates to provide the thermal symmetry of balanced electronic components of the same kind (s. Fig.1), especially the low-noise preamplifiers. Here, water cooling keeps down the noise factor; a decrease of about 0.02 dB/K. have been measured.

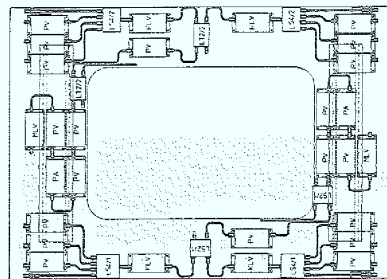


Fig. 1: RF-signal processing plate

The lengths of the coaxial lines interconnecting the RF electronic components have been adjusted to the main energy range of 0.85 to 2.5 GeV. Tab.1 shows the required delays between neighbouring electrode groups, between the centers of the two tanksections of each complete pick-up tank and between pick-up and kicker tanks in the COSY ring. Each branch of the RF-signal path has been measured to correct the sample spread of all components.

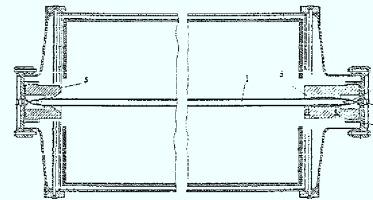
We reached a electrical length uncertainty of about ± 1 mm between the branches over the whole RF band.

kinetic energy /GeV	0,85	1,5	2,5
relative proton velocity	0,851	0,923	0,962
signal length: electrode groups /m	0,282	0,260	0,249
signal length: PU domes /m	2,42	2,23	2,14
signal length: PU-kicker /m	108,05	99,62	95,58

Tab.1: Signal to beam travelling time as vacuum lengths

Beam simulation

An aluminium tube has been centered in the tank to measure the signal behaviour from the center line of the tank to the tank output (s. Fig.2).



- 1: aluminium tube
- 2: vacuum feed-through (N-norm)
- 3: modified CF-flange
- 4: transition element
- 5: PE plastic block

Fig.2: beam simulation under vacuum

Vacuum feed-throughs have enabled RF measurements under vacuum and at cryogenically cooled electrodes bars. The relative positions of the two electrode bars to the aluminium tube were changed by the four stepper motors. The beam simulation corresponds to a relative beam velocity of $v/c_0 = 1$ instead of $v/c_0 = 0.85$ to 0.96 for the real beam. In spite of that small systematical error, we have measured important facts of the signal-processing elements within and outside the tank, like

- RF symmetry of the tank,
- stability of LNAs,
- amplitude and phase fluctuations of the signal-processing channels,
- balance of the channels,
- sum-signal rejection of the channels

Sum and Difference signals

The measurements of transmission from the aluminium tube to the hybrid are represented in Fig.3 for band 1 and Fig.4 for band 2. The isolation between SUM and DIFF output of the hybrid is one of the most important limitations regarding the transverse sensitivity of the system. The insufficient directivity and the large sum signal always pretend an existing diff signal. The diff signal in Fig.3 and Fig.4 represent mostly the isolation of the hybrid and the noise floor amplified by the high gain. The diff-5mm and diff-10mm curves are the corresponding DIFF outputs after moving one electrode bar by 5 (10) mm towards the aluminium tube and the other bar by 5 (10) mm away from the tube.

The transmission curves show that a resolution of 1mm should be possible.

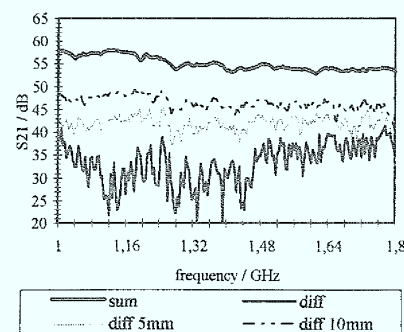


Fig. 3: Transmission measurements, band 1

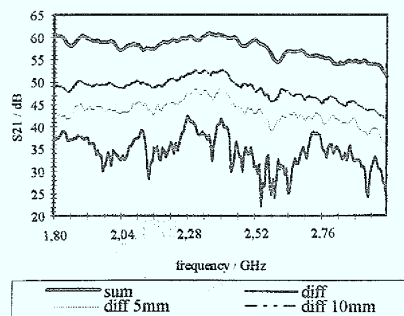


Fig. 4: Transmission measurements, band 2

Behaviour of tank and low-noise preamplifiers (LNAs)

RF characteristics of each group have been proved. The LNAs are the first active parts of the RF signal path [2]. The beam electrodes and the electronic components have been grouped corresponding to $\frac{1}{4}$ -m tank sections.

The LNAs have been connected to the transmission parts of the tank components. RF oscillations occurred in the first HEMT stages of some of the band-2 LNAs greatly affected by the reflection coefficients of the line-type transducers within the tank. The oscillations partly caused a single peak in the RF spectrum, partly changed the DC working conditions although DC uncoupling networks had been interlaced between the LNA inputs and the first HEMT stages. So, changes of the gains by 5 to 10 dB were detected.

The band-2 LNAs now connected to the first PU tank don't show these effects. They have been selected in pairs matched within $\pm 0,3$ dB in order to maintain the sum-signal rejection.

Further, the stability of the LNA input stages is being improved.

Phase behaviour of the diff signal

The phase response from the electrodes up to the hybrid's DIFF output have been measured (Fig.5).

The combining network of the electrodes and the group combining system is optimized for cooling energies between 0.8 and 2.5 GeV [2]. Considering a rel. velocity of 1 in the beam simulation setup the change in phase of the pick-up RF-system is very small (Fig. 5). The diff phases show fluctuations of only $\pm 8^\circ$ in the range of 1.1 to 1.7 GHz.

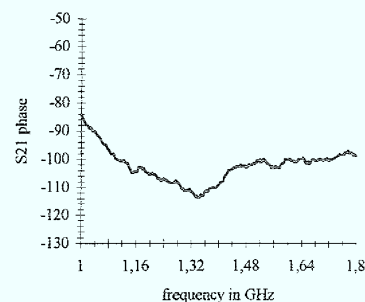


Fig.5: Phase behavior of diff signal

Outlook

All active components of the horizontal pick-up tank are being connected to the COSY-computer system with the aid of the COSY-Control Group. The installation of the first tank into the COSY ring has been performed during the maintenance week before Christmas 1994.

Then it will be possible to take the first RF measurements using the real beam.

Acknowledgments

We like to thank our colleagues of KFA-ZAT and the central workshop as well as the assistants of the other COSY teams.

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During the year 1994 there were two new magnets installed at COSY. One modified quadrupole inside the ring and a small bending magnet right at the beginning of the extraction beamline.

When the extraction of beam from COSY started it turned out that there is a small angular error of the beam behind the magnetic extraction septum. As a quick remedy a spare correction dipole from the extraction beamlines was installed behind the septum magnets to correct for this angular error. With the increase of the extracted momenta, however, it turned out that the bending power of this correction dipole was too weak. Therefore we had to build a new dipole magnet giving a higher field integral. The fastest way to do that was to use the coils of a long correction dipole from the COSY ring. This magnet has an iron length of 227 mm what is more than the effective length of the extraction steerers and a gap height of 180 mm. The extraction beamline requires a gap height of 70 mm only. Therefore making a new iron yoke with a gap height of 70 mm it was possible to gain a factor of 2.6 in bending power. The results from the first experiments show that it now should be possible to correct the angle of the protons extracted up to the maximum momentum of 3.3 GeV/c.

The second new magnet is a modified arc quadrupole. For the internal experiment COSY-11 that will be installed at the beginning of 1995 there was the need to remove a wedge from the sideyoke of quadrupole MQU8. This quadrupole is following the bending magnet MDB8 which is used as spectrometer magnet for the COSY-11 experiment. Machining off a wedge from the sideyoke of the quadrupole gives a large solid angle for the reaction products. To allow for such modifications was one of the reasons for the so called figure of eight design of the COSY quadrupole magnets. In this design the sideyokes are not necessary to carry the magnetic flux of the quadrupole.

As a first step the wedge required by the experiment was machined. This is a wedge in the right sideyoke seen into beam direction. After that the harmonic content of the quadrupole was remeasured. It turned out that at higher excitation when saturation of the iron starts the relative sextupole component is significantly increased. The reason for that is a disturbance of the rotational symmetry by machining only one sideyoke. Therefore as next step the other sideyoke was machined in the same way to keep up rotational symmetry. The harmonic content of the quadrupole was measured again and we found a strong reduction of the sextupole component at high excitation. The values of the sextupole component as well as the sum of the components from $n = 3$ (sextupole) to $n = 10$ (20-pole) before and after machining are given in table 1.

Excitation current	before machining	machining one side	final machining
$n = 3$	$2.6 \cdot 10^{-4}$	$4.9 \cdot 10^{-4}$	$3.1 \cdot 10^{-4}$
$I = 200 \text{ A}$ $\Sigma 3-10$	$2.2 \cdot 10^{-2}$	$2.5 \cdot 10^{-2}$	$2.2 \cdot 10^{-2}$
$n = 3$	$2.6 \cdot 10^{-4}$	$4.1 \cdot 10^{-3}$	$1.1 \cdot 10^{-3}$
$I = 400 \text{ A}$ $\Sigma 3-10$	$2.1 \cdot 10^{-2}$	$2.8 \cdot 10^{-2}$	$2.1 \cdot 10^{-2}$
$n = 3$	$2.0 \cdot 10^{-3}$	$2.2 \cdot 10^{-2}$	$5.2 \cdot 10^{-3}$
$I = 550 \text{ A}$ $\Sigma 3-10$	$2.5 \cdot 10^{-2}$	$5.5 \cdot 10^{-2}$	$2.8 \cdot 10^{-2}$

Table 1: Harmonic content of modified quadrupole MQU8 before and after machining.

Figure 1 shows a sketch of the machined quadrupole MQU8.

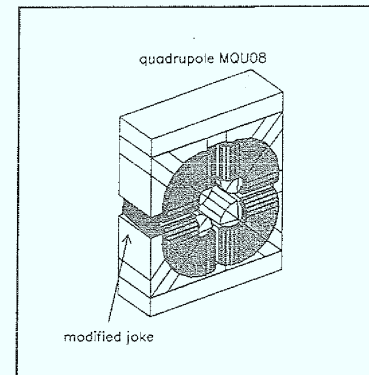


Fig. 1: Quadrupole MQU8 with machined wedge

The modified magnet was installed in the maintenance-week week 31 and aligned to the required position. After the restart of COSY we did not observe any change of the machine performance. A few weeks later COSY reached the design energy with this modified quadrupole magnet.

Proton beam correction scheme

The vertical components of the magnetic field in the COSY electron cooler toroids produce a proton orbit distortion [1]. Fig. 1 shows the proton orbit correction scheme as it is realized. Since the cooler magnets form an upright U, the proton deflection is horizontal and antisymmetric relative to the plane of symmetry in the midst of the drift solenoid. With two steerers at each side of the electron cooler the proton beam can be made coaxial inside the cooler as well as outside beyond the outer steerers so that the electron cooler magnet is "invisible" for the ring. In between the inner and the outer steerers the proton beam passes the compensation solenoid [2] on an oblique trajectory which causes second order vertical deflections. A similar correction using a pair of two steerers is applied which in addition corrects the second order vertical deflections in the toroids and the first order deflections in the dipole coils for the electron beam. Here, in contrast to the horizontal case, the deflection scheme is symmetric. The trajectories shown in Fig. 1 reflect the worst case of maximum angles occurring at injection energy, 37.8 MeV, neglecting the influence of the quadrupoles in between the inner and outer steerers. Since the electron cooler magnetic field is held constant, the horizontal angles are reduced with $1/p$ during acceleration. The vertical angles vanish with $1/p^2$, except the $1/p$

deflection due to the electron beam dipole coil which is being considered for the maximum electron energy, 100 keV. This means that the horizontal steerer field is also held constant, whereas the vertical steerer field has to be controlled according to a function $B(t) = a_1/p(t) - a_2$. The outer horizontal (X) and vertical (Y) steerers are standard components of the COSY ring. The inner steerers had to be designed according to the above correction scheme with maximum deflection angles of 70 mrad horizontally and 5 mrad vertically for 40 MeV p beam energy.

Lay-out and performance

The restricted space in between the cooler toroids and the compensation solenoids had favoured a 0.3 m long, combined X-Y window frame magnet with race-track type coils, Fig. 2. Following a proposition of Halbach [3] the Y coils are split in order to reduce the sextupole component caused by a current distribution not homogeneously filling out the space in between the yokes. Based on a 2-D POISSON calculation maximum magnetic fields of 0.177 T at 285 A for the X steering magnet and 0.0126 T at 40 A for the Y steering magnet were expected. The calculated effective length of 0.35 m in both directions, due to field clamping at both ends, would have guaranteed the required deflection angles at 0.9 Tm beam momentum. Fig. 3 shows

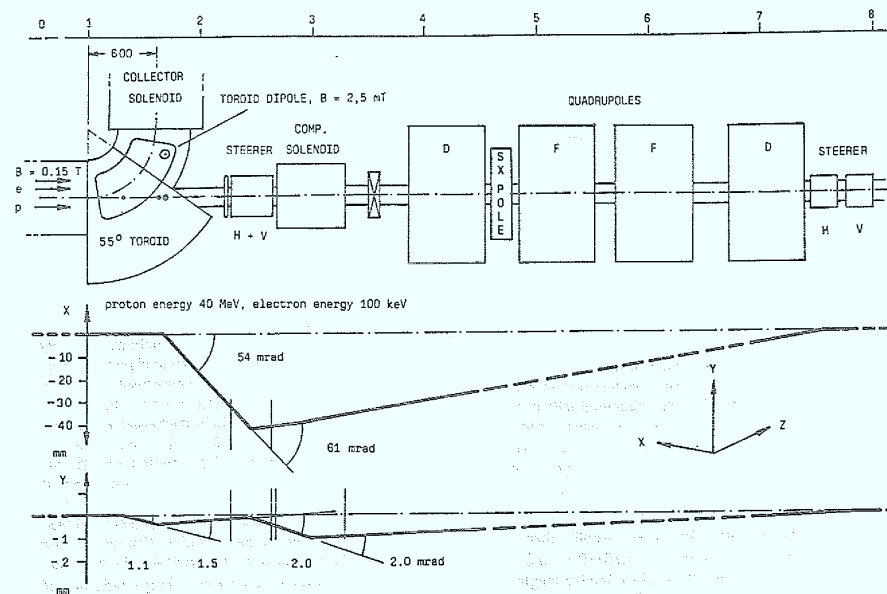


Fig. 1: Proton beam correction scheme for the COSY electron cooler. The angles correspond to the proton injection energy of 38.8 MeV and the maximum electron energy of 100 keV (180 MeV p).

a typical 2-D POISSON field line pattern for the X magnet [4]. Due to the race-track type coils there is an appreciable amount of stray flux. We did not expect saturation in the X magnet since we had checked the field for a double current of 570 A resulting in a maximum B value only 7% lower than linearly extrapolated. The surprise happened when we measured the fabricated device. The X magnet started to saturate already at 100 A reaching only 68% of the design field at 285 A.

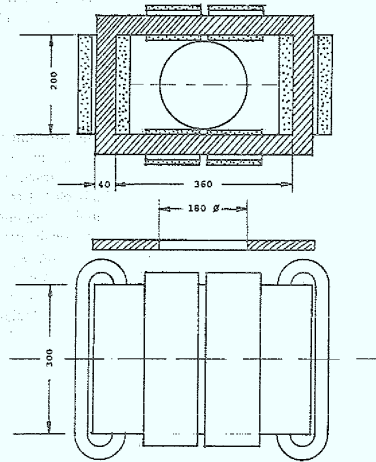


Fig. 2: Mechanical dimensions of the X-Y window frame steering magnet.

Repeating the field calculations with the 3-D computer code PROFI [5], saturation was confirmed showing an excellent coincidence between measured and calculated field, both for the maximum magnetic field, Fig. 4, and the effective length, 0.35 m. The maximum flux density in the X magnet yoke was calculated to be 2.02 T connected with a relative permeability of 20. In all computer calculations the μ values for the real material (swedish dynamo grade iron sheet CK-37) were used.

Discussion

It became obvious that our short window frame magnet with the race-track coils can only be fully understood in a 3-D model. A simplified geometrical argumentation tells us that only the current inside the iron yoke should contribute to the useful flux ϕ_u and all currents outside the iron produce stray flux ϕ_s . In the case of our race-track type coils, Fig. 2, this ratio ϕ_s/ϕ_u is 1:1.85. The useful flux is given by $\phi_u = I_{eff} (w/2) B_{max} = 0.35 (32/2) 0.114 = 0.0064 \text{ Tm}^2$. The flux density in the iron yoke would then be $B_y = \phi_u (1 + \phi_s/\phi_u)/A_y = 0.0064 (1 + 1.85)/(0.3 \cdot 0.04) = 1.5 \text{ T}$, not enough to explain 2.02 T. (I_{eff} effective length, w effective width, B_{max} maximum useful flux density, A_y yoke cross section area). But there are additional field clamps on both front edges acting as magnetic short circuits. From effective length calculations - possible also with a 2-D POISSON procedure [4] - we know that with the given field clamp geometry 26% of otherwise useful flux are lost and

added to the stray flux. The ratio of useful/stray flux is then $(1 - 0.26)/(1.85 + 0.26) = 1:2.85$, resulting in a yoke flux density of $1.5 \cdot 3.85/2.85 = 2.03 \text{ T}$ clearly indicating saturation. The surprising coincidence of this number with the 3-D PROFI result, however, might be accidental.

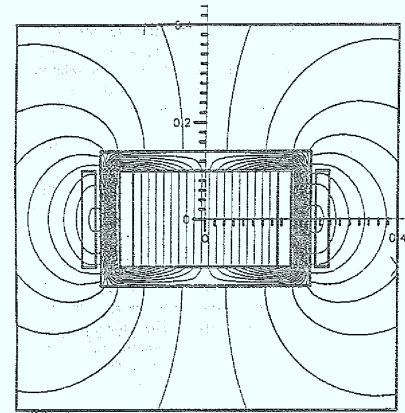


Fig. 3: 2-D POISSON field line pattern for the horizontal (X) steering part of the magnet, 285 A. The ratio of useful/stray flux is independent on the current. It is important to extend the outer field boundary up to about 4 times of the typical dimension of the magnet in order to get a correct result. The ratio of useful/stray flux for this "infinite long" window frame magnet is about 1:1.

Conclusion

The saturation in our window frame magnet is caused by several factors:

- high gap, 200 mm
- high field, 0.18 T
- short length, 300 mm
- race-track type coils
- field clamps on both sides

Saturation could have been foreseen with the above geometrical consideration, but in 1990 nobody of us was aware of the decisive influence of the field clamps. It should be noted that the end plate of the compensation solenoid acts as a natural unavoidable field clamp, Fig. 1. For a proper design under such critical circumstances a reliable 3-D computer calculation like the 3-D PROFI code is indispensable. We have also studied saddle-type coils with the expected result that there is no saturation within the given geometry. If the geometrical edge conditions would permit saddle-type coils, those with longitudinal bendings at the front ends would be the best solution. Our decision now is to increase the thickness of the vertical yoke from 40 to 60 mm making allowance for an even larger ratio of ϕ_s/ϕ_u . The maximum B value is expected to be at least 85% of the original design value and since the effective length does not change this will be sufficient for the operation of the COSY electron cooler solenoid at 0.12 T [1]. By taking away the field clamps at the toroid edge of the X-Y steerer one could increase the effective length, however, one has to check

whether the good homogeneity of the existing magnet, despite its saturation, will be maintained, Fig. 5.

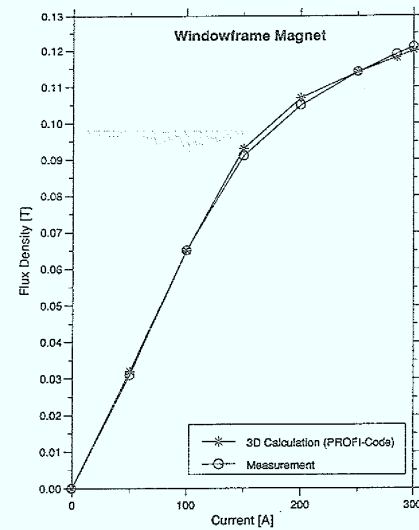


Fig. 4: Comparison of measured and calculated magnetic field B_{max} for the X steering part in the center of the magnet versus current, 108 windings.

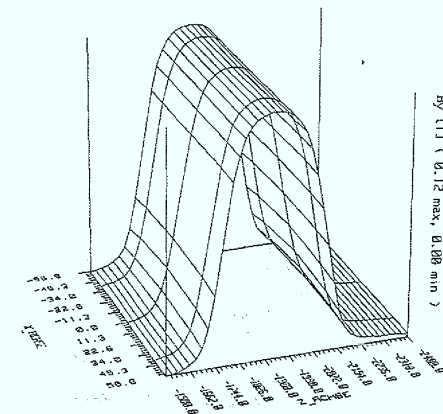


Fig. 5: Measured B field map for the X steering part of the X-Y window frame magnet in the plane $Y = 0$, $X \leq \pm 55$ mm, measured at 280 A. Effective length, 0.35 m, and homogeneity, lower than 10^{-2} , are only little affected by saturation.

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Fast temperature measurement system for the COSY extraction septum magnet

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Technical Requirements

At the position of the electromagnetic septum naturally there are only limited space available. Due to this limited space a very compact design and tied construction of the magnet-system including the magnet coils are essential. As a result of this compact coil design the current densities for our septum magnet are increased up to 80 A/mm² and a power dissipation of 300 kW have to be managed. A very powerful and reliable cooling system is necessary, because of any disturbance of only a part of the cooling system the coil will burn out within 1.2 s. Therefore a very fast temperature control and interlock system is also necessary which can not be setup with standard temperature control components. Due to the construction of several parallel cooling circuits throughout the coil-system any disturbance of only one of the cooling circuits must be detected.

DSP Solution

As a possible solution of a suitable fast temperature control the measuring of current and voltage at the coils are invented and via calculations the resistance of the coil, the resistance temperature coefficient and the temperature are computed. Due to the serial connection of the windings or coils (electrically) and the parallel connection of the cooling circuits the dependence of the resistance deviation by failing of only one cooling circuit is rather low. Therefore not only a fail-safe and very fast but also a very sensitive control system are necessary where current or voltage signal disturbances will be ignored but real temperature deviations must be detected. Furthermore different experimental current settings hat to be accepted.

The hardware solution consists of a Digital Signal Processing system (DSP) which will monitor and control continuously the coil parameter (current and voltage). It performs filter functions (butterworth low pass filter with 1 kHz cutoff frequency; see Fig. 1) and calculates the magnet coil resistance and resistance temperature coefficient as well.

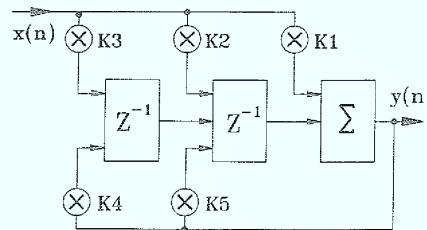


Fig. 1: Butterworth Filter in Z-notation

In a previous performed accurate measurement series the coil parameters of the septum magnet are measured. The calculated resistance and temperature coefficient were saved in DSP memory (ROM). To control the actual temperature behavior the actual measured and calculated values are compared in accordance with the previous performed accurate measurements. The deviation between measured and calculated resistance versus magnet current can be seen in Fig. 2.

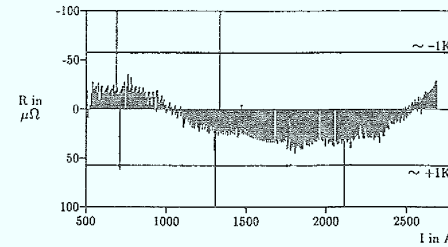


Fig. 2: Deviation between measured and calculated coil resistance

When exceeding a certain threshold value this event is counted. Within a window of 64 measurements the number of counted events must not exceed a programmed number. If this will happen the power supply is switched off. In Fig. 3. the voltage versus temperature are shown (sensitivity of the system).

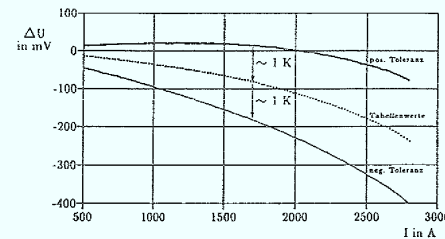


Fig. 3: Voltage deviations and threshold temperature versus current

Moreover due to the COSY extraction magnet where two coils are driven in series the voltage value of both coils can be compared which will give a control function of the symmetrical behavior.

The following overall values, resolutions and criteria were achieved:

resolution of voltage measurement:	$2.5 \cdot 10^{-4} \pm 11.5 \text{ mV}$
resolution of the current measurement:	$5 \cdot 10^{-4} \pm 0.73 \text{ A}$
time of error recognition:	1.3 - 3.4 ms
trip criteria:	1. absolute temperature $\vartheta > 80^\circ\text{C}$
	2. tolerable temperature deviation between both coils $\Delta\vartheta_1 \leq 1\text{K}$
	3. tolerable temperature deviation between calculated and measured value $\Delta\vartheta_2 \leq 10\text{K}$

In case of 40 exceeding events (criteria) within 64 measurements the power supply will trip. The overall tripping time including the power switch reaction time is about 50 ms.

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Introduction

The system for Ultra-Slow-Extraction consists of a pair of digital noise generators [1] and two kicker-structures [2] with the corresponding amplifiers and matching circuits. This system is used for first beam shaping experiments on the proton beam at COSY [3]. The former laboratory system is now optimized for real application.

Signal generation

The bandwidth limited noise is generated with a digital system (see fig. 1) consisting of two numerical oscillators. Their digital outputs are added and fed to a DAC that produces the analog shaped noise. The NCOs are phase-modulated with a complex sequence that is stored in a memory bank. This design has been applied for a patent [4]. A first version of this set-up used 256 Kwords for each NCO, the actual version contains 2 Mwords to store the phase modulation sequence. With a phase modulation data rate of 12.5 Mwords/NCO the frequency resolution of the noise shape is 5.96 Hz. The clock frequency of 50 MHz allows with a 32 bit frequency word an adjustment of the center frequency in 11.6 mHz steps.

A DSP on each memory-board calculates the phase sequence off-line. It is also possible to compute and store sequences on a workstation. The sequence generating the desired noise shape is then transferred via network to the memory bank.

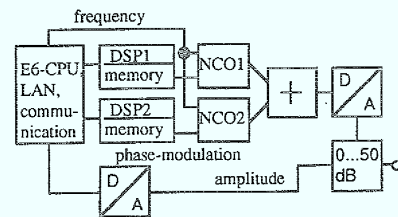


Fig. 1: Layout of the noise generator

The generation of noise is straightforward, although there is a lot of mathematics behind [5]. The pink noise to be generated is given in discrete steps in frequency domain (see fig. 2). This spectrum is treated as a set of oscillators operating at different frequencies with different amplitudes. This information is translated into the two phase modulation-sequences for the NCOs. Both phase sequences contain the phase shift corresponding to the frequency distance of the single oscillator to the center frequency f_0 . The amplitude of this oscillator is then transformed into a phase difference between both NCOs.

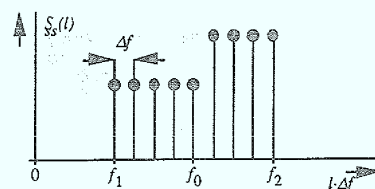


Fig. 2: discrete power spectrum as goal function

In fig. 3a an example of an arbitrarily calculated spectrum is shown and fig. 3b demonstrates the corresponding analog signal, synthesized with the noise generator.

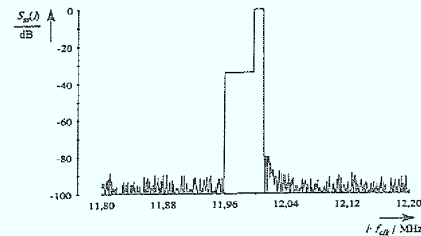


Fig. 3a: Example of a calculated power spectrum

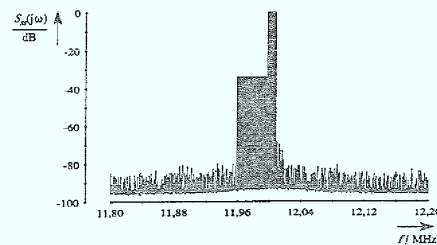


Fig. 3b: the measured spectrum to fig. 3a

As mentioned before there are two noise systems. One operates with a fixed clock frequency of 50 MHz to synthesize a noise at a fixed place in the frequency spectrum. This is used to shape the beam. The second system, with a second amplifier and a second kicker, operates with a variable clock frequency to generate a noise that is swept over a small frequency span. This clock [6] is variable in a range from 40 MHz to 60 MHz, but for normal operation it operated from about 49.9 to 50.1 MHz depending on the desired span.

RF-matching of the USE-Kicker

The noise system is intended to operate on different harmonics of the accelerated beam. The kicker is a band-width limited system, so there are two different modes that are selected with relays (see fig. 4). A wide band mode (see fig. 5a) gives a lower voltage on the USE gap, but allows a frequency range from 2 MHz to 16 MHz. A conventional 50 Ω load resistor acts as a shunt impedance. A narrow band mode (see fig. 5b) produces a higher voltage with less power and works from about 4 MHz to 7 MHz. Therefore a custom made arrangement of resistors was built that gives an impedance of 250 Ω , withstands an rf-power of more than 500 W and does not disturb the rf-characteristic of the kicker structure. The resistor-array is housed in a metallic tube and is air cooled with a fan.

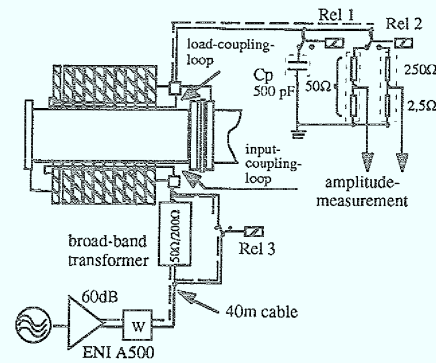


Fig. 4: RF-matching of one Kicker-structure

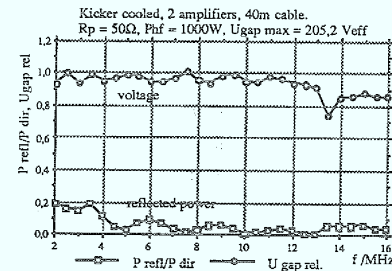


Fig. 5a: Kicker measured in broad-band mode

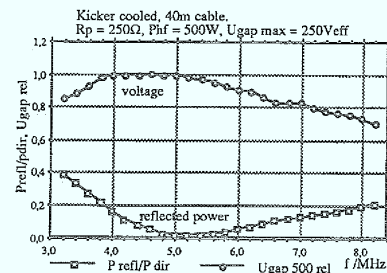


Fig. 5b: Kicker measured in narrow-band mode

First experiments

The experimental set-up for the first experiments [3] is shown in fig. 6 and fig. 7 sketches the involved frequency spectra. Noise generator 1 (DSP1) generates a noise shape with fixed center frequency. The amplitude is controlled with a DAC and an analog multiplier. An attenuator switches this noise signal on or off, depending on the state of a Flip-Flop that is triggered by the COSY-timing system (Timsware 1 and Timsware 2). The second noise generator (DSP2) generates a swept noise. This starts after the fixed noise is switched off. The FFT-spectrum analyzer, a HP 89440A, is triggered with a pulse generator HP 8116A,

to allow several consecutive measurements. This is useful for averaging or to measure the change of the beam distribution as a function of time.

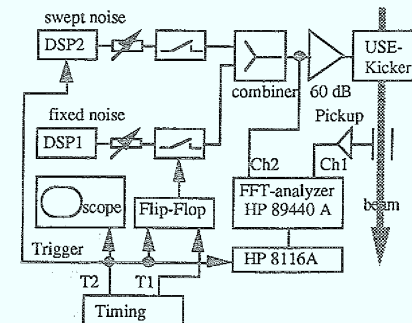


Fig. 6: experimental set-up for first experiments

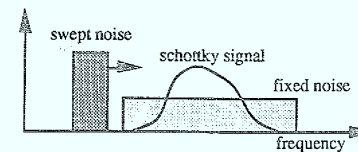


Fig. 7: generated and measured spectra

This system will help to fulfil the requirements for a long duration and constant spill rate of the extracted beam transported to the external experiments at COSY.

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Introduction

In a collaboration [1], [2] between COSY, Jülich and SATURNE, Saclay, we are working on an acceleration station that permits acceleration with a superposition of the fundamental frequency together with higher harmonics at the same time. These non-sinusoidal waveforms are useful for bunch manipulations, acceleration with a flat-top region in the gap-voltage for γ -transition [3], stochastic cooling with a bunched beam, and for faster crossing of γ -transition with a phase jump. Several aspects are different from a cavity operating with a sinusoidal gap voltage. The cavity uses another type of material: VitroVac® instead of ferrites, signal generation is more complex, the phase and amplitude relations between the harmonics have to be compensated and controlled. Finally, tests with high power rf-signal are necessary before such a system can be installed at COSY.

Signal synthesis

We favour a synthesizer that generates the fundamental and at the same time the second and the fourth harmonic.

$$v_{gap} = a_1 \cdot \cos(2\pi f_a t + \phi_1)$$

$$+ a_2 \cdot \cos(4\pi f_a t + \phi_2) + a_4 \cdot \cos(8\pi f_a t + \phi_4)$$

The gap voltage v_{gap} is a fourier series with the amplitudes a_1 , a_2 , and a_4 and the phases ϕ_1 , ϕ_2 , and ϕ_4 .

The "normal" acceleration voltage is described by

$$v(\phi) = v_0 \cdot \sin(\phi)$$

a linearized voltage will be provided by

$$v(\phi) = v_0 \cdot \left(\sin \phi - \frac{5}{32} \sin 2\phi + \frac{1}{256} \sin 4\phi \right)$$

for stochastic cooling with bunched beam one would use

$$v(\phi) = v_0 \cdot \left(\sin \phi - \frac{5}{8} \sin 2\phi + \frac{1}{16} \sin 4\phi \right)$$

and for γ -transition crossing [3] one would apply

$$v(\phi) = v_0 \cdot \left(\cos \phi - \frac{5}{16} \cos 2\phi + \frac{1}{64} \cos 4\phi \right)$$

for about 20 ms. Figs. 1a and 1b show these voltages.

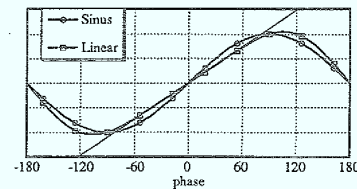


Fig. 1a: sinusoidal and "linearized" voltage

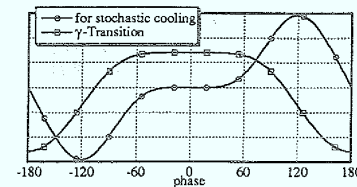


Fig. 1b: voltages for stochastic cooling and γ -transition

The harmonics are generated with a set of oscillators shown in fig. 2. In generating the harmonics with analog oscillators problems arise, because phase shifts, impedance variations in the cavity (a complex load), etc. alter the pulse shape. The problem worsens with a swept frequency. A digital signal synthesis technique (fig. 3) allows frequency ramps preserving the waveform shape, guarantees a phase adjustment of better than 0.02° within a full 360° range, and 16 bit amplitude control for the harmonics. The design bases on a two-stage DDS. An NCO (Numerically Controlled Oscillator) running at 1024 MHz clocks the digital elements. Complex multipliers adjust the complex values delivered by lookup-tables in real time and feed 12 bit DACs. The NCO and the complex coefficients are computer-controlled. The harmonics are added, and the sum drives the power amplifier of the wide band cavity. This setup was the basis for first tests [1], [2].

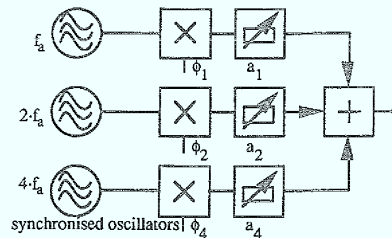


Fig. 2: principle of higher harmonic synthesis

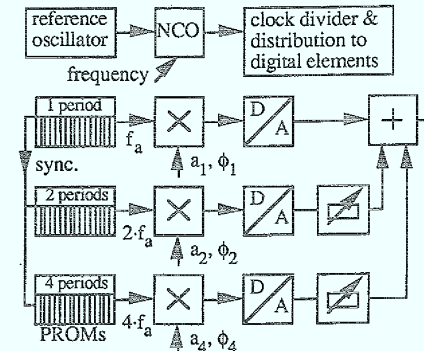


Fig. 3: digital realisation of fig. 2 with synchronized lookup-tables and multipliers for phase-control

Signal synthesis relies on predictable transfer functions of all elements in the signal paths. To minimize errors a fully synchronous design performs the digital processing. The digital clock is an integer multiple of the revolution frequency f_a [6]. A temperature compensated crystal oscillator combined with a NCO (fig. 4) separates the requirements. The TCXO guarantees accuracy and stability over environmental changes and the NCO allows precise tuning.

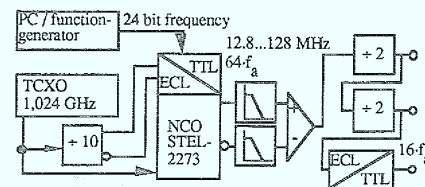


Fig. 4: clock generation with TCXO and NCO

The NCO generates the clock at 64 times the revolution frequency f_a . The clock rate of $32 \cdot f_a$ drives the DACs and $16 \cdot f_a$ all other digital components. This feature is described in more detail in [2] and implies that the algorithm operates frequency-independent.

For tests, the synthesizer was PC-controlled via 96 parallel I/O-lines. The frequency is given with 24 bit, equivalent to a resolution of the fundamental of 0.12 Hz (range: 200 kHz to 2 MHz) and of the fourth harmonic 0.48 Hz (800 kHz to 8 MHz). Amplitude a_h and phase ϕ_h of the harmonic h are normalized to fit into 16 bit two's complement integer numbers: $\{ a_h \cdot \sin(2\pi f_a t \cdot h) ; a_h \cdot \cos(2\pi f_a t \cdot h) \}$.

Behaviour of the VitroVac filled structure

The tuning process of a cavity relies on moving the working conditions on the B-H curve of a non-linear material with a DC-field H_{ext} , demonstrated in fig. 5. The rf-voltage with the field H_{HF} produces sub-loops shown for different values of H_{ext} with tangents a and b. The cavity [1] uses the material VitroVac® 6025F with the characteristic graphed in fig. 6.

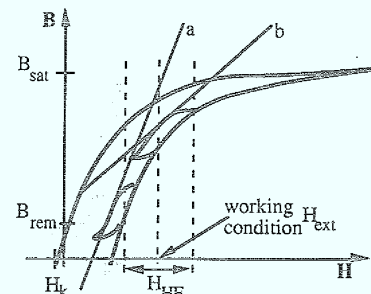


Fig. 5: demonstration of tuning

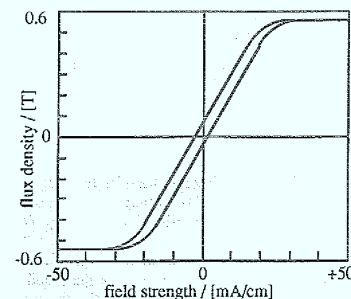


Fig. 6: VitroVac 6025F hysteresis

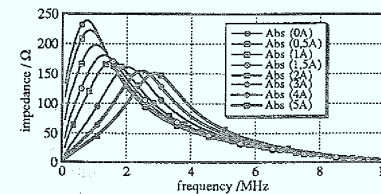


Fig. 7: gap-impedance of the VitroVac®-filled cavity

The impedance pattern of the VitroVac-filled cavity (fig. 7) is broad band and allows a much faster phase jump for γ -transition as shown in fig. 8a compared to the ferrite-filled ($h=1$)-cavity (fig. 8b) installed in COSY. The following table lists parameters of the two different structures:

	(h=1)-cavity	VitroVac-cavity
material	ferrite 8C12, Philips	VitroVac 6025F (VAC)
number	46	24
frequency-range	0.44 ... 2 MHz	0.2 ... 8 MHz
tuning-current	-20 ... +70 A	0 ... +10 A
time duration of a phase-jump	10...50 periods	2 periods
impedance	$\approx 1 \text{ k}\Omega$	$\approx 300 \Omega$
actual applied RF-power	up to 50 kW	up to 10 kW, soon 50 kW
actual maximum RF-amplitude	7100 V peak	2450 V peak
signal-shape	sinusoidal	sinusoidal + harmonics

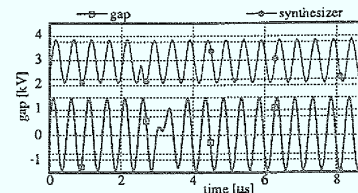


Fig. 8a: phase jump 180° with VitroVac-cavity

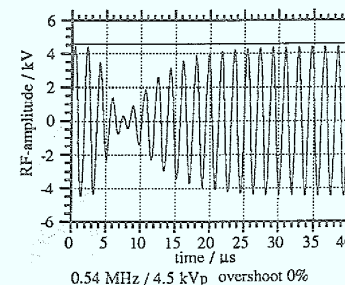


Fig. 8b: phase jump 180° with ferrite-filled ($h=1$)-cavity

Tests with 10 kW rf-power

The VitroVac-cavity was tested to produce higher harmonic waveforms for transition (see fig. 9 [1], and [3]) and for stochastic cooling with a bunched beam (see fig. 10 and [1]). It is important, that the signal at the output of the synthesizer is pre-distorted to compensate the transfer functions of the analog elements. The curve "grid" refers to the signal found on the control grid of the power tubes.

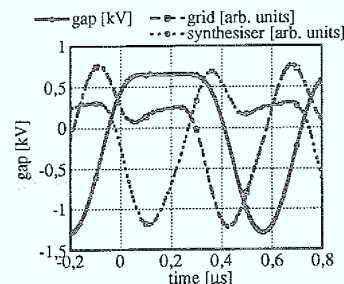


Fig. 9: synthesised voltages for transition

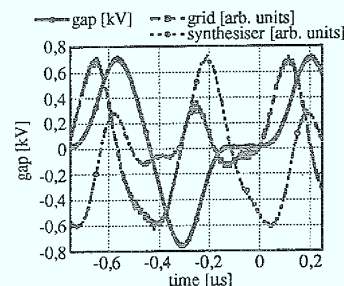


Fig. 10: synthesised voltages for stochastic cooling

To perform these tests a TurboPascal-program controlled all variables and allowed simultaneously ramping of frequency, amplitude, and phase to compensate the influence of the cavity on the signal shape.

Before these higher harmonic signals are applied to the protons parameter studies with a longitudinal program [4] are necessary to calculate the effects on the beam.

Controlling the relations between higher harmonics

Another important aspect is to control and compensate the relations between the fundamental and the higher harmonics. In [6] the synthesizer part and the control loop of COSY are described: the hardware basis. By measuring the phase and amplitude relation of each harmonic component to the fundamental, it should be possible to keep the waveform stable over the full frequency range and to compensate drift effects of the analogue parts. We hope to apply a digital control scheme for optimum feedback philosophy, which is mentioned in more detail in [5] and [7]. Such a feedback-system is now under development and will be based on the digital cavity tuning described in [8].

Conclusions and next steps

A combination of a synthesizer who is able to generate precisely adjusted harmonics in real-time with a broad-band acceleration structure promises interesting acceleration experiments. Some problems remain, regarding the stability of the resonator as a function of temperature and drifts occurring in amplifiers and power supplies.

The maximum voltage at the acceleration gap has to be increased to actual parameters used in COSY. For this reason we have to replace the actual 10 kW tube amplifier with a 50 kW type, almost equal to the one used in the (h=1)-cavity. The mechanical structure was built in the ZAT (KFA), and the tubes and transformers were delivered by Thomson Tubes Electroniques. The assembly is done by the COSY rf-group. We were able to improve the construction of this amplifier, because we drew benefits from the operational experience with the (h=1)-cavity.

It is important to test the VitroVac-cavity in Saclay with 50 kW of sinusoidal signal to guarantee an operation mode comparable to the existing (h=1) cavity. If this test gives positive results, we want to add the higher harmonics.

Another step is to prepare for a test installation at COSY. Involved is a change in the vacuum system and the water cooling system. Power connections are added and control lines have to be installed.

After that, we will try acceleration with this cavity. First with a sinusoidal signal and then with higher harmonics. For the latter, we have to make sure, that we can control the relations between the harmonics in real time. The last step would be to plan the replacement of the existing (h=1)-cavity with the VitroVac-cavity.

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A digital tuning system for the ferrite-loaded coaxial resonator at COSY

F.-J. Etzkorn, S. Papureanu, A. Schnase, H. Stockhorst, and M. Wevelsiep

A ferrite-loaded cavity tuned to the actual revolution frequency [1] in the range from 460 kHz to 1.6 MHz produces the electrical field for acceleration in COSY. The tuning process is described in detail in [2]. The stability of the involved feedback loop is the basis for successful acceleration and is a rewarding object for optimization.

Analog tuning loop

As is pointed out in figs. 1a and 1b with a full cycle from 440 kHz to 1640 kHz and back to 440 kHz the tuning-current of the cavity depends on the frequency, and on the amplitude of the rf-voltage at the acceleration gap, too. The actual frequency is measured and scaled to give a linear and approximate polarisation current following the real value within 10%. A control loop measures the phase difference between the voltages at the acceleration gap and the grid of one power tube and tunes the cavity to a phase difference of 180° so that it behaves like a real load.

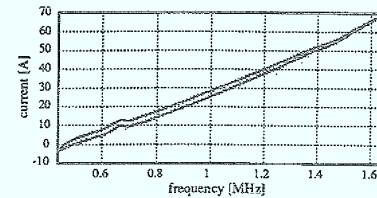


Fig. 1a: Tuning current vs. frequency, "normal" rf-amplitude (DAC-Voltage 4.0 V, COSY-GUI)

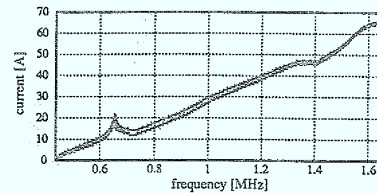


Fig. 1b: "maximum" rf-amplitude (DAC: 9.9 V)

Digital test system

In contrast to the analog type, a digital control loop can be designed to use more knowledge of the process. The measured dependency of the polarisation current on frequency, amplitude and sweep direction (see fig. 1a and 1b) will be stored in a memory to give a better approximation. During each acceleration ramp a digital control loop analyses what has happened to utilise this information for the next ramp. This fits to the operation mode of an accelerator where only small changes from ramp to ramp occur. If the changes are outside a specified interval the DSP assumes that a new ramp is used and the optimization process starts from the initial values stored in memory. The concept of such a digital control loop [3] is verified in a test system shown in fig. 2. The digital phase comparator bases on digital downconversion (fig. 3a) described in [4].

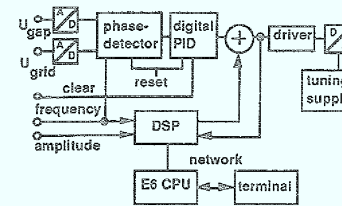


Fig. 2: Digital test system

The circuit operates at 20 MHz with 12 bit resolution ADC and 16 bit processing. The advantages of such a phase-detector are that it uses the full cycle of the input signal (not only the zero-crossings) and operates as a synchronous receiver. So it allows phase measurements at several harmonics, necessary for controlling the cavity for higher harmonics [5].

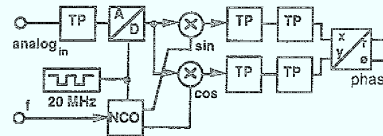


Fig. 3a: Schematic of one digital phase-detector

Two of these circuits and an adder measure the phase-difference between two channels (fig. 3b) in a full 360°-range.

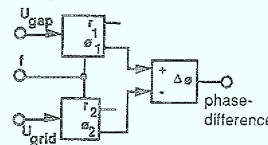


Fig. 3b: Combination of two digital phase-detectors

The digital PID compensation in fig. 2 reacts on the input sequence $e(n)$ with the output sequence $g(n)$:

$$g(n) = P \cdot e(n) + \sum_{i=-\infty}^n e(i) + D \cdot (e(n) - e(n-1))$$

Control inputs serve to reset and freeze the information. The approximation routines for the DSP 32 C (from AT&T, [6]) are written in C on a PC-based development system. The DSP is placed on a VME-Card and communicates via the VME-bus to an E6-CPU. As the DSP is 'slower' than the data-rates delivered from the COSY-function generators, the incoming data is undersampled and the software checks, if the incoming data are consistent. There is a time-out provision that ensures that the last value - for instance in DC mode - is not lost. The driver card is a combination of a 16 bit DAC with Sample & Hold together with optical-coupling to prevent ground loops.

The DSP-program uses the stored function for the first ramp, monitors the behaviour of the PID compensation and consecutively adapts to the following ramps. The stability of this non-linear time-varying system can only be checked with experiments or with simulations, not

with control system theory. The system has to be optimized in a way that small and also big deviations are damped. Important conditions are that the PID compensation does not produce overshoot, and that the stored data points are 'smooth'.

Tests with a cavity model

The ($h=1$)-cavity was not available for experiments, because it was used in production. A test system (fig. 4) simulates the cavity by two original ferrite cores Philips 8C12 which are magnetised by 320 windings for polarisation and 12 turns for the rf-signal. A 100 pF capacitor completes the resonant circuit. With a tuning current in the range from 0 to 8 A the resonance can be shifted from 500 kHz to 1.7 MHz - a frequency-range comparable to the ($h=1$) cavity. The signal synthesizer is the same type used for acceleration in COSY. The function generators are simulated by a PC.

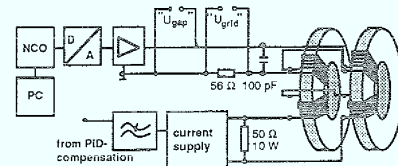


Fig. 4: Test system with simulated cavity

Experiments have shown that there are more than one resonance - this is due to the fact, that the inductors with ferrite cores behave like a complicated system with distributed elements. The current supply had a limited bandwidth (<100 Hz), so that all results can be treated as worst case. The parameters of the PID compensation and of the DSP program were found empirically, a simulation with PSPICE gave only a good starting point.

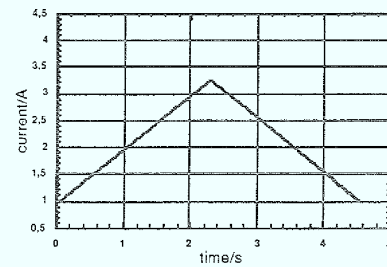


Fig. 5a: estimated current (DSP), first ramp

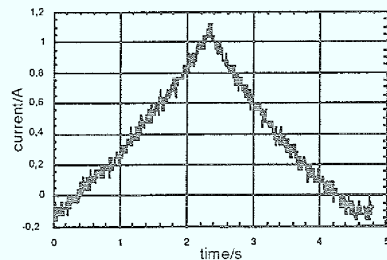


Fig. 5b: output signal PID-compensation, first ramp

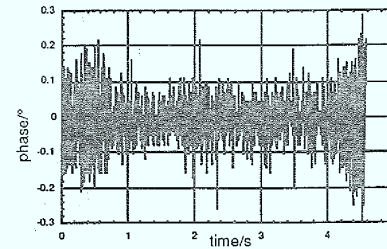


Fig. 5c: phase deviation, first ramp

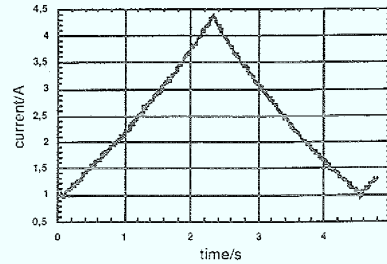


Fig. 5d: estimated current (DSP), consecutive ramps

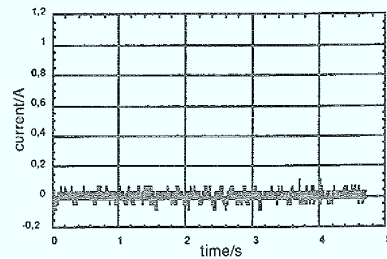


Fig. 5e: output signal PID, consecutive ramps

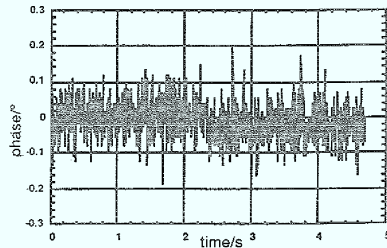


Fig. 5f: phase deviation, consecutive ramps

The diagrams (fig. 5a to 5f) show for the first and the following ramps the current estimated by the DSP, the remaining correction signal of the PID compensator and the resulting phase deviation of the tuning. The frequency sweeps in 2.2 s from 900 kHz linearly to 1.6 MHz and

back again to 900 kHz. It is obvious that the system has 'learned' after the first ramp. The estimated current is more precise and the PID compensator does an easy job. This reduces the resolving phase deviation.

Improvements of the analog ($h=1$)-system

Measurements for the ($h=1$)-cavity parameters have shown that there is some rf-cross-talk between the cables carrying the tuning current and other rf-components. So we added a low-pass filter in the tuning loop, using a VitroVac-toroid and two 100 nF capacitors. Two turns of 25 mm² copper wire yield an inductance of 1.3 mH. Fig. 6 shows the configuration for currents up to 100 A and figs. 7a and 7b show the resulting tuning functions with this filter. The rf-voltage that was present without this filter on the tuning cables (up to ± 50 V depending on the frequency) is now reduced to about ± 200 mV. Additionally a grounded iron tube shields the cables. These experiments have helped to improve the performance of the cavity, and now the digital tuning system will solve the remaining problems with the hysteresis.

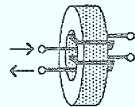


Fig. 6: VitroVac- Toroid 6025F

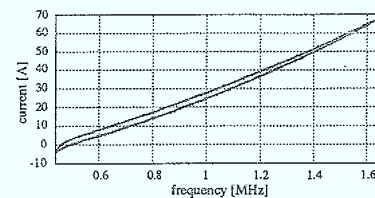


Fig. 7a: DAC: 3.0 V, with filter

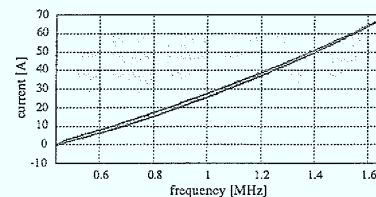


Fig. 7b: DAC: 9.0 V, with filter

Prospects

The results with this type of control loop have to be verified with the real, now improved, ($h=1$)-cavity. When the higher harmonic cavity [5] is tested for application at COSY, we will also need this control loop for stable tuning using non-sinusoidal waveforms.

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Experiments on Longitudinal Beam Shaping Using a Novel Digital Narrow Band Noise Generator

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Abstract

A novel digital rf-noise generator [1,2] was used to excite the COSY-beam longitudinally at a certain harmonic of the revolution frequency. Rectangular shaped noise spectra with bandwidths of the order of 10 kHz and frequency resolution of about 6 Hz were used. The beam distribution was measured during shaping and after switching off the noise source in dependence of bandwidth and amplitude of the rf-noise.

Introduction

The application of specifically shaped noise covers a variety of important investigations in accelerator physics, such as beam transfer measurements [3], sensitive measurements of longitudinal beam aperture, controlled beam heating for electron cooling [4] as well as for ultra slow extraction (USE) [5]. Conventional extraction at COSY will be made by creating a third integer resonant condition by sextupole excitation and sweeping the beam as a whole through the resonance in shifting the horizontal tune of the machine [6]. In this extraction scheme the beam can be considered as a rigid body from which single slices, corresponding to certain momenta, are consecutively taken. Consequently, the momentum of the extracted particles will vary with time. On contrary, the idea behind USE is to move the beam across the resonance by a diffusion process. Firstly, precisely shaped narrow band noise centred on a harmonic of the revolution frequency will be used to transform the initially Gaussian-like distributed beam into a rectangularly shaped distribution. Secondly, to drive the beam towards the resonance, band limited noise, permanently covering the extraction resonance, will be slowly moved across the previously shaped beam distribution. This causes the particles to undertake a random walk within the noise bandwidth thereby filling up any void in longitudinal phase space. Particles reaching the resonance will be extracted from the circulating beam. By properly forming the extraction noise spectra, a diffusion process can be induced which significantly reduces discontinuities in the spill rate, and important as well, will make the average momentum of the extracted particles nearly constant in time with considerably abridged momentum spread. In addition, the speed at which the noise is swept across the shaped beam distribution will control the course spill rate of the extracted particles.

The Digital RF-Noise Generator

Several methods exist to create band limited noise signals. With one exception [7], where the noise spectra are created digitally and a double balanced analog rf mixer is used to convert to the desired revolution harmonic, analog components have been used to build up noise sources and tuneable high quality filters. At COSY a new noise generator [1,2] basing completely on digital signal processing technology has been developed. This offers the possibility to tailor sharply band-limited noise spectra of arbitrary shape, a feature which probably becomes important when considering the transient behaviour of the spill rate. The novel

noise generator will create noise sequences with a frequency resolution down to about 6 Hz and carrier frequencies up to 20 MHz. The bandwidth can be chosen in the wide range from 1 kHz up to 100 kHz. Typical roll-off characteristic of the order of 10 dB/Hz are achieved. Two independent digital noise generators are available. Both allow to control the amplitude of the signal and can be triggered by COSY-Control. One module is capable to deliver swept noise with controlled speed.

Results

Longitudinal beam shaping was studied for debunched beams in flat top of a COSY cycle with approximately one minute flat top length. One noise generator which could be switched on for a predefined time duration was used. After power amplification the output of the noise generator was fed to a longitudinal kicker. The response of the beam to the longitudinal excitation was observed with a beam position monitor of COSY in common mode operation. In addition, the beam current was monitored to adjust the noise bandwidth so that no particle losses could occur during beam heating. Figure 1 to 3 show Schottky scans around the third harmonic of the revolution frequency after 10 s noise excitation for different noise amplitudes with 5 dB/div and 10 kHz span. Note, that the upper lines appearing in the figure are due to interference frequencies and are not affected by the noise.

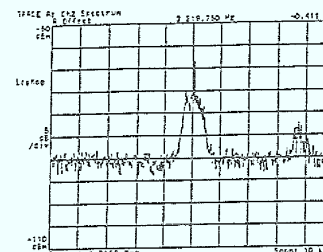


Fig. 1: Schottky scan of the undisturbed beam observed at the 3rd harmonic, 4.219 MHz.

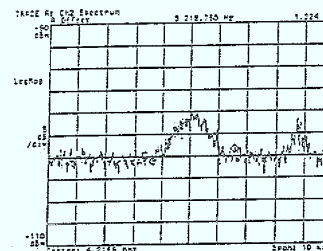


Fig. 2: Schottky scan observed at the 3rd harmonic after shaping with amplitude 2.

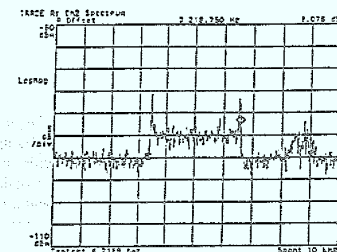


Fig. 3: Schotky scan observed at the 3rd harmonic after shaping with amplitude 8.

The rectangularly shaped noise spectrum of width 4 kHz was centred at the fourth harmonic so that the third harmonic reflects the beam distribution. Figure 1 shows the undisturbed beam spectrum (amplitude 0). The figures demonstrate that the beam shape becomes flattened with increasing noise amplitude. The beam spectrum attains a width of 3 kHz (Fig. 3) as expected according to the equation $\Delta f_m = (m/n) \cdot \Delta f_n$ where n is the harmonic number at which the noise of bandwidth Δf_n is applied and m denotes the harmonic at which the beam is observed. In this case $m = 3$ and $n = 4$. The spectrum in fig. 3 is dog-eared indicating an enhancement of particles near the boundary of the distribution. Random walk simulation show that this is due to the sharp edges of the exciting noise spectrum. This enhancement can be reduced by a corresponding shaping of the noise spectrum. On the other hand the sharp edges could be of advantage for fine longitudinal acceptance measurements. In further tests, noise of constant amplitude but different durations had been applied. From this the expected role that duration times amplitude is constant could be derived. Figure 4 (5 dB/div, span 10 kHz) shows the Schotky scan of the beam distribution at the first harmonic, 1.239 MHz, after 10 s asymmetric noise excitation. The beam was excited at the fourth harmonic, 4.956 MHz, with rectangular shaped noise with a bandwidth of 6 kHz. However, the centre was shifted upwards by 3 kHz.

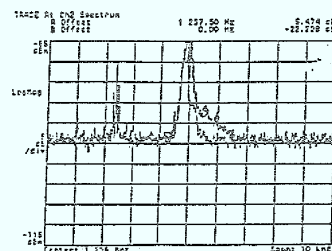


Fig. 4: Schotky scan around the revolution frequency after shaping with 6 kHz noise bandwidth around the 4th harmonic. The noise centre is shifted by +3 kHz.

As expected only particles lying in the right handside of the distribution diffuse into the range of higher momenta. For comparison the figure also contains the spectrum of the

undisturbed beam distribution. The low lying lines in the spectrum correspond to interference frequencies.

Conclusions

The novel narrow-band digital noise generator could be successfully tested for controlled beam shaping. Previous longitudinal beam transfer function measurements [3] using the novel digital rf-synthesizer to excite the beam yielded equivalent results as obtained by applying an analog noise source. This confirms the result following from the current measurements that a frequency resolution of about 6 Hz seems to be sufficient for longitudinal beam shaping applications. Indeed, a controlled diffusion was observed. It was demonstrated that the diffusion coefficient could have been influenced by the amplitude of the applied noise. Depending on the duration of the excitation and due to the sharp cut-off of the noise spectra the resulting beam distributions exhibit a steep slope of the edges superimposed with sharp spikes. Thus, stepping the bandwidth of the excitation noise would allow a fine probing of the longitudinal acceptance of the machine. Future steps have to include improved tailored noise spectra for shaping and swept noise spectra to move the shaped beam towards the resonance. In order to keep shaping time short low noise amplifiers have to be incorporated. The final stage will contain at least two noise spectra of different power density. One which is used for shaping with a fixed carrier frequency, the other one for sweeping across the previously shaped beam. In addition, the application of more powerful noise, centred around the resonance, can serve for reducing ripple in the spill signal.

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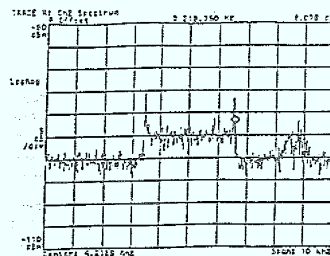


Fig. 3: Schottky scan observed at the 3rd harmonic after shaping with amplitude 8.

The rectangularly shaped noise spectrum of width 4 kHz was centred at the fourth harmonic so that the third harmonic reflects the beam distribution. Figure 1 shows the undisturbed beam spectrum (amplitude 0). The figures demonstrate that the beam shape becomes flattened with increasing noise amplitude. The beam spectrum attains a width of 3 kHz (Fig. 3) as expected according to the equation $\Delta f_m = (m/n) \cdot \Delta f_n$ where n is the harmonic number at which the noise of bandwidth Δf_n is applied and m denotes the harmonic at which the beam is observed. In this case $m = 3$ and $n = 4$. The spectrum in fig. 3 is dog-eared indicating an enhancement of particles near the boundary of the distribution. Random walk simulation show that this is due to the sharp edges of the exciting noise spectrum. This enhancement can be reduced by a corresponding shaping of the noise spectrum. On the other hand the sharp edges could be of advantage for fine longitudinal acceptance measurements. In further tests, noise of constant amplitude but different durations had been applied. From this the expected role that duration times amplitude is constant could be derived. Figure 4 (5 dB/div, span 10 kHz) shows the Schottky scan of the beam distribution at the first harmonic, 1.239 MHz, after 10 s asymmetric noise excitation. The beam was excited at the fourth harmonic, 4.956 MHz, with rectangular shaped noise with a bandwidth of 6 kHz. However, the centre was shifted upwards by 3 kHz.

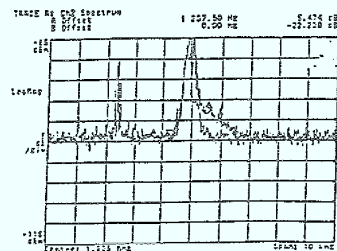


Fig. 4: Schottky scan around the revolution frequency after shaping with 6 kHz noise bandwidth around the 4th harmonic. The noise centre is shifted by +3 kHz.

As expected only particles lying in the right handside of the distribution diffuse into the range of higher momenta. For comparison the figure also contains the spectrum of the

undisturbed beam distribution. The low lying lines in the spectrum correspond to interference frequencies.

Conclusions

The novel narrow-band digital noise generator could be successfully tested for controlled beam shaping. Previous longitudinal beam transfer function measurements [3] using the novel digital rf-synthesizer to excite the beam yielded equivalent results as obtained by applying an analog noise source. This confirms the result following from the current measurements that a frequency resolution of about 6 Hz seems to be sufficient for longitudinal beam shaping applications. Indeed, a controlled diffusion was observed. It was demonstrated that the diffusion coefficient could have been influenced by the amplitude of the applied noise. Depending on the duration of the excitation and due to the sharp cut-off of the noise spectra the resulting beam distributions exhibit a steep slope of the edges superimposed with sharp spikes. Thus, stepping the bandwidth of the excitation noise would allow a fine probing of the longitudinal acceptance of the machine. Future steps have to include improved tailored noise spectra for shaping and swept noise spectra to move the shaped beam towards the resonance. In order to keep shaping time short low noise amplifiers have to be incorporated. The final stage will contain at least two noise spectra of different power density. One which is used for shaping with a fixed carrier frequency, the other one for sweeping across the previously shaped beam. In addition, the application of more powerful noise, centred around the resonance, can serve for reducing ripple in the spill signal.

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The COSY vacuum system is designed for ultra high vacuum operation with a rated base pressure of $1 \cdot 10^{-10}$ hPa in the storage ring. All components of the ring are bakeable to temperatures of 300°C.

After the final assembling of the extraction beam line and the commissioning of COSY in the past year the vacuum system was under operation during the whole year 1994. Exceptions are just some changes in the vacuum system and the EDDA experiment and the installation of the first stochastic cooler tank and one more in line experiment into the COSY beam line. In all these cases we just vented the concerned sections so that the major part of the vacuum system remained under vacuum.

The installation of components as well as changes on the already installed experiments in particular when they are in operation require ventings to air several times per year. In order to operate COSY as much as possible and due to the low vacuum requirements in this early state of maintenance, there was no time spent for bake outs after the ventings to air. The users are satisfied by the vacuum conditions when the system is running under unbaked conditions. The pressure we achieve in this mode is between the low 10^{-9} hPa range and the high 10^{-10} hPa range. This is done by just operating the approximately 70 sputter ion pumps of the COSY ring and leaving the 120 titanium sublimation pumps unemployed.

The status of the vacuum system is permanently controlled by a quadrupole residual gas analyser, ranging up to 10^{-14} hPa per atomic mass unit. Eight analyser heads are served by one control unit in multiplexing mode. Figure 1 shows a typical mass spectrum of the unbaked COSY vacuum system. The total pressure is in the 10^{-10} hPa range. It is achieved after several month of pumping. The dominating hydrogen pressure and the low water vapour pressure indicate a clean system which shows almost the same partial pressure distribution of the residual gas as a baked vacuum system. Hydrocarbon fractions are practically not detectable and far below the specified $5 \cdot 10^{-13}$ hPa per mass unit.

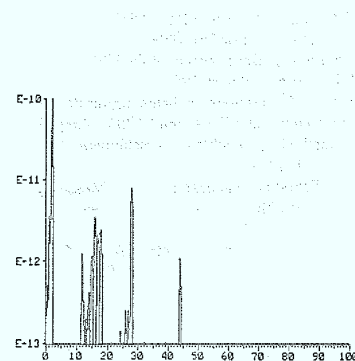


Fig. 1: Residual Gas Spectrum

In 1994 we had no leaks in our system. This success is probably the pay on the effort of giving a final test to absolutely every component exposed to vacuum. We were used to do so during the assembling of the COSY system and we still handle all components of experiments this way.

An additional reason of our successful vacuum maintenance is perhaps the use of the VATSEAL gasket system which is able to resist high bending forces on the flanges of the vacuum chambers. The time saved by having a leak tight vacuum system seems to justify the rather high cost of using this gasket system instead of the common ConFlat system.

The only faults in the operation of the vacuum system we had were caused by lack of electrical supply. But even in these cases the system is still comfortable to handle. Due to pumping only by sputter ion pumps there is no danger of venting the system by error or malfunction and in addition, the degas rates of the vacuum exposed surfaces are so low that still after hours of down time the pressure is below $1 \cdot 10^{-9}$ hPa. This permits to switch the sputter ion pumps back into operation without any roughing.

The COSY vacuum ring is subdivided into eight vacuum sections by gate valves in a manner that each section forms an independent vacuum system. Some of these metal sealed gate valves got small leaks during the last two years probably caused by fibres from the heating jackets of the vacuum pipes. The problem of fibre pollution of the vacuum system is known since the assembling of the vacuum ring but up to now we found no way to avoid this. Fortunately these leaks are just inner leaks and cause no problems during operating COSY.

Up to now we have two in line experiments and the first of four stochastic cooler tanks installed. The attainable vacuum in these areas went down compared with the empty system. This is mainly caused by the large surfaces which are additional installed and the generally higher degas rates of experimental equipment. These areas run unbaked at around $1 \cdot 10^{-9}$ hPa.

At the end of the year a polarimeter tank was installed into the injection beam line. As particles here and in the extraction beam line pass the beam tubes only once, a higher pressure is permitted. So the pressure in the injection beam line as well as the extraction beam line ranges from 10^{-9} hPa at the COSY end to 10^{-7} hPa to the other side.

Alignment Tasks at COSY and the Experimental Areas

U. Bechstedt, W. Ermer, G. Krol, R. Nicia*, D. Rosin, Th., Sagefka

In 1994 most of the alignment tasks were requested by the experimentalists from the external target places. For the Time of Flight Spectrometer (TOF) the rails guiding the different sections of the big vacuum vessel were adjusted in position and height. After that the first elements of the vessel were aligned to bring their axis in coincidence with the axis defined by the last four quadrupole magnets.

For the first beam on TOF in May as well as for the beam-times in September and December there were requests to control and adjust the position of the beam pipe in front of the experiment and to position the targets and the detector arrangements.

To switch the beam between the low energy measurement area (NEMP) and the TOF area the „downstream“ end of the switching dipole has to be moved by approximately 170 mm in radial direction. To allow for that the support of the dipole magnet had to be repositioned and dial gages for the two extreme positions in direction of TOF and NEMP were set.

A lot of work was done for the magnet spectrograph BIG KARL. Besides measurements of the scattering chamber and target positions a complete realignment of all the magnets of the spectrograph was necessary. This was caused by a sag of the support structure under the weight of additional shielding installed after the first alignment.

An additional multi wire proportional chamber (MWPC) was positioned between the last four quadrupole magnets and the scattering chamber of BIG KARL to provide additional information about the focussing of the beam on the target.

The MWPC in the NEMP was adjusted and the grid of fix points was expanded in this experimental area.

For the internal EDDA experiment the newly installed outer detector shell was aligned to the beam axis.

Though the installation of the COSY-11 experiment was delayed to 1995 the measurement concepts for the installation of the vacuum chamber were developed to be ready for the task. In week 31 the modified quadrupole magnet MQU 8 that allows for larger solid angles of the ejectiles from the COSY-11 experiment was installed and realigned.

In preparation of the installation of the H- injection beamline the positions of the vacuum flanges and the support structures for the magnets were measured at the COSY site. After that at the end of the year the components of the H- beamline were installed in the main workshop and measured to their required positions. This allows to check the installation of the modified beam pipe into the magnets before the installation at COSY will start.

Other tasks performed were the measurement of the electrodes inside the spare electrostatic septum and the measurement of the electrodes in the first pick-up tank for stochastic cooling. The position of this tank relative to the COSY-ring was measured after it's installation in the ring in week 51.

For the COSY-13 experiment the positions of targets and guiding structures were measured.

Besides all these activities the group took part in the remeasurement of the COSY level grid and developed new alignment targets and alignment tools.

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A Simulation Program for the particle motion in ELSA and COSY

D. Husmann¹, J. Wenzel¹, R. Maier²

A simulation program for the particle behaviour in the Electron Stretcher Accelerator ELSA of Bonn University and COSY has been developed and successfully tested. It covers physical effects most relevant for the extraction process used at ELSA and COSY. By choosing effective simulation algorithms it is suitable for on-line usage. The simulation program has a graphic user interface and can exchange data with the ELSA control system. First applications of the program give hints to a better understanding of the extraction process and how to optimize it.

1. Requirements

When running accelerator experiments a major amount of time is spent on tuning the machine to the experiment's requirements. To reduce this time an effectively usable control system in combination with diagnostic measurements is necessary. Tuning the machine parameters is always motivated by a physical model describing the particle behaviour. Therefore it is useful to know the model's predictions for diagnostic measurements at the given set of machine parameters when tuning the machine.

To achieve this a simulation is necessary which

- gives results fast enough for on-line usage,
- can read and write the machine's parameters (via the control system),
- displays its results in a form similar to possible results of diagnostic measurements and
- is easy to handle by the user.

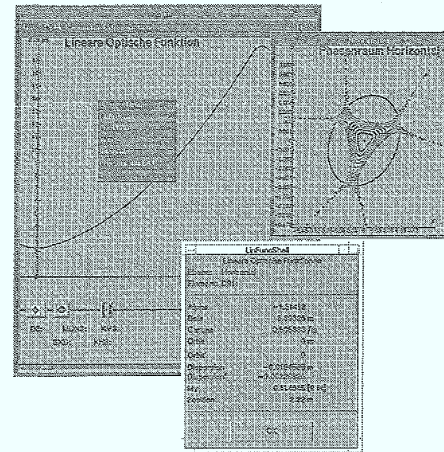
2. Realisation

A simulation program fulfilling these requirements has been developed in combination with the new ELSA control system [1][2]. It is also usable for COSY.

To achieve a quick response-time a special simulation algorithm is necessary which gives analytical solutions for the particle motion. Since the exact equations of motion are not analytically solvable, certain approximations have to be made. This is done within a hamiltonian formalism. Through successive steps of canonical transformations the linear content of the particle motion is extracted. The nonlinear part of the hamiltonian has only major influence on the particle motion in the neighbourhood of resonances. Therefore an ansatz for the hamiltonian which is reasonably exact within the regime of one resonance (e.g. third integer resonance used for extraction) and analytically solvable is used to describe particle motion. Once the ansatz is made the

position (x , p_x) for all particles at all times can be evaluated instantaneously. Further more slow changes in machine parameters are taken into account as adiabatic changes.

Based on the solution of the single particle motion the simulation program can be used to evaluate multi particle effects. Up to now only basic multi particle effects have been implemented, e.g. effects of synchrotron radiation and filamentation due to nonlinearities in the machine.



Abstract

The longitudinal phase space tracking code TIBETANC [1] has now been upgraded with the option of higher rf-harmonics. It thus becomes an auspicious tool to simulate various effects during acceleration of bunched ion beams. It includes the application to space charge and stochastic cooling of bunched beams [2]. In this contribution main emphasis is given to tracking results following a bunch of protons through transition energy. The underlying formalism is presented in detail in [3]. Hardware realizations are exhaustively examined in [4].

Introduction

For certain operation modes of COSY it may be unavoidable to accelerate a bunch of protons across transition energy at which the frequency slip factor η , giving the relative change of the revolution frequency with a relative momentum error, will vanish. In the vicinity of the transition energy, completely determined by lattice functions, the particle motion becomes strongly unstable. I.e. the bunch length becomes very small while the momentum spread becomes very large. In this state the beam is very sensitive against space charge forces, especially for high density beams and a large momentum spread will eventually lead to particle losses. To ensure a stable motion after crossing transition the synchronous rf-phase ϕ_s of the accelerating cavity has to be rapidly changed by $\Delta\phi = \pi - 2\phi_s$. However, conventional rf-phase switching methods suffer from the fact that due to chromatic effects [5], i.e. the dependence of the circumference on the relative momentum spread, particles with different momenta will cross transition at different times and in some circumstances will never cross transition which leads to an increase of particle losses during transition crossing. This makes the timing of phase switching a tagging procedure. Various methods, such as the method of transition energy jump [6] or the "duck under" procedure [5,6] have been developed to overcome this lack and have been studied in detail in [5]. The most challenging procedure is the use of a higher harmonic rf-system. In a neighbourhood of transition one or more harmonics of the fundamental accelerating voltage are switched on for a period of time around transition. The relative phases and amplitudes are chosen in such a way that the resulting voltage will exhibit a broad and flat maximum around the synchronous phase and taking into account for the required energy gain per turn. This will increase the bunch length and confine the momentum spread in an acceptable range. Substantial beam blow up can be drastically reduced.

Double rf-systems have been successfully implemented e.g. at CERN [7] and DESY [8] to cure against space charge blow up. To increase the flat top region of the voltage a variety of technical efforts have been done to develop higher harmonic cavities as is pointed out in [9]. However conventional methods generally require a number of separate cavities which have to be precisely phased in order to achieve the required voltage pattern. To circumvent, a new type of broad band cavity is developed in [9]. In addition, precise phasing becomes available by employing digital

techniques at COSY in synthesizing the higher harmonic voltages [10].

Higher Harmonic RF-Wave Form

Transition crossing is studied using a three octave higher harmonic system, that is, the voltage wave form is composed by the 2nd and 4th harmonic with respect to the fundamental. Since phase focusing at transition is lost [5] a voltage wave which ensures the necessary energy gain per turn and exhibits a broad flattened maximum at the synchronous phase ϕ_s was used. For maximum ramping speed the energy gain per turn in the linear part of the ramp will be $\delta E_s = 1.3$ keV at COSY. Setting the stable phase to be $\phi_s = \pi/2$, maximum flatness will be achieved by the wave form [3]

$$g(\phi) = \sin\phi + \frac{5}{16} \cos(2\phi) + \frac{1}{64} \cos(4\phi) \quad (1)$$

as shown in figure 1. The fundamental peak voltage V_0 is determined by $qV_0 = (64/45) \delta E_s$. With values given above a maximum deviation from the centre value of $< 0.2\%$ for $\Delta\phi = \pm 30^\circ$ is achieved. In this range all particles will receive the same energy gain, provided bunch shearing [3] during transition crossing does not appreciably exceed this value. However, the effect of any deviation can be estimated by tracking using TIBETANC.

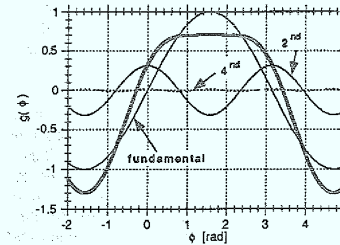


Fig. 1: Single rf-wave forms in comparison to the resultant sum wave form of the three octave rf-system, eq. 1.

Note that the corresponding potential in the Hamiltonian is chair like and will only be focusing for values of $\phi < 60^\circ$. The potential is de-focusing for values above 120° . Similarly, the relative phases and amplitudes of a three octave rf-system suitable for stochastic cooling may be found for given peak voltage V_0 and energy gain δE_s . The potential will then exhibit a broad flattened minimum around the stable phase. Results will be subject to a forthcoming paper.

Conventional Transition Crossing

Phase space plots of transition crossing for maximum ramping speed and top energy of COSY by means of a mere rf-phase jump are shown at three instants, before, at and after transition in figures 2 to 4. A bunch of particles with

an initially parabolic distribution and bunch area $A = 1.09$ eVs is assumed. To fit the bunch well into the bucket the peak voltage of the rf-system was chosen to be 5 kV.

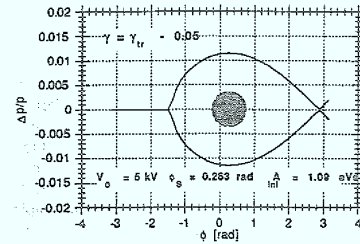


Fig. 2: Bunch and bucket before transition using a single harmonic rf-system.

The phase switching time around transition is 55 μ s. During switching the phase is linearly increased from $\phi_s = 0.263$ rad to $\phi_s = \pi - 0.263$ rad = 2.879 rad. To ensure the necessary energy gain per turn the required peak voltage drops from the initial value 5 kV to 1.3 kV at transition and subsequently increases back to 5 kV. Below transition energy the particles rotate counter-clockwise in longitudinal phase space around the stable phase, $\phi_s = 0.263$ rad, as a result of rf-focusing. In a region close by or at transition, fig. 3, rf-focusing is lost and higher order terms of the frequency slip factor η dominate [3]. Particles with positive momentum deviation become unstable, i.e., they move away from the synchronous phase and form a tail in the direction of positive phases and positive momentum deviations. In the vicinity above transition, now particles with negative momentum deviation will move away from the synchronous phase forming a second tail in positive phase direction and negative momentum deviations. Caused by the non-linearity of the sinusoidal rf-voltage the synchrotron frequency depends on the amplitude of the oscillation. The tails therefore wind up around the core, now clockwise as shown in fig. 4. Thus, even for the fast switching time of the single harmonic rf-cavity at COSY a severe bunch mismatch accompanied with an increase of the effective bunch area may occur. Particles lying outside the momentum acceptance will be ultimately lost. Losses of nearly 50% are to be expected if the momentum acceptance is of the order 5%.

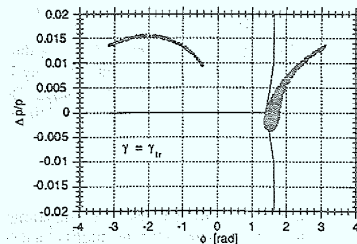


Fig. 3: Bunch and bucket at transition.

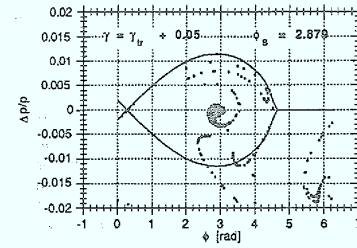


Fig. 4: Bunch and bucket after transition crossing.

Transition Crossing with Two Harmonics

The scenario consists in capturing and accelerating the bunch with a single harmonic system. Just before transition crossing the higher harmonics are switched on for a certain duration. After crossing the bunch will be again accelerated by a single harmonic. With a three octave rf-system (HHS) a bunch of particles with initial parabolic distribution and bunch area $A = 1.09$ eVs was tracked. The initial voltage was now set to 2.5 kV with the corresponding phase $\phi_s = 31.8^\circ$ ($= 0.555$ rad). The bunch with relative momentum spread $\pm 0.2\%$ is shown in figure 5.

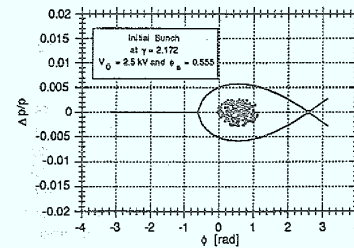


Fig. 5: Bunch and bucket before transition. HHS off.

The bunch is then tracked through transition with the higher harmonic system switched on for different time intervals ± 10 ms and ± 20 ms with respect to the transition crossing time. The necessary peak voltage V_0 is set to 1.87 kV. Bunch and voltage wave form are shown close to transition in figures 6 and 8. Optimum crossing is found for $\Delta t = \pm 10$ ms (fig. 6).

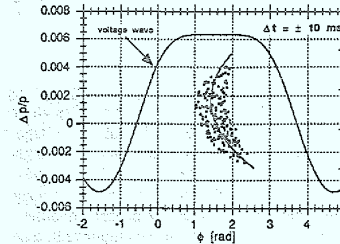


Fig. 6: At transition. The HHS is switched on for 20 ms.

As compared to fig. 4 the relative momentum spread becomes now confined in the interval $[-0.4\%, 0.4\%]$. The crescent bunch shape due to the chromatic effects is visible and is well described by an analytical formula [3] (curve in figure 6).

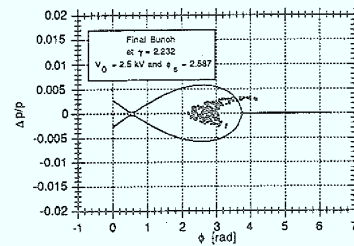


Fig. 7: Bunch after transition. The HHS is switched off and the peak voltage is reset to 2.5 kV.

Also, figure 6 shows that the lower momentum tail of the bunch just enters the decreasing area of the voltage wave form where the voltage is about 2% off the centre value. Particles entering this area are slowly decelerated leading to a small tail in the final distribution after transition crossing, fig. 7. Some of the particles still remain outside the bucket and will eventually be lost. The final relative momentum spread is comparable to that before transition crossing (fig. 5). The effect of a too long on-time is clearly demonstrated in figure 8.

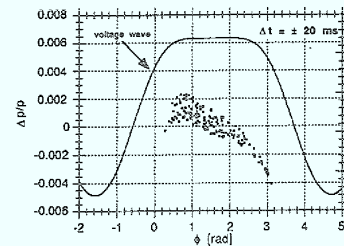


Fig. 8: At transition. The HHS is switched on for 40 ms.

In this case the long low-momentum tail deeply enters the decreasing part of the voltage and particle losses are considerably enlarged.

The figures clearly demonstrate that particle losses and bunch mismatch caused by transition crossing are significantly reduced if a HHS is switched on. Losses can be reduced to nearly 10% depending of the momentum acceptance of the machine.

Conclusions

To conclude, the tracking calculations using the updated version of TIBETANC have shown that transition crossing at COSY by applying a three octave harmonic rf-system would be a challenging procedure. Although the bunch slightly leaves the flat area of the voltage pattern at transition, particle losses are less than 10%. This is to be compared with the conventional transition crossing method

in which particles losses are about 50%. The necessary time interval around transition for the higher harmonic system to work properly is found to be 20 ms. This value is consistent with those derived in [5]. The value is less stringent compared to a fast phase switching necessary in the conventional transition crossing method. Further studies to optimize voltage and chromatic mis-match including space charge can now be done. Additionally, it was mentioned that the same HHS can be used for stochastic cooling of bunched beams [2] as well as for bunch lengthening procedures to cure against space charge effects [7,8], especially of dense cooled beams.

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6. ISOCHRONOUS CYCLOTRON AND ION SOURCES

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For the injection into COSY a well defined beam has to be extracted from the injector over a reasonable period of time. Stability investigations at the cyclotron concerning magnetic field, frequency and amplitude of the accelerating voltage were done simultaneously to the normal operation as injector for COSY.

The NMR based stabilization of the magnetic field, which was implemented in July 1993 [1], operated reliably and proved to be an important precaution for the stable 76 MeV H_2^+ -beam of about 10 μA , which is presently used for the stripping injection into the ring. Long term monitoring showed, that inspite of stabilization small variations of the beam cannot be excluded. Obviously it is not sufficient, to stabilize the magnetic field with respect to the single test point where the NMR detects the field level. Long term drift effects and small changes of iron temperature seem to influence the shape and - as a consequential condition - the average value of the magnetic guiding field.

Since the cyclotron was modified as injector for COSY, a non-interacting diagnostic system is available to monitor continuously phase and intensity of the internal beam at 12 locations [2]. Especially the phase at extraction radius is a very sensitive indicator for the integral influence of magnetic field and frequency of the accelerating voltage (RF instabilities can be neglected in our case). Fig. 1 is a screen copy of the phase detection unit. The upper trace displays the phase, the lower one the intensity of the internal beam over the valley radius between $R=300$ and 1310 mm. Obviously the phase is not well adjusted in this case. The original configuration of the system was modified in one respect: Since not all 12 internal phase probes are necessary for the detection of phase, one electronic channel was connected to a capacitive probe in the transfer beamline to COSY. This allows detection of the external beam current and hence the monitoring of the extraction efficiency.

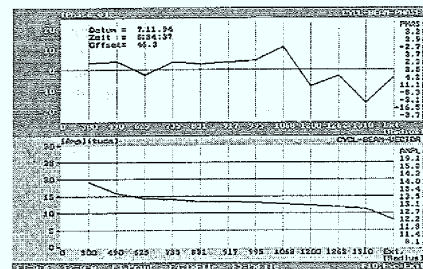


Fig. 1 Display of the phase probe system (see text)

In order to use the phase reading for parameter stabilization, the phase detection system first had to be modified. Originally conceived as a stand alone unit controlled by a dedicated PC, it now had to be equipped with an extension for data exchange in especial with the magnetic field control but also with the newly installed data acquisition system at the cyclotron. The modification was

supported by ZEL (Central Electronics Laboratory of KFA), where the system had been built some years before.

Stabilization now operates in form of superposed loops. The idea is to keep the experimentally optimized situation, which is characterized by a certain phase history, constant over time. We use differences in the reading of the phase probe at extraction radius to calculate changes in the level of the magnetic field required for correction. This new established value is referred to the point where the NMR probe detects the field. The automatic field control system then adjusts the current in an auxiliary main coil winding (6 turns) to achieve the predicted field. Extensive validity checking had to be added to prevent the system from malfunction in case a true phase reading is not possible (no beam, beam chopped in μs -pulses etc.). In such cases the field is stabilized according to the last valid reading. Phase stabilization in the described form, which is in operation since August 1994, contributed remarkably to the stability of the extracted beam.

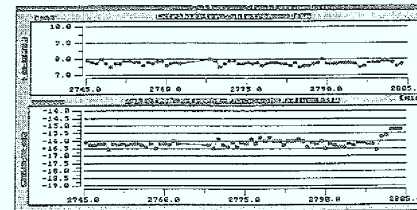


Fig. 2 Display of external beam current (upper trace) and phase (lower trace) over a period of 1 hour

The PC based data acquisition system at cyclotron, ion sources and source beamline was further completed. It combines a number of GPIB IEEE-488 controlled instruments like scanners, DMVs and others with RS-232 connected subsystems like the phase detection system. Dedicated devices like for example beam current instruments are connected via plugged in ADC, DAC and digital I/O cards. The multitasking software VIEWDAC (Keithley Instruments), which is extensively used in the system, allows data acquisition, analysis and display with reasonable performance. The system not only facilitates the operation of the cyclotron but also helps to get information about the long term stability of important parameters and their influence on the beam.

As an example for the effect of phase stabilization Fig. 2 shows an output of the data acquisition system displaying external beam current together with the phase at extraction radius over a period of 1 hour.

* ZEL [Zentrallabor für Elektronik der KFA Jülich]

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W. Bräutigam, G. Blessing*, H.G. Böge, N. Gad, S.M. Qaim*, B. Scholten*

For cross section measurements of charged particle induced nuclear reaction both internal and external beams have been used at the cyclotron. In principle internal irradiations are simple, can be done in a time saving way and should therefore be preferred over using the external beam.

Among the several diagnostic probes available, one main unit can be easily equipped with a variety of probe heads covering the requirements of beam diagnostics as well as internal irradiations. This probe can be moved in and out of the tank by an interlock system without disturbing the vacuum. It can be precisely positioned along the radius from about 150 mm to an extraction radius of 1560 mm, thus providing access to internal beams from very low up to the maximum energy of 45 MeV/nucleon. A possible disadvantage in the use of the internal beam for irradiation is the lack of orbit separation, especially at larger radii where the beam diameter is larger than the orbit separation. This introduces some constraints on target design.

From the irradiation data and the measured radioactivity the cross section of a particular reaction is obtained for one specific projectile energy. In order to determine the excitation function, a large number of irradiations using various probe positions is needed. The work thus involves considerable effort. In contrast, extracted beams can be used with considerable advantage if excitation functions are measured via the stacked-foil technique, which involves a simultaneous irradiation of a series of foils. The projectile energy effective at each foil is obtained by calculation.

At JULIC the internal irradiation equipment can also be used in the extracted beam. Due to the azimuthal position of the main probe unit, which is preceded by the first extraction element, there is access to one separated orbit in

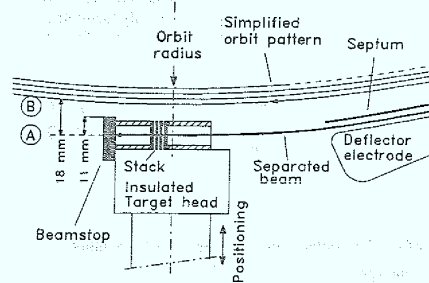


Fig. 1: Simplified representation of the orbit pattern and part of the extraction system. Target head with foil stack in irradiation position

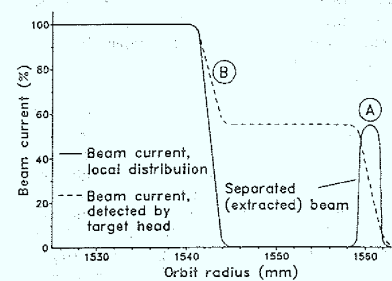


Fig. 2: Location of the internal and the separated (extracted) beam. Positions (A) and (B) refer to the corresponding positions in Fig. 1

the cyclotron. For extraction a part of the internal beam is cut by the septum and then deflected by the electrostatic field in the space between the septum and a high voltage electrode (Fig. 1). This results in an orbit separation of up to 18 mm at the location of the probe.

Since the septum-deflector can be moved radially for optimization purposes, the irradiation position in the extracted beam cannot be defined on a fixed radius scale but has to be found experimentally. Starting at a location in the internal beam, the beam current is detected on the insulated target head (without foil stack) and plotted as a function of radius. This dependence is indicated by the dashed line in Fig. 2. When the probe is moved to larger radius the tip of the beam stop releases the internal beam (position B in Figs. 1 and 2). The reduced external beam then appears according to extraction efficiency and is also detected by the probe head. The probe moving further out finally releases the separated extracted beam and the detected beam current drops to zero (position A). The position of the beam is defined by the slope of the function beam current versus radius. On the basis of known geometries, the probe head with the mounted foil stack can be precisely brought to the appropriate irradiation position.

In preparation on development and production of special isotopes test irradiations were done at the cyclotron in close cooperation with the Institute for Nuclear Chemistry of KFA. The foils to be irradiated are stacked between short Al-tubes which serve as a fixture. An Al-bar finally acts as the beam stop. The whole assembly is mounted on an insulated watercooled target head.

Presently a H_2^+ -beam of 76.4 MeV is used for injection into COSY; so it was also used for the tests. On falling on the first foil of the stack the beam is instantly converted to 38.2 MeV protons of double intensity. The latter then constitutes the primary incident energy on the front foil of the stack.

For beam profile measurements the front foil and another foil in the middle of the stack after irradiation were cut into 1 mm² pieces (in each case a total of about 230 pieces) and the relative activity in each piece was determined. The results are shown in Fig. 3 for the front foil. The various bars depict the intensity distribution of the beam. Evidently, the beam is well focused and centered almost in the middle of the front foil. While passing through the stack the beam is somewhat broadened but is still well suited for cross section measurements.

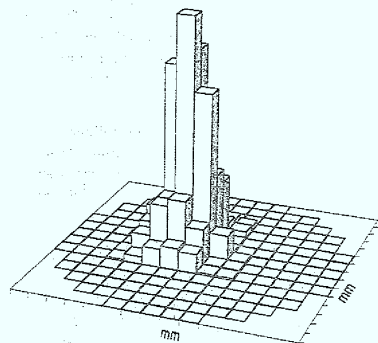


Fig. 3: Beam profile in the front foil determined via radioactivity mapping.

The beam energy degradation within the stack was calculated on the assumption that the beam falls perpendicularly on the stack. For a check of this assumption several Cu monitor foils were placed in different positions within the stack and the irradiation was done at a low beam current of about 52 nA. The following three nuclear reactions, having different reaction thresholds and excitation functions, were chosen for monitoring the beam: $^{63}\text{Cu}(p,n)^{63}\text{Zn}$, $^{63}\text{Cu}(p,2n)^{62}\text{Zn}$ and $^{63}\text{Cu}(p,x)^{58}\text{Co}$. The radioactivity of each product radioisotope was determined via high-resolution γ -ray spectrometry. The decay data were taken from the literature [1]. The proton flux was then calculated using the well-known activation equation. The excitation functions adopted for the calculations taken from [2] are reproduced in Fig. 4(A). The proton flux values derived using the three monitor reactions are shown in Fig. 4(B). Except for some larger deviations in certain energy regions, originating presumably from larger uncertainties in cross sections in those energy regions, the results are generally consistent over the whole investigated energy range of 38 to 19 MeV. The mean flux value of 54 nA (shown by the solid line) is in excellent agreement with the nominal value of 52 nA. The proton energy of 38.2 ± 0.2 MeV incident on the first foil was confirmed by ratio measurements of activities of the monitor reactions [3].

The irradiation facility has been successfully applied in the measurement of $^{124}\text{Te}(p,n)^{124}\text{I}$ and $^{124}\text{Te}(p,2n)^{123}\text{I}$ reaction cross sections over the proton energy range of 20.5 to 31.2 MeV [4]. The smooth transition from the data measured at $E_p < 20$ MeV at the compact cyclotron (CV 28) to the

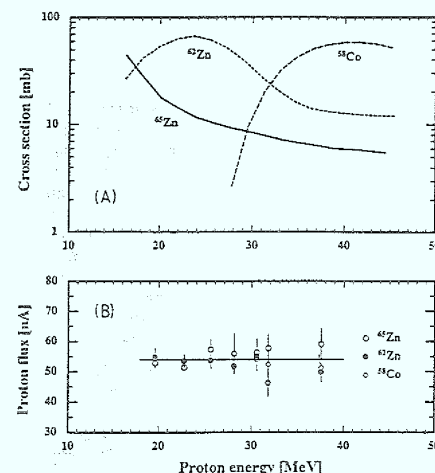


Fig. 4: (A) Excitation functions for the formation of ^{63}Zn , ^{62}Zn and ^{58}Co in proton irradiation of ^{63}Cu , used for beam monitoring in the energy range of 16 to 45 MeV. (B) Proton current values determined via the three monitor reactions at different projectile energies.

higher energy region covered by the present irradiation system, as well as the good agreement between our data and the literature values, validate the various techniques involved in work using the present facility. Cross section measurements on several other medically important radioisotopes, produced via (p,2n) and (p,3n) reactions, are underway.

We conclude that from the viewpoint of both beam profile and beam positioning the irradiation facility described here is very suitable for cross section measurements via the stacked-foil technique.

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Beam injection into COSY requires a pulse structure of the cyclotron beam adjustable in wide ranges. An appropriate pulsing system is installed in the low energy beam line between external ion sources and the cyclotron, which makes use of electrostatic deflection plates. It is triggered by the COSY control system and provides beam pulses from continuous beam down to $1...2 \mu\text{s}$. For COSY machine development, but also for most experiments up now an often used operating mode is characterized by a pulse width of 10ms and a repetition rate of 10s, hence a duty cycle of 1:1000. This results in a net beam time for injection into the ring of less than 90s per day.

So it is obvious to take the partial use of the cyclotron beam for other purposes, i.e. radioisotope production for medical application, into consideration. In this case the beam configuration defined by the injection into COSY has to be accepted by the irradiation experiment. Presently H_2^+ beams of 76 MeV, which is less than the maximum energy from the cyclotron (45 MeV/amu), and intensities in the range of $7...10 \mu\text{A}$ are routinely used. It is planned for a number of reasons to replace this injection scheme by a system for negative ions in the near future. As could be already demonstrated in tests an unpolarized H^- beam of 45 MeV and a beam current of about $10 \mu\text{A}$ is available from the cyclotron [1]. So the basic conditions for the partial use of the cyclotron beam will not be severely influenced by this modification.

An irradiation arrangement has to be provided in the cyclotron vault as a separately shielded area. This is a precaution to use the beam partially, when COSY is in operation and to use it to 100% in times when the ring has to be accessible for the service of internal experiments or machine maintenance.

There are two basic methods to expose a target periodically to the beam: a) use of a relativ fast switching magnet (laminated yoke), to deflect the beam into a separate target station, b) insert the target into the beam mechanically. The first, more flexible alternative would require the replacement of one dipole in the injection beam line inclusive the corresponding power supply. Since it can be accepted to spend a fraction of a second for target positioning, we prefer the mechanical solution. This has the advantage of simplicity and - even more important - does not influence the characteristics of the beam delivered for injection into the ring.

The unit was designed according to the following main specifications: fast alternating motion, vacuum interlock to insert or remove probes without breaking the vacuum, large aperture for flexibility in the shape of probes, watercooled probe head, beam current measurement, visual control of probe under irradiation. The arrangement of the target and target support in Fig. 1 is an exemplary solution only and can be easily adapted to specific needs. In the shown case a tube with built-in watercooling is attached to an insulated bushing by a clamping sleeve. Beam current monitoring is performed by connecting the tube to an appropriate integrating instrument. The end cap is fixed in place only by atmospheric pressure. This design allows the dismounting of an irradiated target from a safe distance with manipulating tools. After removal of two screws the tube with the target is pulled out and the valve can be closed. Then after ventilation the end cap with the whole target arrangement can be taken away to a shielded place for further treatment.

As appears from Fig. 1 angular motion was chosen in preference over lateral positioning. This minimizes the deformation in the vacuum bellow for increased lifetime which is a precaution for continuous operation. The angular motion, adjustable up to $\pm 10^\circ$, is achieved by a crank mechanism. It is driven by a DC-motor with a flat rotor (lamellar conductors). This device with a very low moment of inertia moves the target into and out of the beam in less than 100ms so that during COSY operation up to 98% of the beam time can be used for target irradiation.

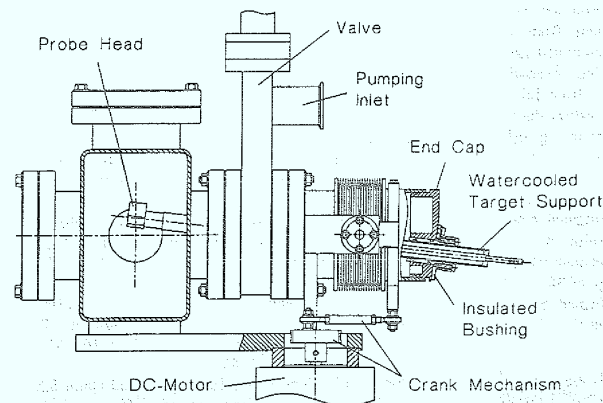


Fig. 1: Irradiation system for the partial use of the cyclotron beam

[1] W. Bräutigam, H. Beuscher, H.G. Böge, R. Brings, et al. Acceleration of Negative Ions in JULIC, IKP Annual Report 1993, p. 204

Optimizing the H^- Production for the COSY Operation

H. Beuscher, W. Hoffmeister, A. Müller

H^- Source

Studies of the performance of the H^- source [1] have been made to optimize the source for a future continuous operation of COSY with stripping injection starting from negative particles.

One disadvantage of the multi-cusp H^- source is the limited lifetime of the filament. To increase this lifetime the arc voltage has been lowered from normally 150V to 80V. This reduced the sputtering of the wire material remarkably and the lifetime increased by about a factor of 2 (two weeks instead of one). With the lower discharge voltage still enough H^- current could be extracted.

Additionally slow pulsing should further extend the lifetime of the wire. Since the injection time into COSY is only about 10ms whereas the whole cycle of the accelerator takes several seconds, the current from the source is only needed for less than 1% of the time. Source pulsing has been tested by pulsing the arc voltage. Using an on-time of 20ms 80% of the DC H^- current could be achieved in pulsed operation. A system for continuous pulsed operation is in preparation.

The output H^- current can be increased without raising the arc voltage by adding a second gas. The addition of 10% N_2 increases the H^- current by 30% and of 10% Ar by 45%. This is a possibility to achieve higher currents, if necessary, without raising the discharge power. It is expected that the second gas improves the H^- production by reducing the recombination of hydrogen atoms.

Emittance measurements have been performed for the H^- source, since this parameter mainly determines the percentage of the beam intensity, which can be accepted by the cyclotron. For the following source settings, which produce the "COSY beam", an emittance of 50mm mrad has been determined: 6mm diameter of the extraction hole, 4.5kV extraction voltage and an extracted current of 250 μ A. The expected dependence of the emittance on the extraction hole radius and the extraction voltage was found, what proves the reliability of the results. The measurements have been performed with a wire-slit system and the results have been verified for one case by measuring beam size and angle directly. A slight increase of the emittance with increasing current and no dependence of the emittance on the arc voltage have been observed. The results are in agreement with the data given by the IBA company, which built the source.

Before the test operation of COSY with H^- particles will start the H^- source shall be set up near the actual position of the H_2^+ source but connected to the beamline via a 25° bending magnet (Fig.1). Then it is possible to switch from the H^- source to the polarized H^- source without "breaking the vacuum". The quick change from one source to the other is especially important when testing the performance of the low current polarized source with the cyclotron and the COSY ring. The former location of the ECR source in the upper hall is still connected to the cyclotron, but will

mainly be used for source testing. A source in this location has a 90° analysing system available and can run without interfering with the source, which is continuously running for COSY. The 25° bending system has been tested at the test source location simulating the optical conditions of the final source position in the basement. The necessary optical matching could be achieved without adding further magnetic elements.

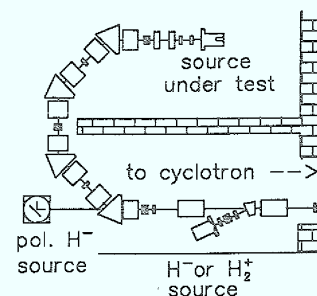


Fig. 1: New source locations at the cyclotron beamline

Microwave H^- Source

To overcome the disadvantage of the filament in the multi-cusp H^- source a microwave source based on our H_2^+ source [2] has been tested. Radio frequency H^- sources are known, which produce H^- currents in the mA range [3]. But they need a lot of RF power and always also contain destructible parts. It was not the idea to build such a source, but to look for the performance of an ECR based type of source with low microwave power injected and without destructible parts. The tests have been performed in the housing of the multi-cusp H^- source by replacing the filament by a 2.45GHz microwave generator. The source had a magnetic multiplet and a magnetic filter but no magnetic solenoid. To start the plasma a small piece of wire has been installed, which in the beginning absorbs the microwave power and starts the plasma by providing electrons. After the plasma has started the microwave power is absorbed in the plasma and the wire stops glowing. The largest H^- current extracted from this type of source was 60 μ A. But this current could only be achieved with 10% of N_2 as support gas. The result is, that this kind of simple source is not able to produce sufficient H^- current (>150 μ A) for our purpose.

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The polarized ion-source for COSY has been built and set in operation by a collaboration of three groups from the universities Bonn, Erlangen, and Köln (cf. Ann. Rep. 1993). The source of the colliding-beams type (CBS) is designed to deliver polarized H^- or D^- within the acceptance of the injector cyclotron JULIC during the 10-20 ms injection period of COSY. The negative ions are produced in a charge exchange type of ion source, which comprises three major components; the atomic-beam part (Bonn), the charge exchange region (Erlangen), and the Cs-beam part (Köln).

The atomic beam part prepares an intense, neutral nuclear polarized hydrogen or deuteron beam. Firstly, gas molecules are dissociated to atoms by means of a RF-discharge. In a second step two 6-pole magnets and RF-transitions prepare the nuclear polarization of the atomic-beam. Subsequently the neutral polarized atomic-beam enters the charge exchange region and meets inside a solenoid magnet a fast neutral Cs-beam. As a consequence, the charges are exchanged according to $\bar{H}^0 + Cs^0 \rightarrow \bar{H}^- + Cs^+$. The negative ions are accelerated at the exit of the solenoid, extracted, and bent by an 90° electrostatic deflector. The ions pass a Wien Filter, an einzellens and an electrostatic quadrupole-doublet in order to match the $\bar{H}^-$ beam to the beam line of the injector cyclotron.

The fast Cs-beam originates from a Cs-oven and is thermally ionized, presently at a potential of about 40 kV. Most of the time the Cs-beam is deflected to the Faraday Cup in the deflection chamber. About 150 ns before the injection period of COSY a Cs-cell is filled with neutral Cs-vapor. The charged Cs-beam, passing through this cell is neutralized and enters subsequently the charge exchange region through an orifice of the 90° deflector.

During the past year the collaboration focused its work to:

- improve the neutral beam intensity
- improve the efficiencies of the RF-transitions and the neutralizer
- improve the transmission of the Cs-beam, in order to reduce sputtering and thermal load in the Cs beam-pipe
- complete the remote-control of the source
- set up a low-energy polarimeter

The H^0 intensity has been measured to be $> 4 \cdot 10^{16}$ atoms/s. From the drop of the H^- beam-intensity, when the RF-transitions are switched on and off, an efficiency of more than 90% (for each of the RF-transitions) could be derived. By adjusting the neutralizer oven-temperature,

the neutralizer efficiency could be improved to exceed 90% also.

After the quadrupole triplet has been set in operation, the Cs-transmission from the Cs-source through the triplet, the neutralizer, the 90° deflector, and the charge exchange solenoid was $> 55\%$. This latter result is unmatched by any other CBS-source, especially in view of the distance of about 2m the Cs-beam has been transported. The collaboration expects to improve even this transmission in the near future. The possibility to transport the Cs-beam with little losses over 2m allows a fast maintenance and is a necessary condition for a high availability of the source.

Up to now an H^- current of about $1.8 \mu A$, with less than 45 nA background could be observed in the beam-line to the injector cyclotron.

A low-energy polarimeter that had been built at the university of Krakau has been assembled and installed in the transfer line between the cyclotron and COSY. The polarimeter has been tested with an H_2^+ -beam successfully.

Presently, the collaboration is waiting to get cyclotron run-time in order to analyze the $\bar{H}^-$ beam with the polarimeter precisely, before the beam is offered to COSY. Only with the help of this polarimeter the Wien-filter can be adjusted properly and a reliable information of the $\bar{H}^-$ beam polarization can be achieved.

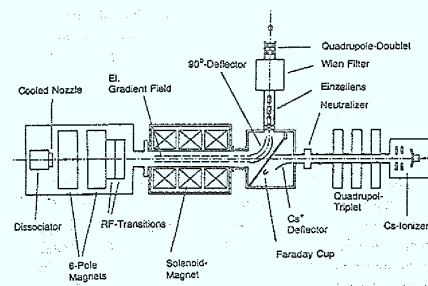


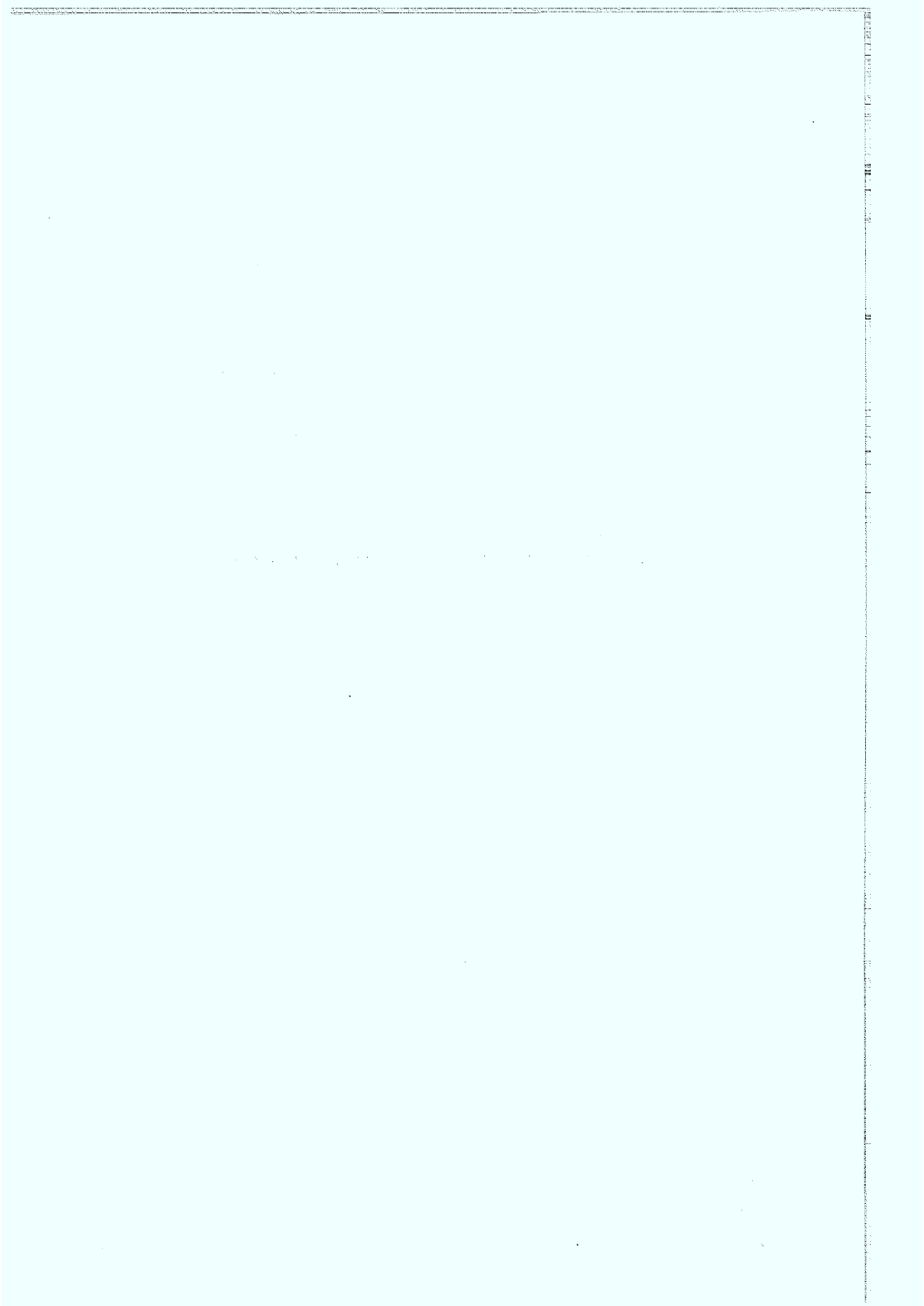
Fig. 1 Schematic of the polarized Source

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*** Institut für Kernphysik, Universität Köln

7. SPECTROMETER BIG KARL



Magnetic Spectrometer BIG KARL

J. Engel, R. Jahn, K. Kruck, H. Machner, R. Maier, P. v. Rossen, P. Schmidt, R. Tölle, V. Vau,
GEM- and MOMO-Collaboration

In several runs performed by the GEM- and the MOMO-Collaboration the spectrometer was used for experiments studying pion production in different reactions. As there had been an extensive remodeling of the ion optical entrance region the first measurements focused besides others also on verifying the success of these modifications. Equally important was the running-in of the specific detector components for the GEM and MOMO experiments and their newly developed data acquisition systems.

Work around the spectrometer aimed at improving the handling and the control of the settings. By measuring in an independent way all currents of ion optically critical components and following proven procedures reproducible adjustments with high fidelity are possible. This system also enables to monitor any critical instability at an early stage.

A new support structure for the position resolving detectors in the focal plane was put into service. Low friction bearings allow for easy movement of the chambers, which is essential when mounting/removing the protective covers for the chambers and the fragile exit window of the spectrometer. Most important, the rails allow to position the chambers back to the original position with high accuracy. An additional benefit is the possibility to use the easy and precise linear movement for calibration purposes of these chambers and checks of the ion optics.

As set-ups in the target area change quite frequently it was important to provide a reliable means for positioning target and detectors systems precisely onto the optical axis of the spectrometer. An essential requirement in this context is that this had to be done with the liquid hydrogen/deuterium target cooled down to ca. 14 K and the target section under vacuum as both parameters influence the mechanical position. A system has been installed in front of the first dipole that consists of a mirror mounted under 45° that can be moved with high precision onto the central axis of the magnetic spectrometer. Using appropriate reference marks it is now possible to align and check components to approx. 0,1 mm under exact measurement condition.

The hodoscope system following the position sensitive detector planes has been significantly enhanced. There are now four independent layers which allow to shape the trigger according to the experimental demands. To raise the resolution of the time of flight measurement of the hodoscope the distance between the two double layers has been enlarged to 2.5 m. This sharpens the condition for selecting specific reaction channels and thus reducing disturbing background. According to the experimental needs the general usage of the four layers is quite diverse. For instance, by using adequate absorbers in front the fourth layer, this layer has been very effective during data acquisition by using its

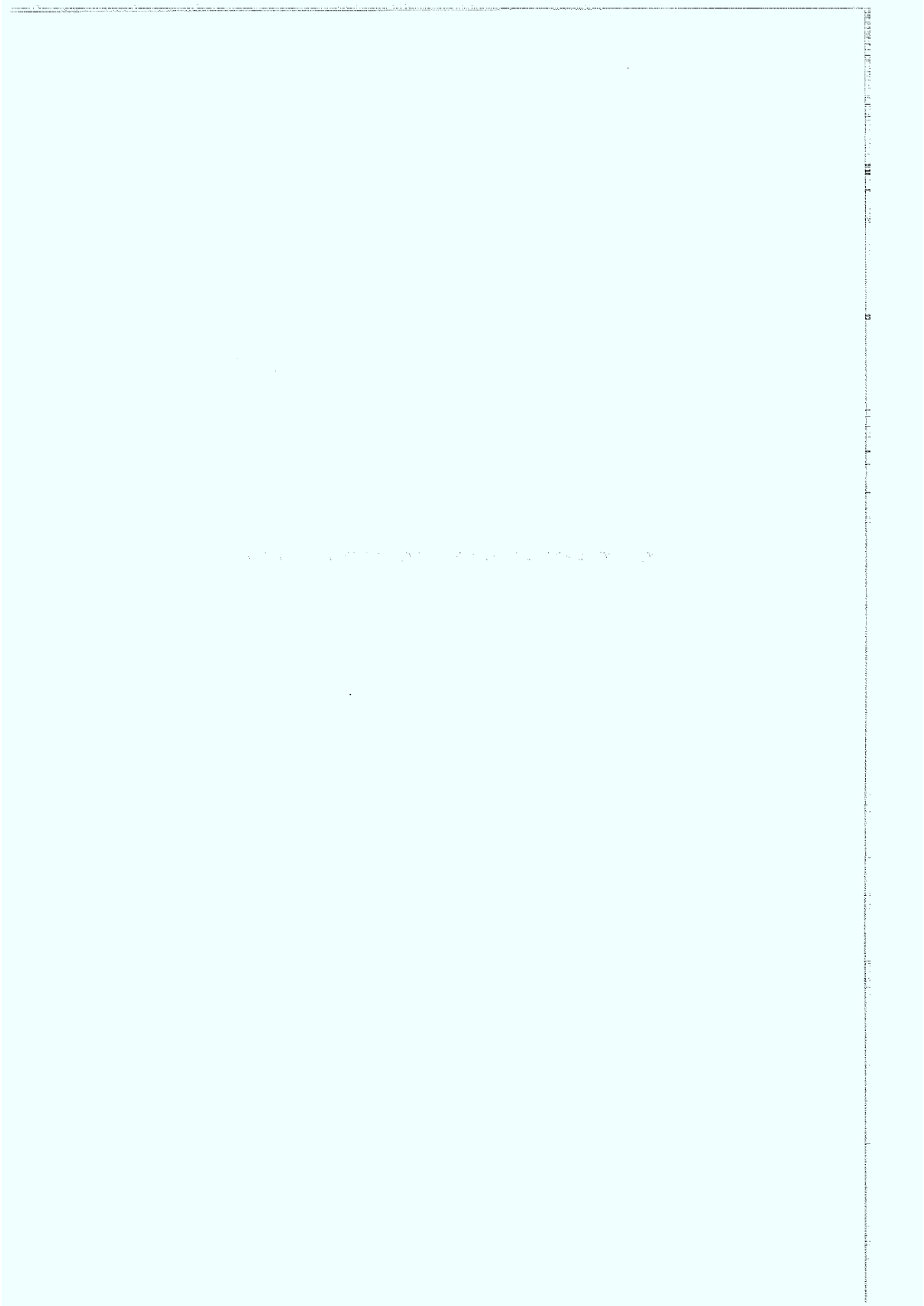
signal in anti-coincidence. The first layer gives the primary signal for time of flight processing while the following layers are helpful for making cuts on the energy loss of the observed ejectiles.

The precision of the data obtainable now at the spectrometer has been greatly enhanced through continuous and substantial improvements of the extracted COSY-beam. Beam diameters of 1.2 mm FWHM at the target have been achieved during data taking runs. The GEM collaboration which for added precision was using a doughnut shaped veto detector with a 3 mm inner hole in front of the target was able to verify this kind of performance throughout the entire run. The goal in those initial runs was to study the reaction $pp \rightarrow d\pi^+$ close to threshold detecting the emitted deuterons in the focal plane of BIG KARL. This reaction was especially suited to check out the ionoptical resolution of the spectrometer. A dedicated detector close to the target registered the pions in coincidence. The paramount questions was whether it would be possible to see the kinematical characteristics of this reaction with sufficient resolution by tracking from the focal plane the particles back to the target point. This important milestone was passed successfully with very satisfying results opening the way for many other experiments that depend on this capability. You find the details concerning this measurement in articles describing the progress of the GEM-experiment in this annual report.

The MOMO scattering chamber at the pivot point of BIG KARL was temporarily removed to make place for the GEM scattering chamber. This measure served the mechanical fine tuning of the components of this chamber and the study of electronic signal processing under the real conditions of electromagnetic interference faced in the forthcoming data taking. Parallel to this the MOMO-chamber underwent its final upgrade. After all vertex modules had been manufactured and tested they were mounted at the rear wall of the MOMO-scattering chamber. This detector system comprises now 672 independent scintillating fibers and is ready to measure coincidently the $\pi^+\pi^-$ and the K^+K^- -mesons. In a run after the successful completion of this final component first data were accumulated in the reaction $pd \rightarrow {}^3\text{He} \pi^+\pi^-$ registering the coplanar pions in coincidence with the ${}^3\text{He}$ in the focal plane. More information on this experiment is laid down in a dedicated contribution found in this report.

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8. RADIATION PROTECTION



As usual the radiation protection duties also in 1994 were very manifold. Some topics shall be treated in the following.

Licenses with reference to the atomic law

At the end of march the license to operation of the low energy measuring area (NEMP) was given. This was the latest license to the operation of the external experimental areas of COSY. Similar to the injunctions of the license to COSY and the other external experimental areas the personnel safety system (PSA) of NEMP had to be tested by the experts of the authorities before the first setting in operation, too. These acceptance tests for NEMP and also for TOF were carried out in 1994. Since during beam operation in the experimental area TOF the gallery and the staircase to the gallery in the experimental area BIG KARL can be entered, the situation cannot be excluded that persons were exposed by increased radiation there. In order to prevent that the PSA was modified provisionally so that no access to BIG KARL is possible during beam operation in TOF. This restriction will be canceled if in the BIG KARL area the measures are finished which prevent an access to the staircase during beam operation in TOF (installation of a door equipped with contacts which are logically tested in the PSA).

Personnel safety system (PSA)

The personnel safety system (PSA) is currently improved and optimized; amongst others the confirmation of the patrolling by means of a magnetic card reader was included in the logic of the PSA:

After the pushing of the quit buttons in the closed areas the authorized persons must type the numbers of the concerned rooms in the keyboard of the PSA and, as the latest measure, confirm the quitting by using a magnetic card related to the person. By that the number of the card and the room, the date and the clock time is stored, till at the next time these data are overwritten.

If during beam times in the COSY-ring this room has to be entered frequently (for instance because of modifications of the setting parameters) much time is consumed, because every time the ring area must be scoured and the quit buttons must be pushed. On our motion the authorities have permitted that the PSA can be modified in such a manner that at the existence of definite criteria the quitting may be resigned: this status can be initiated only by means of a key switch, the key is stored up at the radiation protection officer and delivered up to an authorized person only by signature, the status is finished in case of a duration of more than 15 minutes, at most two persons may enter the COSY-hall. Further modifications refer to the improvements of the information about the status of the PSA and of the individual components as well as to the inducement of measures by means of the keyboard of the Simatic in the operating room of COSY.

Radiation monitoring

The existing radiation monitoring system was expanded by additional neutron monitors required for the operation of TOF and NEMP. The software was modified accordingly. Besides the software for the measurement and indication of the radiation level in the external experimental areas was produced and tested. For these measurements neutron monitors are used. The above named software effects an

illumination of a transparent with the inscription "Increased radiation level", if the dose rate of 7.5 $\mu\text{Sv/h}$ is exceeded. Moreover the software for the so called fast neutron monitors was produced, tested and set in operation. Therewith now the doses of radiation level of very short length but very high can be determined.

Shielding and dose rate measurements

It was agreed with the experts of the authorities that in the same manner as for COSY and BIG KARL also for TOF and NEMP the radiation protection officer of the institute before the official, in the license demanded dose rate measurements (for the confirmation of the effectiveness of the shielding) determines the radiation level at the outer shielding walls during experiment operation. Such dose rate measurements were carried out once for NEMP and twice for TOF, the latter once with the installed neutron monitors during a time interval of 30 hours. The measurements showed that out of TOF the experimental areas for the energies and intensities used during the experimental operation no neutron and gamma dose rates above the background could be ascertained.

The knowledge of the dose rates in the inner part of the COSY-hall during beam operation in the ring, especially at the doors, is very important for the estimation of a possible access of this part of the COSY-hall during beam operation. For that reason these measurements were continued also in 1994. The results show that in case of extreme situations dose rates of above 2 mSv/h can be measured; but there are beam times (for instance the beam operation on the experimental area EDDA with protons of 2.5 GeV) in which very low dose rates (max. 2.9 $\mu\text{Sv/h}$) were measured in the inner hall. Then there are no reasons of radiation protection which forbid a stay in the inner hall.

The gamma dose rate measurements at the beam line of the COSY-ring and at the injection beam line after shutdown were continued. Up to now the gamma dose rates in the ring are less than 10 $\mu\text{Sv/h}$ some days after shutdown. Only at few positions this value is. The maximum value at the injection beam line is 1.6 mSv/h.

Air activity monitoring

The activity emitted together with the spent air from the cyclotron vault (essentially the radio nuclides produced by neutrons in reactions with air components) had to be determined by means of a graphic evaluation. This treatment of the air activity monitoring was modified. Now the monitoring is carried out by means of the automation system Simatic S5. In 1994 the software needed for the monitoring was produced and tested; then the device was set in operation: now, at the begin of a month a list is printed, in which the daily activity release of ^{13}N , ^{15}N and ^{41}Ar of the preceding month is written in. Moreover a continuous surveillance of the regular device function (criterion: the activity release rate is less than a threshold value dependent on the background) and the supervision of the activity permissible for one hour (criterion: the activity release rate exceeds a threshold value) are carried out. If necessary the error messages are documented by means of a printer. A surveillance of the air release rate of the system of the cyclotron building and a continuous test of the regular sucking of the activated air are planned.

IV. Technical Developments

9. DATA ACQUISITION AND PROCESSING

COSY Experiment Data Processing

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In collaboration with institutes inside the KFA, other research centers and external companies the activities due to the enforcement of COSY experiments have been proceeded /1,2,3/.

The main topics in this intention are the development of modular data acquisition systems for COSY experiments, the preparation of the data processing infrastructure and user support in hardware and software, especially implementation support for test systems and for the buildup of data acquisition at the experiments.

Therein, the realization of a data processing structure for COSY experiments comprises the extension of computer networks /4/, the buildup of Server Workstation Systems and the preparation of resources for the analysis of experiment data. In software, support is provided for the operating systems OSF/1, ULTRIX, OS9 and VMS, for various libraries as well as for extensive applications.

The existing local area VAX cluster /5/ is applied to several scientists groups for offline processing, especially for the preparation of COSY experiments.

For the development of COSY experiment data acquisition a UNIX Workstation System has been installed, consisting of 64 bit ALPHA Workstations serie 3000 operating with OSF/1 in conjunction with ULTRIX DEC-Stations serie 5000.

The buildup of Data Acquisition Systems for the experiments GEM, COSY11 and ZDF is in progress. Therefore, the workstation data acquisition software will be provided. In frontend, the installation of hardware will be supported for various Instrumentation Systems and for special modules in CAMAC and FASTBUS including the controller, master, VME CPUs and VICbus interfacing.

For the ZDF experiment, the first extension of the data acquisition was put into operation /6/. This system consists in frontend of one FASTBUS crate, covering the Fastbus Multiblock Instrumentation System, the CHI master and the belonging GSI-Synchronisation modules. Two VME CPUs and one EXABYTE tape, one CPU configured as Subevent Reader and the other one as Eventbuilder are implemented for readout and recording of event data. The DEC-Station 5000/240 with ULTRIX vers. 4.4 has been installed as an experiment workstation. This system will be extended with CAMAC in the near future.

A similar system, equipped with extended online data analysis and visualisation features based on PAW, is in operation at FZ Rossendorf for detector tests.

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COSY Experiment Data Acquisition

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During this reporting period, the activities due to development of the modular data acquisition system for COSY experiments were proceeded in the continual development of the data acquisition software, in the implementation of frontend hardware and the buildup of data acquisition systems at the experiments /1,2,3,6/.

I. Data Acquisition Software

In progress of the establishment of the 64 bit ALPHA AXP workstations at the COSY experiments, the Data Acquisition Workstation Software previously developed on ULTRIX was ported to the OSF/1 operating system. Therefore, in the software mainly the data types have been defined uniformly to support both, the 32 and 64 bit CPU architecture. At now, one common source code of the Data Acquisition Workstation Software exists for the ULTRIX and the OSF/1 operating systems respectively.

The development of the Data Acquisition Software was proceeded. Additional Instrumentation Systems have been implemented and new features in the experiment handling, especially run control functions will be provided. For multicrate configurations, the GSI-Synchronisation System used for the synchronisation of data readout has been implemented in hardware and software. At now the Dedicated Synchronisation Procedure mode (DSP mode) for a single crate or the GSI-Synchronisation mode is supported in any experiment configuration.

The software system has been improved in the OSF/Motif programming. Problems in the dynamical behaviour were eliminated by use of Xevents, especially for interprocess communication.

According to the experience in real experiments, the user interface will be designed more functionally and user friendly.

A. Master Control Process

The user interface has been improved, so that the experiment run can be controlled by one OSF/Motif window. In this window, the major control functions in particular

- start/stop data acquisition,
- open/close tape file,
- enable/disable data analysis,
- enable/disable data recording

are realized by PushButtons. The PushButtons will be updated corresponding to the current experiment state. All other not frequently used experiment preparation and run control functions are activated by PullDown menus. Additionally, the dynamic monitoring of event and data rates are involved in the main window (Fig. 1).

To allow a flexible handling of experiment run, the run control functions can be combined to hierarchical sequences of control functions as follows:

- functions to be effective immediately:
 - write user files at open tape file,
 - open/close tape file at start/stop data acquisition,
 - initialize and start/update and stop scaler, if configured at start/stop data acquisition,
 - update started scalers, if configured at open/close tape file,
 - start/stop data analysis at start/stop data acquisition,
 - set stop data acquisition at end of tape,
 - periodical write file to tape, if tape file is opened,
 - stop data acquisition and close tape file at exit of Master Control,
 - enable or disable of Instrumentation Systems according to data taking,
- functions to be executed at restore time:
 - restore saved setup files of the Synchronisation and of all configured Instrumentation Systems,
 - enable/disable of Instrumentation Systems,
 - monitor run control functions settings,
 - start monitoring scalers, if configured,
 - start selected Instrumentation System Control Processes.

Because of the hierarchical structure of this functions, not concatenated functions can be executed separately. The existing variation in the execution of control functions will be obtained. All selected options can be saved in a file (default RunControl.save) and restored or executed at the

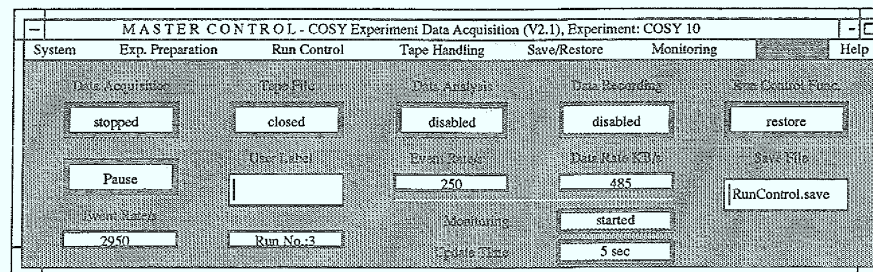


Fig. 1: Main OSF/Motif window of Master Control process

next run. Thereby, a quick experiment restart according to the selected options will be achieved. During the run, each defined sequence can be disabled and enabled again.

To control the state of the frontend hardware, the dynamical monitoring of the setup status of the Synchronisation System and all configured Instrumentation Systems including the status of each configured scaler will be activated at start of Master Control. Additionally, monitoring of statistics will be involved in the Master Control and already existing monitoring functions will be improved.

For the recording of event data on EXABYTE cartridges /4/ multifile volumes are supported.

At the beginning header records and at the end of file trailer records are written to the tape file. Both, the header and trailer records consist of main header record and the hardware configuration file, the module type file and the save files of all configured Instrumentation Systems as well as of the Synchronisation System.

Contrarily, during data taking on request any user file will be written to the tape file. The file records of all to the tape written files are formatted to tape records according to XDR data structure /7/. Therefore, each of such tape records is provided by a header covering the record type and subtype, the Instrumentation System subtype and the run number. The tape record size is limited to 65532 byte. Files exceeding this limit will be stored in concatenated tape records marked in the record subtype as tape records belong to one file.

User tape records are ordered according to the following record types:

- main header record,
- hardware configuration file,
- module type file,
- save file of the synchronisation system,
- save files of configured instrumentation system subtypes (one save file for each subtype),
- user files.

Hence, several events are stored in one tape record, each ending at event boundary. The data of one event are formatted in a header and the collection of subevent data of all fired Instrumentation Systems. In the event header, the event length, the event and trigger numbers as well as the number of following subevents are stored. Accordingly, in the subevent the corresponding Instrumentation System number and the subevent length are marked, followed by subevent data.

The assignments of Instrumentation System numbers to the Instrumentation System Subtypes are part of the experiment database. Of course, each Instrumentation System Subtype has a singular data format, depending on the supported module types and the readout sequence.

B. Instrumentation Systems

In order to achieve a high degree of flexibility in the configuration of Instrumentation Systems according to setup, control and subevent readout, the handling of the frontend procedures have been reorganized. Generally, procedure lists, each consisting of a set of procedures, its version and

the belonging parameters, have been implemented in the workstation software instead of single procedures.

In a common ASCII file (default ISProcList.dat) to each subtype of Instrumentation System a set of procedure lists will be defined corresponding to following functions:

- setup list, consisting of procedures for download of module parameters,
- control list, covering procedures for pedestal or thresholds building by digitizing systems or control operations for auxiliary systems,
- read setup list, with procedures for readout of module setup parameters,
- readout lists, collecting procedures for subevent readout.

In one list, each procedure is dedicated to one module type or a group of different module types. In addition, each procedure is signed by a procedure code covering the function type and a mark whether parameters are required or not.

For each subtype of Instrumentation System one procedure list of the type setup, control or read setup can be defined. Hence, several readout lists can be defined for each subtype of digitizing system for subevent readout. Corresponding to the configured instrumentation systems, the appropriate procedure lists will be loaded in the experiment database from the common Procedure List file at startup time. During the setup, the procedures will be computed according to the procedure code.

For each digitizing system, several readout procedure lists, each covered with a priority and a different trigger pattern, can be defined for subevent readout at experiment run. In one readout list each readout procedure can be enabled or disabled additionally.

The Instrumentation Control Processes will be improved by standardizing many functions, especially the handling of the procedure lists, monitoring and the save and restore operations. Building pedestals and/or thresholds by software trigger using the feature of GSI-Synchronisation module has been realized for Instrumentation Systems supporting modules needed pedestals and/or thresholds setup. Additionally, pedestals and/or thresholds processing can be executed by external trigger alternately.

At this time eight different subtypes of Instrumentation Systems are implemented in hardware and software, carrying by extensive tests, in particular the camac fera system /5/, the fastbus multiblock system, the camac tdc2277 system, the fastbus LeCroy1800 system, the camac scaler system, the camac drams system, the camac tdc2228 system and the general multipurpose system.

B1. Camac Scaler System

The camac scaler system previously designed as an auxiliary system for the LeCroy Mod.4434 has been extended in addition to the digitizing system. The operating mode will be defined during the setup procedure. In both operating modes the counts of all scalers of one Instrumentation System are stored in a variable located in the VED. At update time or scaler overflow the counts in the variable will be updated appropriately.

In digitizing mode, the subevent is built by scaler counts accumulated between two triggers. Depending on the setup of readout procedure lists, the subevent is fed to the event-

path only or will be sent by unsolicited message to the control process for monitoring respectively.

Two modes of scaler monitoring are realized for each configured scaler, the periodical and triggered mode. During periodical monitoring the scaler readout and monitoring update time are defined by the user. Hence, the triggered monitoring is applied to the digitizing mode, the monitoring update rate is synchronous to the trigger rate. This monitoring mode should only be used for sufficient low trigger rates. In auxiliary mode only periodical monitoring is supported. Certainly, in digitizing mode both monitoring modes can be chosen for each scaler alternately. Each scaler will be monitored in a separate window. Of course, all configured scalers in one and different Instrumentation Systems can be monitored simultaneously. In the software for each monitoring mode a dedicated monitoring process is realized.

In addition to the 32-fold LeCroy scaler Mod. 4434 the implementation of the 12-fold scaler Mod. 2551 in this Instrumentation System is still in progress.

B2. Fastbus LeCroy1800 System

In the Fastbus LeCroy1800 System, a selection of the serie 1800 modules produced by the LeCroy Company are supported, in particular the 96 channel TDC1879, the dual range 64 channel TDC1875A, the multihit 96 channel TDC1876, the 64 channel ADC1881 and the calibration and trigger module CAT1810. The implementation of the 96 channel multihit TDC1877 is currently in progress. In one Instrumentation System the mixed mode operation of any or of all supported modules is implemented. Accordingly, one CAT1810 module can be used for several Instrumentation Systems inside one crate.

During this reporting period, the communication to frontend regarding the setup of parameters of all supported module types has been realized according to the new structured server functionality and the OSF/Motif user interface has been improved respectively. Additionally, several monitoring functions have been involved too.

The data readout of different module types, each type marked by its own characteristics, is oriented according the following criteria:

- for each module type a dedicated readout procedure is used,
- to each module type producing event data a variable is assigned. The position of set bits in this variable corresponds to the slot number of module locations of the same type.
- to modules, equipped with Sparse Data Scan, Broadcast Classes are assigned, one readout Broadcast Class for each of this module types and one pattern of Broadcast Classes for each module, to specify the Broadcast Classes to which the module will respond.

Combining the sequence of the readout procedures in the procedure list in order to the conversion time of different module types with the variable covering the location of module of the same type and the Broadcast Class patterns, the readout of different module types will be optimized. Besides, all modules of the same type can be enabled or disabled from data taking by enabling or disabling the appropriate readout procedure and any module of one type

can be disabled from readout by manipulating the Broadcast Class patterns.

B3. Fastbus Multiblock System

In the Fastbus Multiblock System the 32 channel QDC type 10C2 and TDC type 10C6 modules produced by Phillips Scientific Company are supported. Both module types are characterized by a short conversion time (10microsec.) and are equipped with onboard hardware-driven data reduction (thresholds, pedestal subtraction, zero and overflow suppression) as well as fast Multiblock data readout, and thereby predicted for experiments with a high event rate.

Detected hardware problems by the QDCs during the test period, were eliminated in cooperation with Phillips Scientific. Therefore, various test software has been developed, based on Fastbus-NIM routines, on the Fastbus Master under the operating system OS9.

This Instrumentation System has been implemented in hardware and software in the data acquisition system and approved by testing detectors.

In course of this reporting period, the server process running in the Fastbus Master CHI was structured according to the EMS server concept and the GSI-Synchronisation mode has been involved too. The setup and readout procedures have been improved appropriately, especially the pedestal building by software trigger using the GSI-Synchronisation module was realized.

The Fastbus Multiblock Control process has been standardized according to the extended data structures and the handling of OSF/Motif user interface.

B4. Multipurpose System

To perform an unexpected and quick implementation of new, by other Instrumentation Systems not supported hardware in the data acquisition, especially during the experiment time, the Multipurpose Instrumentation System has been developed. Any modules, supported by procedures in the frontend server processes can be collected to one Multipurpose Instrumentation System. Generally, this condition is fulfilled for all CAMAC and FASTBUS modules by the implementation of the standard CAMAC Esone and FASTBUS NIM routines in the server processes at the Readout Controllers (ROCs) (Fig. 2).

In the Multipurpose control process, module test and data acquisition modes are supported.

In data acquisition mode the Instrumentation System can be defined as an auxiliary or digitizing system. In both operating modes setup, control and for digitizing systems the readout procedure lists are represented by simple ASCII files. Each procedure list file consists of one or more procedure lists. Therein, to each procedure list a set of procedures is attached. Procedures are defined by the procedure name, the procedure version and parameters according to the format ProcedureName;vers(p1,p2,...,pn). By omitting the version, the highest version of the procedure is taken. Omitted parameters assume procedures without parameters.

In the file, procedure lists are separated by the character %. Comments if inserted, beginn with an exclamation mark. Comment and blank lines are permitted.

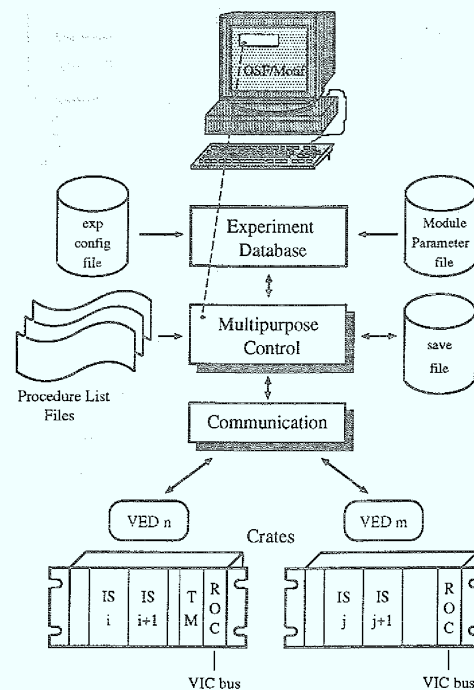


Fig. 2: Structure of the Multipurpose Control process

In digitizing systems, for setup and readout, different procedure list files are required. To perform definite subevent decoding in case of digitizing systems the following restrictions are recommended:

- all in the Instrumentation System defined modules should produce event data with clear defined structure,
- the readout sequence of modules in the readout procedure list file should correspond to the order in Instrumentation Data Set of the hardware configuration file.

The setup handling including the save and restore operations are similar to other subtypes of Instrumentation Systems.

For module test the execution of single procedures as well as the execution of defined procedure list files are implemented.

II. Offline Eventdistributor

The decoding of experiment data requires the knowledge of the hardware configuration, the data formats of all configured Instrumentation Systems and of the readout sequence of modules inside each Instrumentation System. To lighten the experimentalists in data decoding and experiment documentation, for the offline data analysis a software tool, the Offline Eventdistributor has been developed.

The functionality of the Offline Eventdistributor is the tape handling, the decoding of raw event data and the distribution of decoded event data to the dedicated data analysis processes as well as monitoring of various data saved at experiment run (Fig. 3).

Tape handling comprises open and close tape, tape motion functions (skip file or record and rewind tape), and monitoring regarding tape directory and tape status.

For data analysis the experiment database will be configured according to the state at the experiment run time. Therefore, all file header and user records will be formatted to disk files and the required files will be loaded into the experiment database. The decoding of raw event data takes place automatically corresponding to the hardware configuration at experiment run time, as required in the Hardware Oriented Event Format (HOEF) or Detector Oriented Event Format (DOEF), like during online data analysis [1]. To the Eventdistributor connected data analysis processes will be provided by event buffers of requested formats.

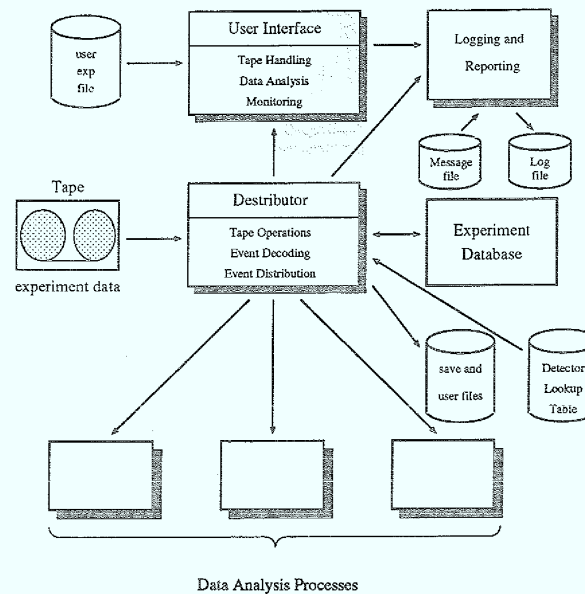


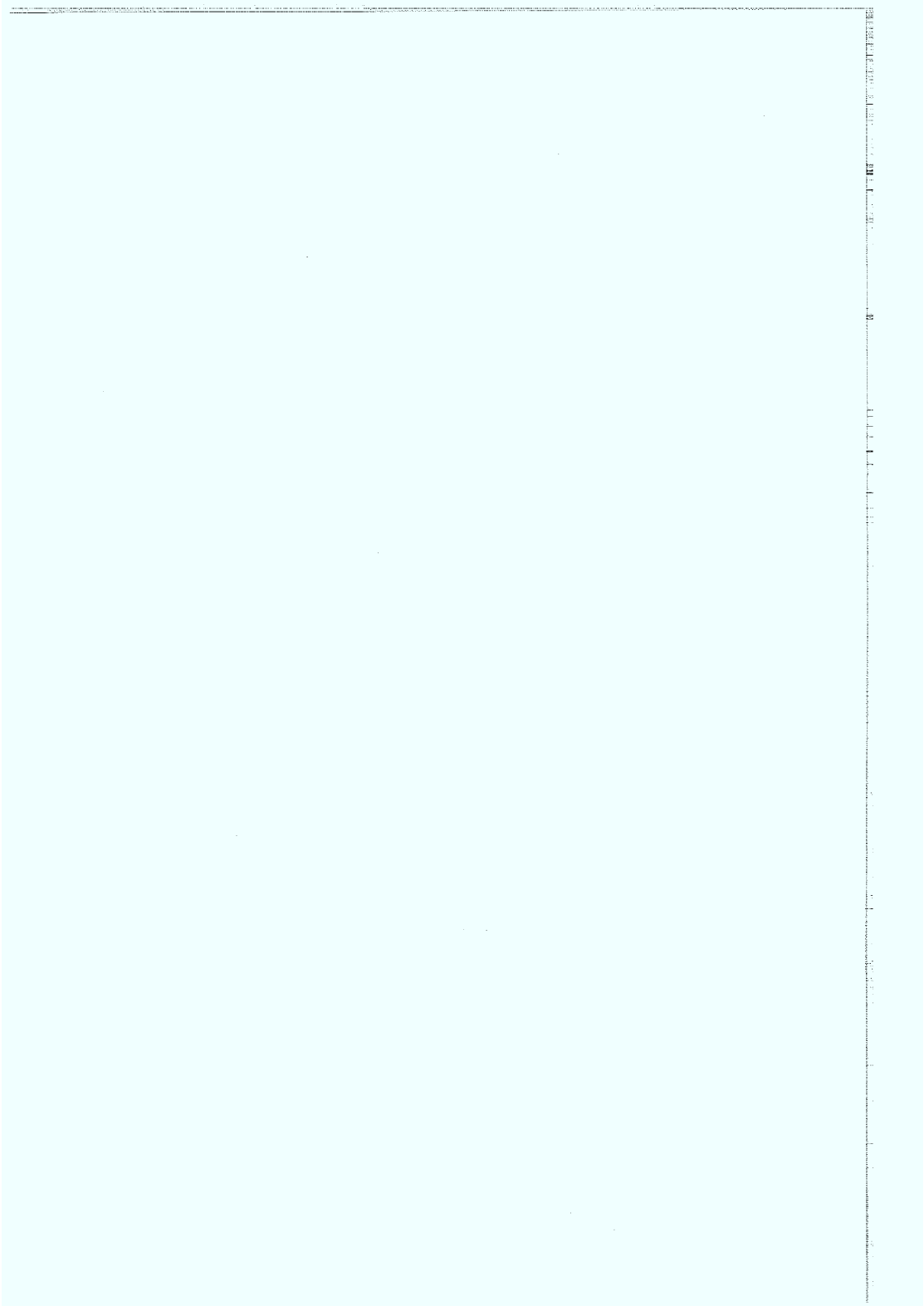
Fig. 3: Structure of the Offline Eventdistributor

In software, the Offline Eventdistributor is structured in the OSF/Motif user interface, the distributor process and the logging and reporting system respectively. At this time, the communication with the data analysis processes is arranged by semaphores; event data buffers are mapped into shared memory sections. For the user written data analysis processes a template is provided.

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10. ELECTRONICS



Electronics

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The activities were concentrated on the diagnostic equipments in the whole COSY-area (source beamline, cyclotron, injection beamline, synchrotron COSY and extraction beamlines) and on the nuclear experiments each with about 50% of our manpower.

Some of the activities are described with more detail in this Annual Report. In the following the most important are summarized:

- The multi-wire proportional chambers (MWPC) in the extraction beamlines have been used extensively. Several improvements concerning lower gas consumption, on-line monitoring, mechanical retouchings, and larger high voltage stability are necessary and were started. Solutions for these problems could be found and will be realized. A calibration measurement with 38 MeV protons from the cyclotron is prepared to extend the usability of the MWPCs also for absolute intensity measurements [1].
- For the BPM-electronics [2] systems at COSY, which process the monitor signals of the 29 beam position monitors and need a Local Oscillator (LO)-signal in the narrowband mode of operation, a PLL-controlled tracking LO-synthesizer [3] was developed and manufactured. The reference input is the COSY-RF. The frequency of the generated output signal can be chosen, by GPIB-interface, equivalent to harmonics ($n=1-10$) of the COSY-RF and hence of the revolution frequency of the particles. By this, the needed LO-frequency is, e.g. in closed orbit measurements, automatically the correct one and an individual frequency setting is no longer necessary. Therefore closed orbit measurements can be performed at any time during the COSY-cycle, especially also in the acceleration ramp. Special care was taken about the spectral purity of the LO-output signal, which could be realized to better than -60 dBc at >8 Hz frequency offset in the 10 kHz - 500 MHz range.
- For test and repair of the analog BPM-modules (BPMcP and BPMcA) [2] a computer-controlled measurement equipment was finally completed. The essential module parameters and properties can be checked and documented. The electronic devices used (function generator, programmable attenuator, oscilloscope, network analyzer) are controlled via GPIB by the VEE-system. The settings of the devices are selected by graphical user windows and the data analysis is performed by own-developed application programs.
- The electronics systems of the Stripline Unit (STR) [4] and the Schottky-Pickups (SchPU), which are used until now for the tune-measurements, were modified and essentially extended for the measurement of the Beam Transfer Function (BTF) [5]. The BTF corresponds, in consideration of the transfer impedances of the STR and the SchPU, to the ratio of the pickup-signal and the excitation amplitude. The extension of the STR-electronics contains: noise source with remote switchable bandwidth of 50 kHz and 1 MHz, mixer-electronics to transform the center frequency of the noise spectrum to the desired excitation range, programmable attenuator to

control the excitation power, fast switch (GaAs) with optically coupled trigger circuit and remote controlled coax-relay to select the excitation signal from the noise source or the tracking generator of a spectrum analyzer. The controls can be performed remotely via an integrated GPIB-interface. For the analysis of the SchPU-signals the BPM-electronics [2] was modified and extended. In the BPMcA-module the demodulator in the narrowband channels was bypassed to make available the 10.7 MHz-intermediate signal and in the BPMcD-module the FIFO-memory size and the sampling rate were enlarged to gain higher resolution and frequency range in the FFT-spectrum. In order to make possible the digitalization of the 10.7 MHz-signal and its conversion into the frequency domain by FFT, the 10.7 MHz-signal (with 50 kHz bandwidth) is transformed down into the 100 kHz-frequency range. The mixer electronics was built up in an additional VXI-module. The Local Oscillator-signal for the mixer circuits is supplied by a commercial synthesizer, its frequency is remotely preset via a GPIB-interface.

The most important activities for the experiments were the following:

- The motor control (angle resolution $3 \cdot 10^{-5}$ in the 8° range) for the low-energy X-ray spectrometer has been expanded to a twin system. Computer control and programming of the single system, realized by use of the program-system LABVIEW, was reorganised for this purpose.
- For the internal Lambda-fission-experiment and additionally for preparing the target positioning of the 0° -facility-experiment a stepper motor control was developed to realize a fast but smooth movement of a thin strip target which is attached mechanically only on one side.
- For the GEM-experiment at BIG KARL various converter modules to adapt levels and impedances have been built as well as a high voltage distribution system for 100 photomultipliers similar to devices which have been built in 1993 for the COSY-11 -experiment.

Furthermore, all groups in IKP were supported in electronics matters by consultation, repair of devices, acquisition and stock-keeping of electronics components.

References

- [1] H. Labus et al., Beam Profile and Intensity Measurements in the COSY Extraction Beamlines, this Annual Report
- [2] J. Biri et al., Beam Position Monitor Electronics at the Cooler Synchrotron COSY Jülich, Real Time '93 Conference, Vancouver, 1993
- [3] J. Bojowald et al., Tracking LO-Synthesizer for the BPM-Electronics at COSY, this Annual Report
- [4] J. Bojowald et al., Longitudinal and Transverse Beam Excitations and Tune Measurements, Annual Report 1993, p. 174
- [5] J. Bojowald et al., Electronics for BTF- and Tune-Measurements at COSY, this Annual Report

Measurements with the Wall-Current Monitor

J. Bojowald, K. Bongardt, J. Dietrich, H. Lawin, G. Lürken, I. Mohos

The broadband wall-current monitor (WCM) [1,2,3] is used to investigate the bunch shape of the COSY beam. Furthermore from the pulse distances and the average current measured by passive integration the number of circulating particles is deduced. The data analysis is performed by the HP-VEE system [4].

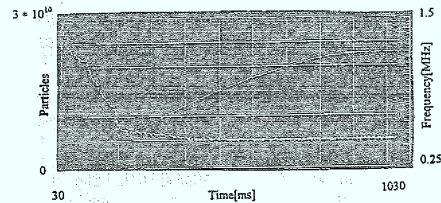


Fig. 1: Number (left scale $5 \cdot 10^9/\text{div}$) and revolution frequency (right scale $0.2 \text{ MHz}/\text{div}$) of the circulating protons dependent on time ($0.1 \text{ s}/\text{div}$) after the start of the COSY cycle

In Fig. 1 the number and the frequency of the circulating particles is shown for a typical process of accumulation and acceleration up to 1030 MeV/c dependent on the time after start of the COSY cycle (COSY 0). The injection occurs 26 ms after COSY 0. The maximum number of particles is observed 34 ms after COSY 0, for this case being $3 \cdot 10^{10}$. Strong particle losses can be seen before the start (103 ms) and the first phase (till about 200 ms) of the acceleration ramp. After that time the proton number stays approximately constant being $5 \cdot 10^9$. In order to clear up the losses during injection and acceleration further investigations and simulation calculations [5] are in progress.

In Fig. 2 the debunching effect after switching off the COSY-RF voltage (1307 ms after COSY 0) is demonstrated. For this case overlapping of the individual bunches starts about 3 ms after switching off the RF-amplitude. At latest after further 20 ms coasting beam is existing in COSY.

References

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- [2] H. Lawin, J. Bojowald, K. Bongardt, J. Dietrich, U. Hacker, G. Lürken, R. Maier, I. Mohos, K. Sobotta, IKP Annual Report 1992, p. 195
- [3] R. Maier et al., Proc. of EPAC 90, Nice, 1990
- [4] E. Brökel, J. Dietrich, H. Labus, G. Lürken, I. Mohos, IKP Annual Report 1993, p. 197
- [5] H. Stockhorst, private communication

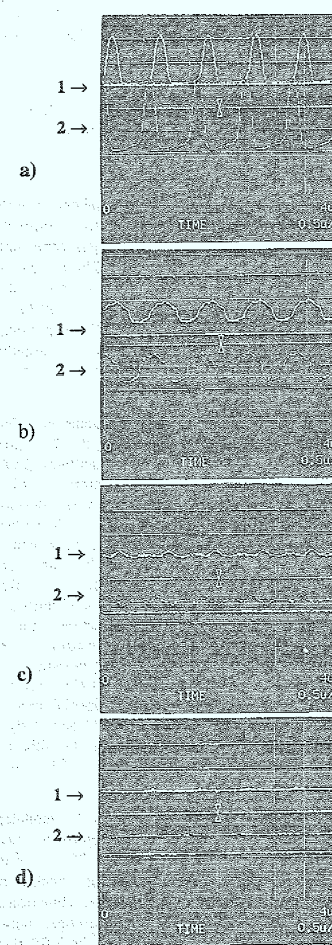


Fig. 2: Bunch shapes measured with the standard BPM electronics (7 MHz, curve 1) and the broadband amplifier without baseline restorer (70 MHz, curve 2) showing the debunching of the beam 1.30 s (a), 1.31 s (b), 1.32 s (c), 1.33 s (d) after the start of the COSY cycle (COSY 0). The accelerating voltage was switched off 1.307 s after COSY 0.

To measure the beam position, 29 monitors are installed along the COSY accelerator ring. The preamplified signals of the monitors are processed by the analog unit of the BPM electronics. The total amount of the electronics (two sets for the X and Y planes at each monitor and reserve units) is 80 sets. The BPM electronics consists of three main parts: analog preamplifier (BPMeP), analog processing modul (BPMeA) and digital control and data acquisition modul (BPMeD).

The testing of the analog BPM electronics (in Fig. 1 the block diagram is shown) needs a lot of complex high frequency measurements in the frequency range DC-200MHz and the amplitude range 5uV-1V. A test in all modes needs about 1700 single measurements. The processing of the test data means simple computations in large quantities. The long time for manual measurements and the higher probability of human failures are not acceptable for serial tests. For this reason an automated test station was developed and realized.

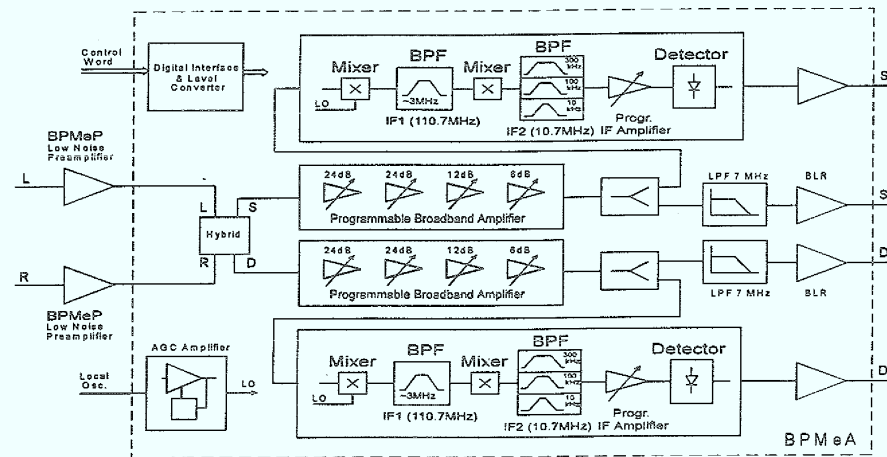


Fig. 1: Block diagram of the analog BPM electronics (BPMeP and BPMeA)

For the complete checking of the analog specifications signal generators, attenuators and oscilloscope are necessary with accuracies much better than the specification of the BPM electronics and with possibilities for interfacing to any bus system. The signal generators and measuring instruments, the VXI crate for the unit under test and the computer (IBM/AT) are placed in a 19" rack. All equipments are interfaced to the PC via standard instrumentation bus (GPIB). The VEE testprogram [1] running under the operating system OS/2 controls the measurements and processes the test data. For each functional part in the BPMeA stand alone program packages are developed, they can be started independently or also combined together according to the main menu.

Test segments are available for measuring the

- equivalent input noise of the preamplifiers
- gain accuracy of the preamplifiers
- accuracy of the sum and difference signals

- gain accuracy of the programmable amplifier
- broadband amplifier's frequency response
- frequency response of the low-pass filters
- frequency response of the IF filters
- linearity and offset of the level detectors.

The test results with tolerance limits appear on the display also graphically. The measured data, the unit identifier and datum are stored in file.

Ref. 2, describes the BPMeA's operation modes and specification. Fig. 2 shows the block diagram of the automatic test station.

References:

- [1] VEE = Visual Engineering Environment (primarily HP)
- [2] J. Biri et al., Beam Position Monitor Electronics at the Cooler Synchrotron COSY Jülich, Real Time '93 Conference, Vancouver, 1993.

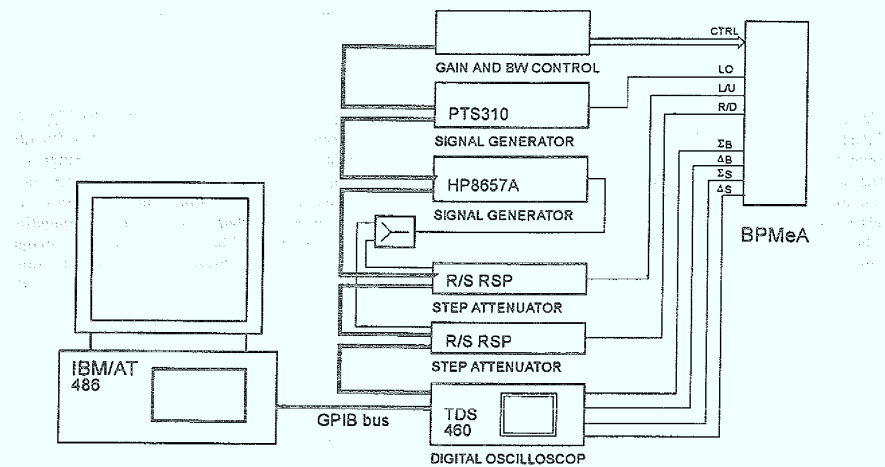


Fig. 2 Block diagram of the automatic test station

For the measurement of the beam positions, 29 beam position monitors (BPMs) [1] are installed in COSY. By the accompanying BPM-electronics [1] the monitor signals can be processed (i) in the broadband mode (time domain) with 7 MHz bandwidth and (ii) in the narrowband mode (frequency domain) with presettable band pass filter widths of 10, 100 and 300 kHz in the 0.2 - 60 MHz frequency range. The broadband mode reproduces the bunch shape and is used for the analysis of the beam turn by turn; but the resolution is limited due to noise, especially at small deviations of the beam from the centre position (small Δ -signals) typical for the aspired ideal beam orbit. The narrowband mode results in much better resolution due to the higher signal to noise ratio, but is applicable only with stored beam under „stable“ conditions due to averaging over several turns caused by the filter response times. This is fortunately the most often case of beam position measurement, e.g. in routine determinations of the „closed orbit“, the instantaneous equilibrium orbit of the beam in COSY.

During narrowband measurements the BPM-electronics need for their mixer circuits a Local Oscillator (LO) - signal, which makes possible by correct frequency setting the selection of the frequency content in the input signal to be measured. Until now the LO-signal was generated by a commercial synthesizer, the frequency of which was set by the COSY-control system via the GPIB-interface. Now a tracking LO-synthesizer [3] was developed and constructed, which allows to preset the LO-frequency equivalent to harmonics of the COSY-RF and hence of the revolution frequency of the particles in COSY. By this the needed LO-frequency is, e.g. during closed orbit measurements, automatically the correct one and an individual frequency setting is no longer necessary. The tracking speed is >10 MHz/s and exceeds the maximum frequency slope in course of the acceleration ramp, so narrowband closed orbit measurements can be performed at any time in the COSY cycle. Much care was taken about the spectral purity of the synthesizer output signal, which could be realized to better than -60 dBc (at >8 Hz frequency offset in the 10 kHz - 500 MHz range) and renders possible to suppress the disturbing influence of neighbour peaks in the frequency spectrum and thus to enable a measurement accuracy according to the good signal to noise ratio in narrowband measurements.

Fig. 1 shows the block diagram of the tracking LO-synthesizer, the most important part of which is an extended and especially optimized phase locked loop (PLL). This includes besides the voltage controlled oscillator (VCO) and the phase detector, additionally, a mixer and stable crystal oscillator (XTCO) to shift the LO-signal frequency by the amount of the intermediate frequency (110.7 MHz) of the mixer circuits in the BPM-electronics. Furthermore the loop contains a presettable frequency divider by which the LO-frequency can be

chosen equivalent to harmonics ($k = 1 - 10$) of the COSY-RF frequency. The LO-frequency therefore is defined by

$$f_{LO} = f_{ZF} + k \cdot f_{HF} = 110.7 \text{ MHz} + k \cdot f_{HF}$$

with f_{ZF} the intermediate frequency of the BPM-electronics, f_{HF} the frequency of the acceleration cavity which corresponds to the particle revolution frequency and k the harmonic number.

In order to gain the needed loop settling speed a loop filter must be adjusted according to the harmonic number. The setting of harmonic number and loop filter is carried out by an internal microcontroller which is connected with the COSY-control system via the GPIB-interface. The harmonic number can also be set by encode switches at the frontpanel. At time of turning on the supply voltage the control is set to „local“ and after arrival of the first GPIB-command switched to „remote“; the status is indicated by a LED at the frontpanel.

The „PLL locked“-status is also indicated by LED. With COSY-RF signal present at the input the loop is closed automatically. In the case of missing COSY-RF signal the internal microcontroller periodically scans the total LO-frequency range; the LED „PLL locked“ stays non-illuminated until the COSY-RF signal is present again.

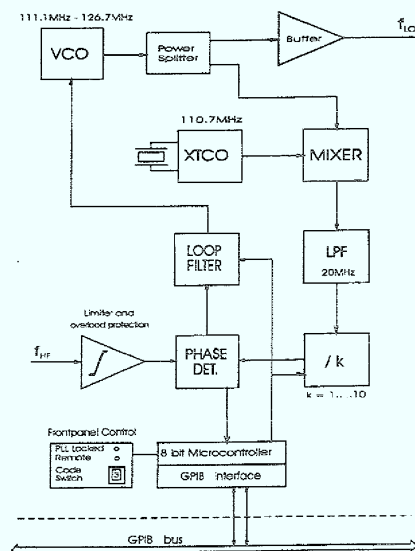


Fig. 1: Block diagram of the tracking LO-synthesizer

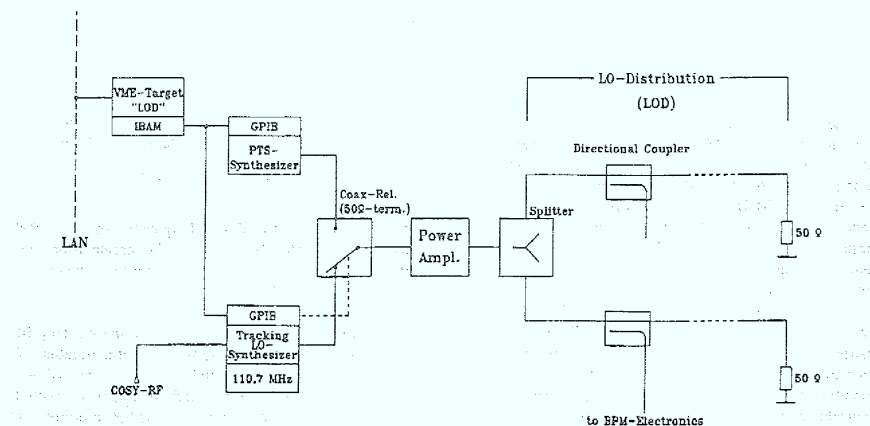


Fig. 2: Schematic overview of the LO-signal system for the BPM-electronics at COSY

Fig. 2 gives an overview of the LO-signal selection and distribution (LOD) system for the BPM-electronics at COSY. The LO-signal can be selected from the tracking LO-synthesizer or a precision synthesizer (from company PTS) via a coaxial relay controlled by the microcontroller and GPIB-interface. The tracking LO-synthesizer is used during usual accelerator operation in measurements to monitor the COSY beam, e.g. in closed orbit measurements. The revolution frequency of the particles or harmonics of it are analysed in these cases and the tracking LO-synthesizer always yields the correct LO-frequency at any time during the COSY cycle. The PTS-synthesizer is used in measurements with fixed LO-frequencies different from harmonics of the revolution frequency, e.g. for monitoring the cyclotron beam (27 MHz-structure) injected into COSY or for the measurement of machine parameters (e.g. tune), which

are determined by the transverse oscillation of the beam particles (betatron oscillation) and are deduced from the frequencies of the sidebands in the frequency spectrum of the transverse monitor signal.

References:

- [1] R. Maier et al., Non-Beam Disturbing Diagnostics at COSY-Julich, EPAC'90, Nice 1990
- [2] J. Biri et al., Beam Position Monitor Electronics at the Cooler Synchrotron COSY Julich, Real-Time'93 Conf., Vancouver 1993
- [3] I. Mohos, private communication; G. Schnitzler, Diploma Thesis, 1994

Fig. 1 shows the blockdiagram of the total BTF-measurement system. The STR and the SchPUs are installed in COSY about 65 m apart from each other. Directly at the STR are mounted also the 50 Ω -terminators (100 W) for the striplines and the 50 Ω -attenuators (100 W; att. 10 dB) for coupling out the exciting signal as reference for the BTF-measurements. The subdevices to supply the STR, i.e. electronics for signal generation according to the selected excitation mode, power amplification and computer control, are installed in the inner part of the COSY-hall near the STR, but outside the COSY-tunnel. The „Schottky-electronics“-systems for processing the pickup signals and the reference signal are installed inside the COSY-tunnel directly at the place of the SchPUs. Only the computer-controlled synthesizer for the Local Oscillator (LO)-signal of the mixer circuits in the BPMcA-modules [3], which is used to select the investigated frequency range in the input signal, is situated

In the tune-measurements performed up to now the tune was determined only from the estimated center frequencies of the side bands in the frequency spectrum of the pickup signals. The precision of the resulting tune values is not sufficient enough (about 1%) due to the inaccurately defined center of the broad and partly interrupted peaks. The side band frequencies and consequently the tune values can be deduced more precisely in BTF-measurements from the phase response, because a phase jump occurs in the side bands due to the resonance

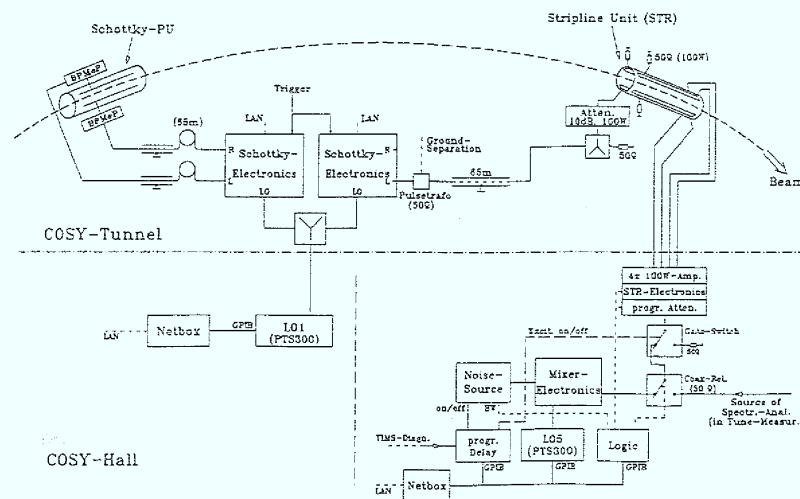


Fig. 1: Blockdiagram of the BTF-measurement system

process. Therefore as result of the extended measuring equipment a more precise tune-measurement gains at the same time. But the BTF-measurements moreover allow further important measurements, e.g. the determination of the particle momentum distribution in the beam and of the accelerator impedance, and they yield significant indications about possible instabilities and the magnitude of the necessary feedback signal, if instabilities have to be overcome.

The extension of the STR-electronics contains the following developments: A noise source (white noise) was realized to be used as source of the beam exciting signal with remote selectable bandwidths of 50 kHz and 1 MHz, the center frequency of which can be transformed by a mixer circuitry and a remote controlled synthesizer to the center frequency of the desired excitation range within the 0.4 - 60 MHz frequency range; the diagram of the circuitry is shown in Fig. 2(a). A fast trigger electronics with programmable presetting of time delay and excitation time was built up, which facilitates by GaAs-switch to turn on the beam excitation for short but defined time at any moment during the COSY-cycle. Thereby several excitation periods can be released in turn with presettable time distance within the same COSY-cycle (for average purposes). For adjustment of the excitation power a programmable RF-attenuator was inserted. And a program-controlled coaxial relay was incorporated to select as source of the exciting signal the noise source or the sweep oscillator of a spectrum analyzer, respectively. For the connection to the computer system a GPIB-interface was integrated. This also controls now the signal settings

for the striplines according to the selected excitation mode, as described in Ref. [1].

The Schottky electronics (see Fig. 1) is a modified and extended BPM-electronics [3]. The modification concerns in the BPMcA-module the bypassing of the demodulator in the narrowband channels, so that the 10.7 MHz-intermediate frequency signal is available now, and in the BPMcD-module the enlargement of the FIFO-memory size (from 4K to 64K) and the adaptation of the sampling rate (200 kHz instead of 1 MHz with 100 kHz-bandpass filter in the BPMcA-module), in order to expand the signal recording time and thereby to gain higher resolution (now 3 Hz) in the FFT-spectrum. The extension of the Schottky electronics consists of a mixer circuitry, which down-converts the 10.7 MHz-signal with 50 kHz-bandwidth to the 100 kHz-frequency range, so that it can be digitized in the BPMcD-module or alternatively also in a FFT-analyzer and than transformed by FFT to the frequency domain. The mixer electronics was built up in an additional VXI-module; the functional diagram is shown in Fig. 2(b).

References:

- [1] J. Bojowald et al., Longitudinal and Transverse Beam Excitation and Tune Measurement, Annual Report 1993, p. 174
- [2] J. Bojowald et. al., Schottky-Pickups, Annual Report 1991, p. 238
- [3] J. Biri et al., Beam Position Monitor Electronics at the Cooler Synchrotron COSY Jülich, Real Time'93 Conf. Vancouver, 1993

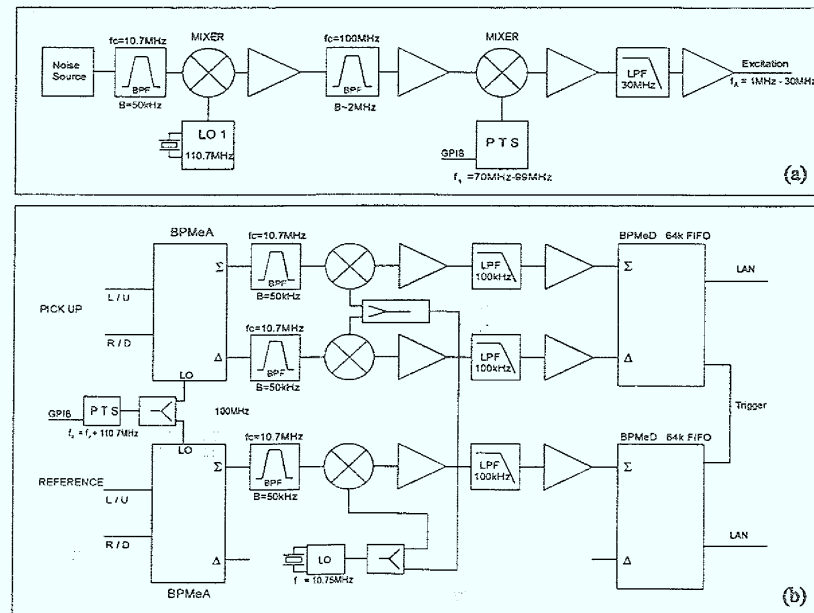
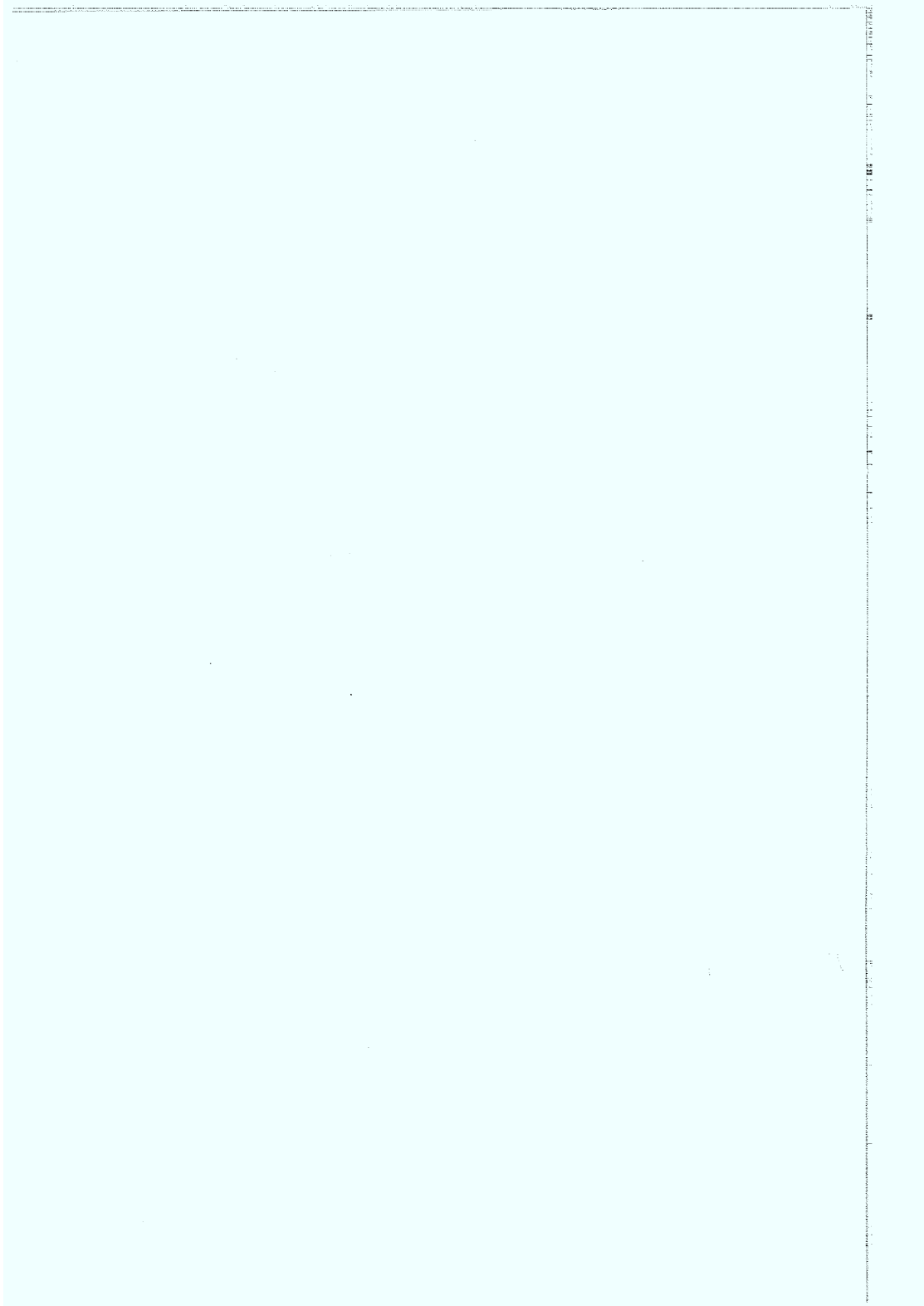


Fig. 2: Functional diagram of (a) the noise source used as signal source for beam excitation and of (b) the mixer electronics in the extended Schottky electronics for down-conversion of the 10.7 MHz-signal to the 100 kHz-frequency range.

11. DETECTORS



G. Fiori, A. Hamacher, S. Igel, H. Machner, H. Metz, D. Protić

As in the previous year the activities of the group were mainly devoted to the construction of the detector systems called Germanium Wall for COSY-experiments [1], [2]. All components needed for the first version are completed, tested and are ready for final mounting.

The cryostat provided with feedthroughs made from flexible cooper leads on Kapton foils was mounted in its position in the experimental area together with a LH_2 -target. After cooling down to 77 K a pressure of 7×10^{-7} mbar could be achieved with the open valve to BIG KARL. Closing the valve improves the vacuum inside the cryostat by an order of magnitude. A separate vacuum system for fixing a LN_2 -dewar was constructed. Now the detectors can be cooled using LN_2 -dewar or liquid nitrogen flowing through the pipes inside the cryostat.

A Quirl-detector with 200 spiral elements was manufactured and tested in a modified test-cryostat. Using special connectors it is quite easy to put the detector from one cryostat into the other. This also applies to the E-detector (Figure 1) of the Germanium Wall. Having a volume of

135 cm^3 it is, to our knowledge, the largest planar germanium detector ever built. The detector was manufactured from a piece of high purity germanium with a diameter of 102 mm and a thickness of 17.8 mm. A central conical hole was bored after preparing the n-contact, about 1 mm thick, by Li-diffusion at 350°C . 32 radial segments were defined by plasma-etching the grooves through the boron implanted p^+ -contact. Two additional Quirl-detectors were manufactured and will be placed in the mountings for GEM-experiments to be used as spare detectors. Several germanium wafers are currently being processed for further Quirl-detectors.

All electronic components for the read-out of approximately 500 detector elements are completed and tested. Various modifications were performed to improve the properties of the components. The sensitivity of the pre-amplifiers and amplifiers was adapted corresponding to expected detector signals. Electronic parts of the sum amplifier were soldered on printed boards. Every 16 amplifiers and discriminators were mounted together with two sum amplifiers inside a NIM-unit and then tested.

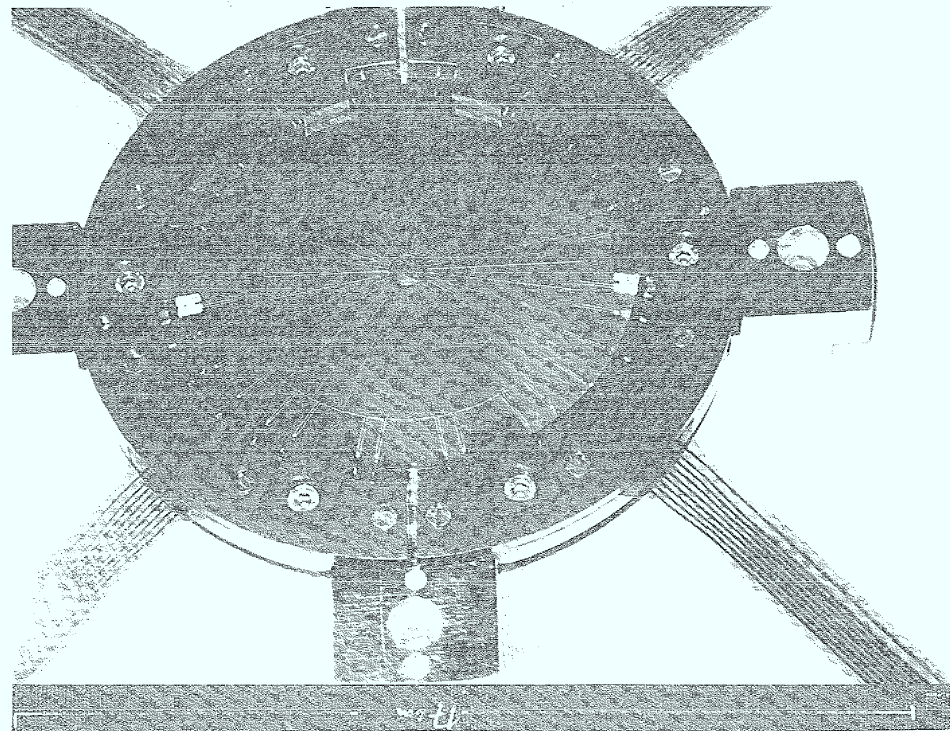


Fig. 1: E-Detector for the Germanium Wall

The development of thin silicon annular ΔE -detectors having a position sensitive structure defined by spirals on both sides has been started. These "Quirl"-detectors will be used for TOF-experiments. A series of 500 μm thick silicon wafers with a central hole was prepared for etching before ion-implantation. Several diodes with implanted layers on both sides, manufactured at INP-Cracow, were tested after evaporation of Al-contacts. This development will be continued in collaboration with INP-Cracow and ISI.

Extensive tests were performed with the prototype of a germanium strip detector [3], together with the collaboration of E.S.R.F., Grenoble. Three elements of $0.2 \times 200 \text{ mm}^2$ separated by 35 μm wide grooves and surrounded by a guard-ring were connected to the first stages of preamplifiers mounted inside the cryostat. The rest of the preamplifiers, made by SMD-technology and developed by KFKI-Budapest, was placed outside the cryostat. As expected, the energy resolution was determined mainly by the noise of the input FET and amounted to $0.85 \pm 0.05 \text{ keV}$ [FWHM] for a 60 keV γ -line. After the tests performed by the detector group in Grenoble one can state that:

- a) neither the reverse current of the detector nor the resistance between the strips limit the energy resolution of the prototype,
- b) there is only a small influence of the strip structure on the charge collection and therefore on the shape of the peaks and
- c) cooled input FET's, selected to match strip capacitance, should be used for further improvement of the energy resolution.

The next steps in the collaboration with KFKI and E.S.R.F. will be devoted to the development of the strip detectors with a resolution of better than 0.5 keV [FWHM]. This can be achieved using special miniature preamplifiers with cooled input FET's mounted as close as possible to the strips. Furthermore, a germanium strip detector with a special test-structure should be used to study the influence of geometrical parameters on the detector properties, for example width and depth of the groove.

Three germanium detectors are manufactured for the zero degree facility of CELSIUS in Uppsala, a collaboration with the University of Stockholm:

1. ΔE -detector with both contacts made by ion implantation having an area of 40 cm^2 and a thickness of 1.1 mm,
2. E-detector with a Li-diffused rear contact and boron implanted entrance contact of 50 cm^2 and 15 mm thickness and
3. ΔE -detector, similar to the first one, but position sensitive in xy-direction with 66 strips on one side and 18 strips on the opposite side. Each of them will be connected to an amplifier chain outside the CELSIUS-ring. The first two detectors are already tested in Jülich and Stockholm. Soon the position-sensitive detector will be mounted and the strips bonded to the finished printed leads on Kapton-foils. After mounting it in a test-cryostat the response of each strip will be examined with its own amplifier chain. All detectors and mounts are so constructed that they can sustain a temperature of up to 470 K inside the CELSIUS-ring.

In collaboration with ISI a series of P-I-N and staircase-diodes made from AlGaAs/GaAs was tested with α - and x-radiation. Regarding the staircase-diodes special attention was paid to the internal multiplication and its influence on the energy resolution. For the 13.96 keV x-ray the best resolution of 0.9 keV [FWHM] could be achieved at room temperature with an internal amplification factor of 4.2.

Commercial silicon and germanium detectors were tested or repaired in urgent cases. Several groups at KFA were supported with some special technical procedures (plasma etching and bonding). A group from the University of Bonn was supported during the development and fabrication of special targets by making use of facilities for photolithography and ion-etching.

References:

- [1] Annual Report IKP 1993, Jül-2879, p. 242
- [2] Annual Report IKP 1993, Jül-2879, p. 7
- [3] Annual Report IKP 1993, Jül-2879, p. 241

Targets

J. Pfeiffer, D. Protić

The most important activity was the preparation and mounting of carbon foils to be used as targets for the COSY injection system. A technician from the COSY crew has been introduced to basic techniques of carbon target preparation. Very large targets of 60 mm in

diameter and a thickness of $20 \mu\text{g}/\text{cm}^2$ were manufactured for the EMA of the injection beam line. Various targets were prepared for IKP-groups and outward laboratories. Various materials (scintillators, foils) were coated by evaporation of different metals.

The production mechanisms of the mesons π^0 , η , and η' in p+p- and in p+nucleus reactions can most effectively be studied through the direct detection of the 2 photons, into which these neutral pseudoscalar mesons may decay. A large acceptance, granulated photon detector, which would be very helpful for this type of research [1] at COSY, cannot be realized immediately. Nevertheless, we are building a test detector module, consisting of 25 pure CsI scintillator elements, which can be used to determine the energy of electromagnetic showers, produced by η decay photons, which hit the central element of the 5*5 scintillator array. The elements are 30 cm long and have the shape of truncated pyramids, with the dimensions ≈ 3.3 cm² on the front side, and ≈ 6.6 cm² on the rear side. Photons originating from η mesons, which were produced close to the NN threshold and which decay symmetrically to the beam axis, have energies of ≈ 400 MeV in the laboratory system. Simulation calculations show, that photons of 400 MeV, which enter the center of the array, produce electromagnetic showers, which deposit 96 % (average value) within the module. The width (fwhm) of the distribution is expected to be 3% [2].

On the basis of this value, we aim to limit the systematic error of the shower total energy determination, which is caused by imperfect energy calibration, to $\pm 1\%$. The energy region, for which the response function (charge at photomultiplier anode / energy deposit in scintillator) must be determined, and the desirable precision of the energy calibration, follow from the distribution of the energy deposits per single element. Photons of ≈ 400 MeV cause energy deposits in one crystal from 1 MeV to 360 MeV. Highest precision is needed for the large energy deposits, which occur close to the center of the shower. In this region we may have 360 MeV in one crystal, but also amounts of ≈ 100 MeV distributed over 4 crystals, if the center of the shower is located close to the edge of a crystal. We conclude, that for η detection we should have a calibration precision of 1% in the region from 400 to 100 MeV and tolerate -as a first rough estimate- an absolute error of ≈ 1 MeV at lower energies. In η' experiments, the high energy precision region has to be extended to ≈ 1 GeV.

For the energy region up to 1 GeV a pure CsI photon spectrometer is best read out with photomultiplier tubes (= PMT). It is well known, that PMTs are non linear for many reasons, in particular in the region of large pulse heights. How can we obtain the shape of the response function and an absolute calibration for energy measurements up to 1 GeV?

Energy deposits in one scintillator element, which are precisely known, can easily be produced at energies in the 1 MeV region using standard γ ray sources. If we want to derive absolute calibration of 1% at high energies from low-energy reference points, we must determine the detector response function for the complete energy range with the same precision. After one has succeeded to

identify π^0 -, η - or η' mesons in an experiment, one will certainly use the well known masses of these mesons to determine corrections for the original calibration parameters.

Well defined energy deposits in the 100 MeV region could be produced with protons, of course. However, due to the high ionization density, the scintillation light output is much less than it is in ionization processes, caused by the electrons of an electromagnetic shower or by electrons produced in low-energy Compton processes. Thus calibration with protons or heavier ions must be excluded.

As indicated in ref. [3] we have examined a light pulser system, which can simulate the scintillation light of pure CsI. The pulse shape is adapted to the ≈ 8 ns time constant of the dominant fast component of CsI light and the intensity can be varied with a precision of $< 1\%$ in a range corresponding to energies ≈ 1 MeV to ≈ 1 GeV deposited in a pure CsI reference scintillator. The light pulser system consists of a fast pulser (risetime=2.5 ns, falltime=15ns, amplitude=10 V), a visible light laser diode ($\lambda = 635$ nm), a lens which produces a parallel beam, 2 dichroic linearly polarizing filters, which can be rotated, and a lens, which expands the beam to cover the entrance window of the PMT [4]. The light intensity behind the polarizers is given as $I = I_0(\cos^2\Theta + R)$. The angle between the principal axes is denoted Θ . In the crossed position we obtained $I = R = 1 \cdot 10^{-3}$. If scattered light was sufficiently suppressed, the precision of the adjusted intensity depends only on the stability of the laser diode and on the error of the angular setting. We use a stepper motor driven rotator with $\Delta\Theta = 0.01^\circ$, which makes the relative intensity error $\Delta I/I < 1\%$ over the full dynamic range of 1000. This result was verified with the help of a PIN-photodiode, which is linear to a very high degree [5].

The 25 element pure CsI-detector module will be read out by the Hamamatsu R4551 PMTs, which have a $\Phi=60$ mm photo cathode with UV transparent entrance window. Fig. 1 shows the dependence of the peak anode current from the pulser light flash intensity. This intensity is expressed through the equivalent energy deposit of γ radiation in a

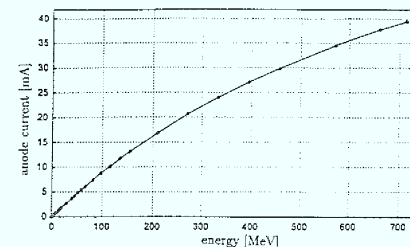


Fig. 1: Peak current at the anode of a Hamamatsu PMT vs. radiation caused scintillation flashes in pure CsI, measured with light pulser system

pure CsI-reference crystal. The voltage at the PMT was -1000 V, linearly distributed between cathode and anode. The pulse shape was such, that we obtain for the ratio (total charge) / (peak current) = 27 ns. The measured response function corresponds perfectly with older results obtained with another calibration method, which is relatively complicated and time consuming [5]. For the operation conditions given above, the non linear response of the PMT starts at ≈ 10 MeV. At 700 MeV the output signal amounts to only 57 % of that of a linear system. Thus careful determination of the response function for the complete dynamical range is unavoidable.

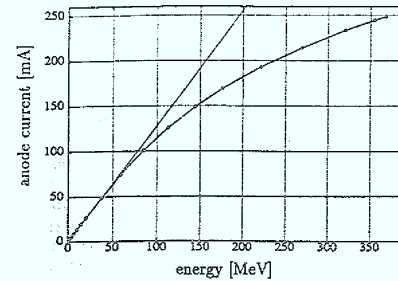


Fig. 2: Anode current at Philips XP2020Q for U=1500 V

We have also started to measure the response function of the Philips XP2020Q PMT, which is widely used to read out fast scintillators. The standard voltage divider base S563 was used. Fig. 2 shows strong non-linearity of the peak anode current for light intensities corresponding to >70 MeV. The anode-cathode voltage was U=1500 V. This effect may be due to space charge in the tube. At U=1000 V (see fig. 3) the charges are lower by a factor ≈ 60 . The response becomes linear up to ≈ 300 MeV, in spite of the smaller electric fields in the tube.

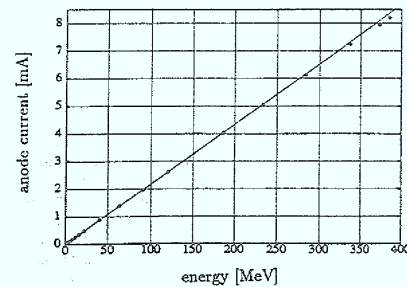


Fig. 3: Same as fig. 2, yet U=1000 V

We conclude that the described pulser system is well suited for precise determination of the current response of PMTs to light flashes. The application of the method is in particular interesting for scintillator systems, which consist of a large number of elements, as the variation of the light intensity over the complete dynamic range can be done remotely through computer control.

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Comparison of Measured and Simulated Features of the EUROBALL CLUSTER Detector

D. Kutchin, R.M. Lieder, W. Gast

The prototype of the EUROBALL CLUSTER detector has been developed in a collaboration between the KFA Jülich and the Cologne University. The Cluster detector consists of seven large, encapsulated Ge detectors in a compact geometry which are surrounded by a BGO escape-suppression shield consisting of 18 BGO scintillation detectors. Details about the CLUSTER detector design and the encapsulation are given in [1]. At the initial stage of its design, Monte-Carlo simulation calculations were carried out to optimise the geometry of the CLUSTER and to determine its detection features, namely, the full-peak efficiency and the peak-to-total (P/T) ratio [2]. All detectors and the crystal are now available and the first tests of the complete system have been carried out. The efficiency, add-back probability and P/T ratio of the CLUSTER detector resulting from both the Monte Carlo simulation calculations and the measurements using the prototype are compared in Table 1 [3].

Table 1: Measured and simulated features of the EUROBALL CLUSTER detector for $E_\gamma = 1.33$ MeV.

CLUSTER features	Measured	Calculated
$\epsilon_{ph}\Omega_{Ge} [10^{-4}]$ (at 43 cm)	28.8	31.9
$(P/T)_{un-suppressed}$	0.37	0.40
$(P/T)_{suppressed}$	0.61	0.72
add-back factor	1.44	1.48
1 detector response	0.694	0.675
2 detector response	0.271	0.283
3 detector response	0.033	0.039
4 detector response	0.002	0.003

The full-peak efficiency and the P/T ratio of a single semi-hexagonal tapered Ge detector and a CLUSTER consisting of 7 such detectors have also been calculated as a function of γ -ray energy. The results are shown in Fig. 1.

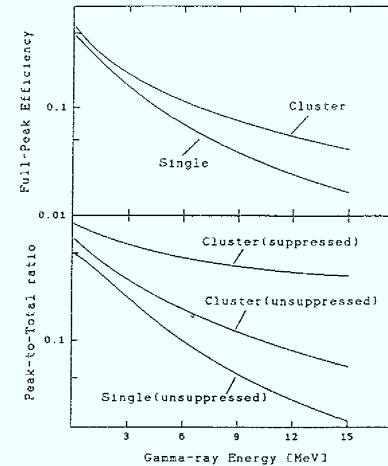


Fig. 1: Calculated full-peak efficiency and P/T ratio of a single semi-hexagonal tapered detector and the cluster of 7 such detectors plotted against the energy of the incident γ -rays.

It can be seen that the increase of the full-peak efficiency and P/T ratio of the CLUSTER detector as compared to those of the single Ge detector gets larger at higher γ -ray energies. For 15 MeV the improvement amounts to a factor of 3. This makes the CLUSTER detector useful also for experiments in which high-energy γ -rays are studied.

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12. EUROPEAN SPALLATION

SOURCE (ESS)

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In order to achieve quite ambitious design goals, for the ESS neutron source it was decided to study a linear accelerator for acceleration of primary particles, H^+ or H^- ions, up to an energy of 1.35 GeV with a time average current of 3.7 mA (i.e. 5 MW), in pulses of 1 μ s duration with intervals of 20 ms.

With tungsten or tantalum as spallation material thermal time-averaged neutron fluxes of more than 10^{14} n/cm²s can be reached, i.e. thermal neutron fluxes during the pulses will be higher than 10^{17} n/cm²s.

The ESS target was proposed as a "rotating double wheel" geometry to facilitate cooling. But disregarding the more progressive idea of a wheel target Conrad /1/ believed that from the conventional solutions the bent plate target is the most promising. The essential features of this proposed target version are:

- bent plates to stabilize the geometry
- axially varying plate thickness in accordance with the varying power density
- radially varying coolant gap width (sickle shaped), coarsely equalizing the cooling conditions

Although it is obvious that a rotating solid state target has its advantage with respect to heat removal, lower neutronic dilution by coolant and lower specific radiation damage such a target will be larger and technically more challenging and probably more expensive than a non-moving target. If the life-time of a stationary target were solely governed by the specific radiation damage of the window, a factor of 50 or more is easily gained employing a rotating target. Thereby an enormous reduction of the number of service shut downs is gained as well. On the other hand, if a rapid exchange were devised, a non-moving target could be an attractive alternative. In order to have a sound basis for a final decision, we investigated the fundamental requirements of a stationary target.

Thermohydraulics of a stationary target

Investigation of the rotating target results for energy deposition, coolant flow, temperature rise, and material damage led to the design of a stationary target system: The stationary bent disc type tungsten target (Fig. 1).

The first calculations had to estimate the energy deposition in the target materials to study the cooling problems and the dynamic stress of the materials. Energy deposition in the targets, moderators, and reflectors is rather sensitive to the exact shape of the proton beam. We assume that the beam is centered on target in a circular beam spot. The beam profile is Gaussian. The standard deviation is $\sigma = 2.5$ cm and the beam is truncated at radius $r = 5$ cm.

We calculated the energy deposition in the target plates along beam axis in radial rings of 1 cm thickness (Fig. 2). The highest energy deposition value of 3800 W/cm² is

reached in the fourth tungsten plate in the inner zone of 1 cm radius.

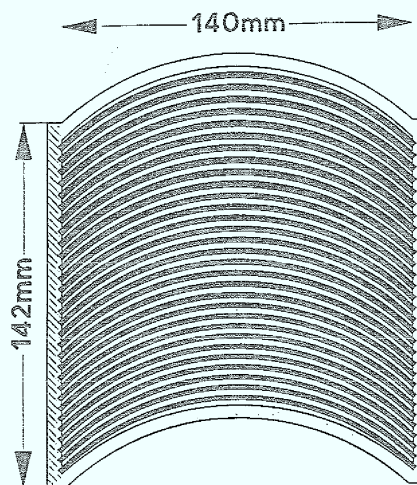


Fig. 1: Vertical cut of target 1 of bent disc target system

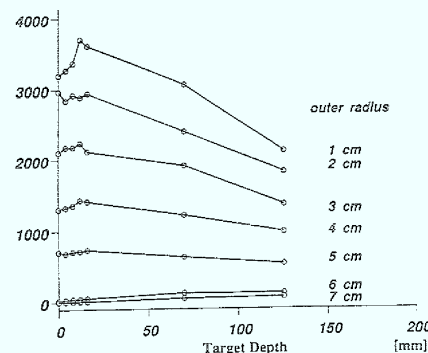


Fig. 2: Distribution of energy deposition in target 1 of bent disc target at 5 MW and 1.35 GeV proton beam

Figure 3 gives the energy deposition per 2 cm² surface of tungsten target plate. The coolant power has to be almost constant along the beam axis but there is still some optimization possible. The investigated surface is of radius 2 cm around beam axis. As result we have got to remove a power of 400 W/cm² from the tungsten surface.

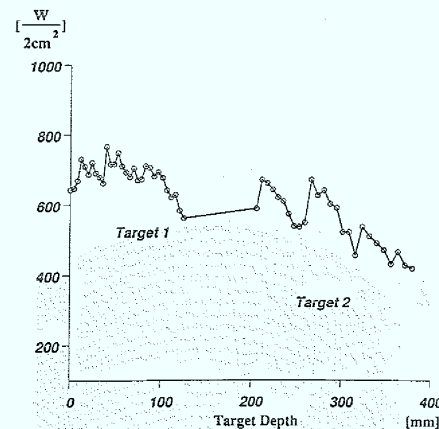


Fig. 3: Energy deposition along central area of radius $r = 2$ cm of bent disc target at 5 MW and 1.35 GeV proton beam

Neutron Flux

Figure 4 gives an overview about thermal neutron flux in a moderator for 1350 MeV. With this energy we calculated the double wheel target, a flat plate target, and a bent disc target and found a step by step decrease of the moderator flux values.

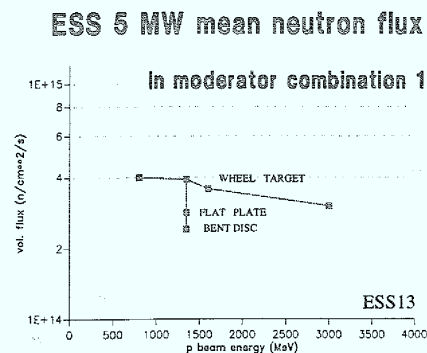


Fig. 4: Thermal neutron flux in a fast moderator for different target systems

To reduce the energy load on the target it is necessary to widen the proton beam. Therefore, it is interesting to know how the neutron fluxes and neutron currents of the moderators are influenced by variation of proton beam diameter. A widening of the proton beam from pencil beam to 14 cm diameter reduces thermal neutron volume flux as well as surface flux by 8%, whereas back scattered thermal neutron current is reduced by only 2% which is not a significant difference to the current from pencil beam. A

widening of the proton beam from pencil beam to 14 cm diameter increases average fast neutron flux (volume flux) by 3% and reduces fast surface flux by 9%, whereas backscattered fast neutron current is increased by 25%. This might be an important increase of background for the measurements.

Calculation of Target Shield Thickness

To perform calculations for the region in and around a spallation target one should use a computer code that has a detailed geometry description and treats all particle types (i.e. protons, pions, neutrons etc.). A code with these features is the Monte Carlo Code HETC /2/. For the calculations of deep penetration problems, on the contrary, one needs a fast running code that can calculate the gamma ray and neutron flux, because the high energy neutrons are the deep penetrating particles and the neutrons and the gammas are the particles which determine the dose outside the shield. A detailed description of the geometry is not necessary for the bulk shielding calculations. So for the shielding calculation the one-dimensional transport code ANISN /3/ can be used.

We decided to combine the advantages of both codes to the coupling procedure /4/, which was used to perform a first estimate of the shielding thickness of the ESS target station. As target we used a system of tungsten cylinders with total active length of 50 cm and radius 5 cm. The target is surrounded by a lead reflector. This model contains the essential features of the proposed wheel target /5/ which is described in the Annual Report 1993. For the shielding calculation the model target region is surrounded by an iron sphere with inner radius 1 m. 10 cm inside this sphere the coupling surface is defined. The calculations were performed for proton beam energy 1350 MeV and different angular intervals. The radial dependence of the dose rate is shown in Fig. 5 for a shield of variable thickness of iron with about 30 cm concrete as outer wall. In all cases the beam power is 5 MW and the dose rate outside the shield is demanded to be less than $7.5 \mu\text{Sv/h}$. The energy dependent flux-to-dose conversion factors are taken from ref. 6. The resulting shield thicknesses are shown in Fig. 6.

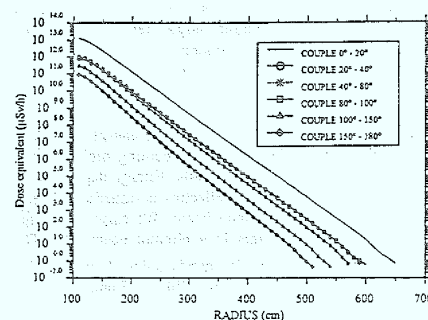


Fig. 5: Dose rates for all angular intervals versus radius of the shielding sphere

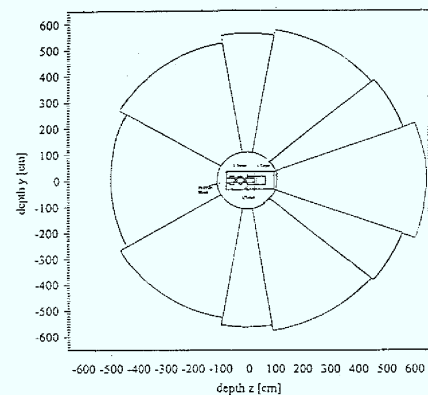


Fig. 6: Resulting shield thickness for different angular intervals to get dose rates less than $7.5 \mu\text{Sv/h}$ outside the shield

Calculation of Radioactivity

For spallation facilities it is important to know the amount of radioactivity caused by beam losses in the accelerator, because this allows to decide whether hands-on-maintenance is possible or remote handling systems have to be foreseen.

To answer this question we calculate the radioactivity in the following way: The residual nuclei produced as result of the spallation process are calculated using the Monte Carlo Code HETC [2]. This code computes the spatial dependence of residual nuclei using a intranuclear-cascade-evaporation model.

The residual nuclei distribution is input for the code ORIHET [7] which solves the equation of nuclide production and decay. The code needs a data base containing half-lives, decay probabilities, binding energies, and gamma energies for all nuclides. So for each time the amount of nuclei, their radioactivity, the power in the system, and the gamma ray spectrum can be calculated.

The cavity walls of 3 mm thickness can be hit by protons lost from the beam. Because of the curved walls it is possible that the beam loss protons can cross the wall several times. So the effective thickness can be much larger than 3 mm. To take this into account and to avoid too many calculations we performed a sensitivity study for the important cavity materials (see Table 1). The aim of this study was to be able to estimate the radioactivity for a given energy and effective thickness using the results of our parameter set.

Target-Material	$1/\Sigma_{\text{geo}}[\text{cm}]$
Steel (X2CrNiMo 17 13 2)	6.57
Ni	7.85
Cu	8.16
Nb	10.62
Al	16.16
Al_2O_3	19.12
Al_2Mg_3	19.85

Table 1: Accelerator structure materials and their geometrical cross sections

At first we calculated the residual nuclei production per source proton for the different materials and for different energies using the so called thin-target-option of HETC. In this case no geometry is used but every incoming proton hits a target nucleus. The results of these calculations are given in Table 2. Then we performed HETC calculations

Target-Material	Incident Proton Energy [MeV]	Residual Nuclei Production per Source Proton
Al	800	0.4459
Al	1350	0.4121
Al	3000	0.3587
Al_2Mg_3	800	0.4674
Al_2Mg_3	1350	0.4685
Al_2Mg_3	3000	0.4554
Al_2O_3	800	0.4378
Al_2O_3	1350	0.4370
Al_2O_3	3000	0.4223
Steel(X 2 CrNiMo 17 13 2)	800	0.5670
Steel(X 2 CrNiMo 17 13 2)	1350	0.5684
Steel(X 2 CrNiMo 17 13 2)	3000	0.5589
Cu	800	0.5818
Cu	1350	0.5841
Cu	3000	0.5430
Ni	800	0.5748
Ni	1350	0.5799
Ni	3000	0.5664
Nb	800	0.6258
Nb	1350	0.6225
Nb	3000	0.6152

Table 2: Residual nuclei production per source neutron for different proton energies (calculated by HETC with thin-target-setup)

for different energies and wall thicknesses and compared the amount of produced nuclei with the values got from the calculations with thin-target-setup. The result is found in Fig. 7, where ζ is the effective wall thickness multiplied by the macroscopic geometrical cross section Σ_{geo} of the target material. These values are given in Table 1. Next step was the calculation of radioactivity with increasing time of irradiation (see Fig. 8). Last step was activity behavior

after shut down of the beam after 1 year irradiation time (see Fig. 9).

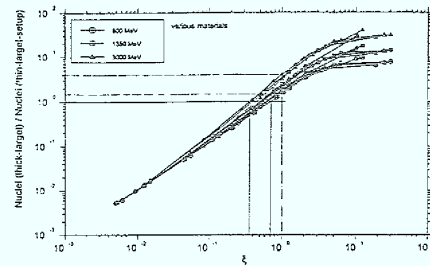


Fig. 7: Ratio of nuclide production rates (thick/thin target) versus relative wall thickness ζ

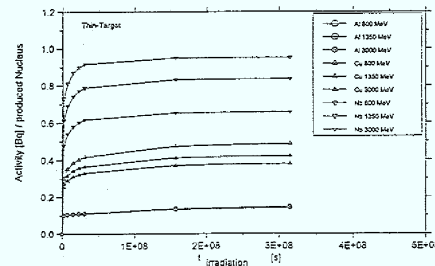


Fig. 8: Induced activity per produced nucleus versus irradiation time

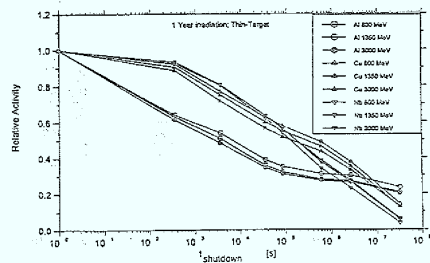


Fig. 9: Activity decrease after shutdown for materials irradiated 1 year

As result of the sensitivity study one finds that for nearly all cases the maximum radioactivity is produced for the lowest proton beam energy and long term irradiation. So one can calculate for an effective thickness d the ζ value and ther. estimate the radioactivity per proton/sec as

$$a_{\text{est}} = r_o \cdot V_{\zeta} \cdot \alpha_{\text{ind}} \cdot V_{\text{dec}}(t_{\text{shutdown}})$$

where r_o = nuclide production rate for 800 MeV (see Table 1)

V_{ζ} = ratio of nuclide production rates for thick/thin-target for abscissa value ζ (see Fig. 7) with $\zeta = d \Sigma_{\text{geom}}$ (from Table 1)

α_{ind} = activity value for 800 MeV and long term irradiation (see Fig. 8)

$V_{\text{dec}}(t_{\text{shutdown}})$ = ratio of activity after t_{shutdown} to activity at $t_{\text{shutdown}} = 0$ for 1 year irradiation time (see Fig. 9)

As an example we estimate the activity per proton/sec for copper wall with effective thickness 4.08 cm and proton beam energy 2200 MeV. The exact calculation using HETC and ORIHET gives for 3 months irradiation time at 10 hours after shutdown

$$a_{\text{calc}} = 0.078 \text{ Bq/s/proton.}$$

The estimated value is $a_{\text{est}} \leq 0.19 \text{ Bq/s/proton.}$

So it is shown that with this method for given beam losses the activity can be calculated using our parameter set.

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from Jan. 25 to Feb. 22, 1994
from April 6 to May 6, 1994
from June 16 to July 14, 1994
from Sept. 17 to Oct. 14, 1994
(Jagellonian Univ. of Krakow, Poland)

P. Moskal (E1)
from April 6 to May 6, 1994
from Aug. 1 to Sept. 2, 1994
from Sept. 19 to 24, 1994
from Nov. 15 to Dec. 15, 1994
(Jagellonian Univ. of Krakow, Poland)

Dr. V. Nelybin (E2)
from May 29 to June 26, 1994
from Nov. 6 to Dec. 11, 1994
(PNPI Gatchina, GUS)

Dr. J. Nemchik (Th)
from April 24 to May 23, 1994
(Slovak Academy of Sciences, Kosice, Slow.)

Dr. M. Nioradze (E2)
from Nov. 21 to Nov. 28, 1994
(HEPITU Tiflis, Georgia)

Dr. B.C. Pearce (Th)
from June 1 to June 30, 1994
(Univ. of Adelaide, Australia)

Dr. A. Petrus (E2)
from Nov. 20 to Nov. 27, 1994
(JINR Dubna, GUS)

Prof. H.S. Plendl (E1)
from April 19 to June 18, 1994
(Florida State Univ., Tallahassee, Florida, USA)

Prof. S. Popov (E2)
from Nov. 20 to Nov. 27, 1994
(Budker Institute Novosibirsk, GUS)

Dr. A. Puzynin (E2)
from Nov. 20 to Nov. 27, 1994
(JINR Dubna, GUS)

Dr. K. Pysz (E2)
from March 21 to March 31, 1994
from April 18 to May 14, 1994
from Sept. 1 to Sept. 30, 1994
from Dec. 1 to Dec. 22, 1994
(Jagellonian University Krakau, Polen)

Dr. Z. Rudy (E2)
from Nov. 16 to Dec. 15, 1994
(Jagellonian University Krakau, Polen)

Dr. T. Rzaca-Urban (E1)
from Oct. 7 to Dec. 6, 1994
(Inst. of Exp. Physics, Warsaw, Poland)

Prof. M.G. Sapozhnikov (E2)
22.11.-27.11.1994
(JINR Dubna, GUS)

Dr. A. Selikov (E2)
from Nov. 6 to Nov. 26, 1994
(Moskau State University, GUS)

Dr. A. Sibirtsev (E2)
from June 19 to July 13, 1994
from Aug. 22 to Sept. 27, 1994
from Nov. 20 to Dec. 11, 1994
(ITEP Moskau, GUS)

Dr. R. Siudak (E1)
from July 15 to Aug. 15, 1994
from Nov. 15 to Dec. 15, 1994
(Jagellonian Univ. of Krakow, Poland)

Dr. J. Smyrski (E1)
from Jan. 1 to Feb. 18, 1994
from April 23 to 29, 1994
from June 1 to Sept. 30, 1994
from Nov. 7 to Dec. 7, 1994
(Jagellonian Univ. of Krakow, Poland)

Prof. A. Strzalkowski (E1)
from Jan. 25 to Feb. 21, 1994
from July 8 to August 4, 1994
(Jagellonian Univ. of Krakow, Poland)

Dr. A. Szczurek (Th)
from Febr. 28 to March 25, 1994
from August 15 to Sept. 14, 1994
(INP, Krakow, Poland)

Dr. V. Tchernychev (E2)
from May 8 to May 16, 1994
from Nov. 20 to Nov. 27, 1994
(ITEP Moskau, GUS)

Prof. A.W. Thomas (AvH-Awardee) (Th)
from Aug. 4 to Aug. 25, 1994
(Univ. of Adelaide, Australia)

Prof. D. Toporkov (E2)
from Nov. 20 to Nov. 27, 1994
(Budker Institute Novosibirsk, GUS)

Dr. S. Trusov (E2)
from May 9 to July 8, 1994
from Nov. 20 to Dec. 2, 1994
(JINR Dubna, GUS)

Dr. R. Tzenov (E1)
from Nov. 26 to Dec. 9, 1994
(University of Sofia, Bulgaria)

Dr. R. Tzolov (E1)
from May 8 to June 15, 1994
(Bulgarian Academy of Sciences)

Prof. T. Udagawa (Th)
from May 10 to June 30, 1994
(Univ. of Texas at Austin, USA)

Dr. T. Venkova (E1)
from March 2 to May 30, 1994
from Oct. 14 to Dec. 13, 1994
(Bulg. Acad. of Sciences, Sofia, Bulgaria)

M. Wojciech (GG)
from Dec. 5 to Dec. 23, 1994
(Jagellonian University Krakau, Polen)

Dr. B.G. Zakharov (DFG-fellow) (Th)
from Sept. 1 to Oct. 31, 1994
(Landau Inst. for Theor. Phys., Moscow, Russia)

Prof. D. Zawischa (Th)
from July 4 to July 8, 1994
(Univ. Hannover)

Dr. M. Ziolkowski (E1)
from Jan. 15 to 29, 1994
(Jagellonian Univ. of Krakow, Poland)

Dr. V. Zoller (DFG fellow) (Th)
from Nov. 1 to Dec. 24, 1994
(ITEP, Moskow, Russia)

DP P. Zolnierczuk (E1)
from Jan. 13 to Dec. 31, 1994
(Jagellonian Univ. of Krakow, Poland)

Dr. I. Zychor (GG+E2)
from Aug. 1 to Aug. 31, 1994
from Nov. 16 to Dec. 14, 1994
(Soltan Institute Swierk, Polen)

Technical and Administrative Staff:

P. Birx (LI)
H.G. Böge (LI)
W. Borgs (E2)
H. Borsch (LI)
R. Brings (LI)
G. Brittner (Ec)
P. Brittner (LI)
E. Brökel (Ec)
U. Burck (Ad)
J. But (Ws)
L. Conin (LI)
B. Dahmen (LI)
W. Derissen (Cd)
N. Dolfus (Ec)
F. Drees (LI)
R. Enge (LI)
J. Engel (LI)
P. Engels (LI)
W.R. Ermer (E2)
W. Ernst (Ec)
K. Esser (Ad)
H.P. Faber (LI)
G. Fiori (Dt)
H.-W. Firmenich (Ws)
G. Gad (LI)
D. Gehsing (LI)
S. Geisler (Cd)
G. Göbbels (Ws)
M. Gringel (Cd)
Ch. Haberbosch (LI)
H. Hadamek (Ws)
A. Hamacher (Dt)
J. Wimmer (Sd)
T. Hillenberg (Cd)
H. Hintzen (Rp)
W. Hoffmeister (LI)
M.G. Holona (Ws), since June 13, 1994
K.D. Jach (LI)
H.M. Jäger (E1)
H.J. Jansen (Ws)
R. Janssen (Ad)
A. Jeglorz (E2), until April 30, 1994
M. Karnadi (Da)
R. Klein (E1)
S. Köhn (Da)
K. Krafft (Rp)
M. Kremer (Ws)
G. Krol (LI)
K.P. Kruck (LI)
M. Küven (Ws), since Aug. 8, 1994
K.G. Langenberg (LI)
W. Lorenz (Ad)
G. Lürken (Ec)
H. Metz (Dt)
A. Müller (LI)
M. Müskes (LI)
R. Nellen (Da)
S. Papaspyrou (Ad)
J. Pfeiffer (Dt)
H. Pütz (LI)
A. Retz (Cd)

A. Richert (LI)
U. Rindfleisch (Cd)
DI D. Rosin (LI)
N. Rotert (LI)
D. Ruhrig (LI)
Th. Sagefka (LI)
M. Schaaf (LI)
F. Scheiba (LI)
P. Schmidt (LI)
Jos. Schmitz (Ws)
Jürg. Schmitz (LI)
H. Schneider (LI)
H. Schröder (Da)
F. Schultheiß (Ws)
R. Seidemann (E2)
M. Simon (LI)
H. Singer (LI)
K. Sobotta (LI)
W. Spieß (LI)
D.W. Spölgén (Ws), since May 16, 1994
C. Strauch (Ad)
J. Strehl (Ws)
E. Tesch (Ad)
K.P. Wieder (E2)
K. Winkler (Ec)
DI J.-D. Witt (LI)
H.W. Zens (LI)

(E1) Institute for Experimental Nuclear Physics 1
(E2) Institute for Experimental Nuclear Physics 2
(Th) Institute for Theoretical Nuclear Physics
(LI) Large Nuclear Physics Instruments
(Ad) Administration
(Cd) Construction and Design
(Da) Data Acquisition Group
(Dt) Detector and Target Laboratory
(Ec) Electronics
(Rp) Radiation Protection
(Sd) Simulation and Detector Development
(Ws) Mechanical Workshop

IX. PUBLICATIONS

IKP-100194
Abrahamson, A.; Källberg, A. et al.*; Meuth, H.; Schnase, A.
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Altman, J.; Weil, D.; Sala, P.; Luig, M.; Eschenauer, M.; Albers, S.; Dewald, A.; Zell, K.O.; v. Brenlano, P.; Gast, W.; Lieder, R.M.; Bazzacco, D.; Rossi-Alvarez, C.; de Poil, M.; Maron, G.; Rico, J.; Bedovalto, G.; Biasi, N.; Lo Bianco, G.
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IKP-100494
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IKP-100594
Anagnostopoulos, D.; Borchert, G.L.; Lenz, S.
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Phys. Lett. B, in print
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Arneodo, M.; Arvidson, A.; Badelek, B.; Ballintijn, M.; Baum, G.; Beaufays, J.; Bird, I.G.; Björkholm, P.; Botje, M.; Brogini, C.; Brückner, W.; Brüll, A.; Burger, W.J.; Ciborowski, J.; van Dantzig, R.; Dyring, A.; Engelen, H.; Ferrero, M.I.; Fluri, L.; Gaul, U.; Granier, T.; von Harrach, D.; van der Heijden, M.; Heusch, C.; Ingram, Q.; Janson-Prytz, K.; de Jong, M.; Kabuś, E.M.; Kaiser, R.; Ketel, T.J.; Klein, F.; Kullander, S.; Landgraf, U.; Lindqvist, T.; Mallot, G.K.; Mariotti, C.; van Middelkoop, G.; Milsztajn, A.; Mizuno, Y.; Most, A.; Mücklich, A.; Nassalski, J.; Nowotny, D.; Oberski, J.; Paic, A.; Peroni, C.; Povh, B.; Rieger, R.; Rith, K.; Röhrich, K.; Rondio, E.; Ropelewski, L.; Sandacz, A.; Sanders, D.; Scholz, C.; Seitz, R.; Sever, F.; Shibata, T.-A.; Siebler, M.; Simon, A.; Stalano, A.; Szleper, M.; Tzamouranis, Y.; Virchaux, M.; Vuilleumier, J.L.; Walcher, T.; Windmolders, R.; Witzmann, A.; Zetsche, F.
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- IKP-101994
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Jüli-3035
February 1995
ISSN 0944-2952