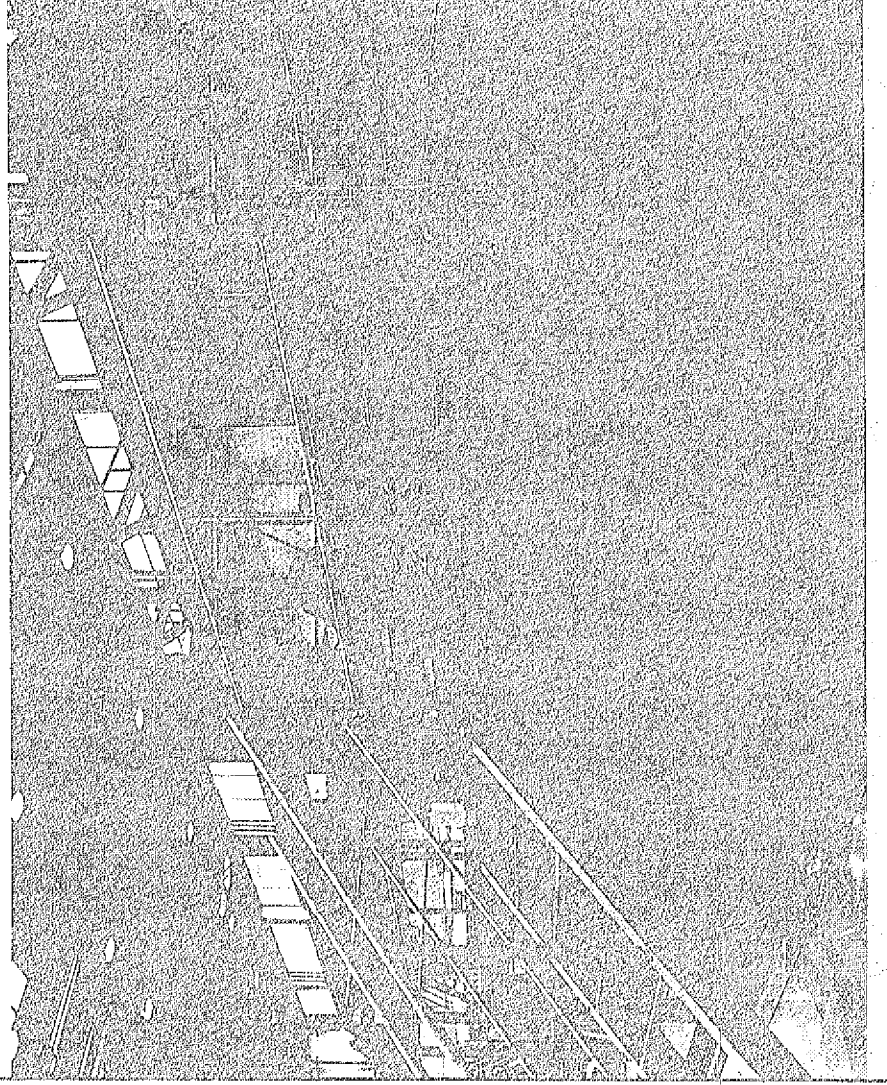
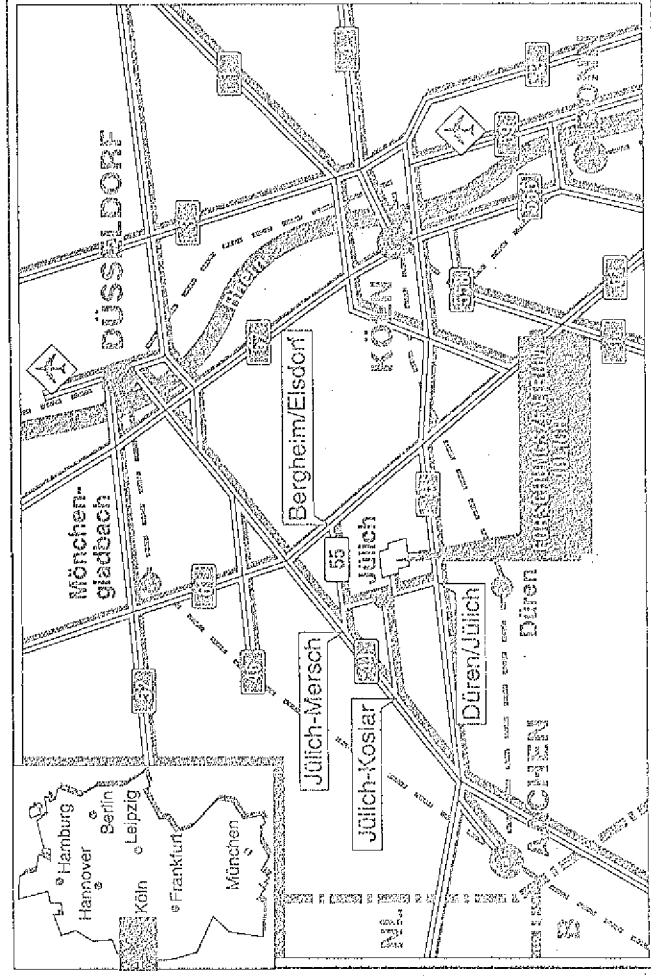


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Pictured on the front cover is the electron cooler at its test site located in the accelerator ring's inner area. It will "cool" the circulating protons to obtain a highly uniform beam. This central part of the accelerator ring is the most distinguishing component in comparison to preceding accelerator generations. It allows to shape the beam quality such, that new experimental techniques will be feasible. Its prominent role even stimulated the euphemistic acronym COSY, which was derived from Cooler-Synchrotron. At its present position the electron cooler will undergo a thorough test of all its components before its final insertion in the ring. A detailed report is given on page 223.

THE UNIVERSITY OF CHICAGO

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During the year 1991 the work of the members of our *Institut für Kernphysik* (IKP) in the *Forschungszentrum Jülich* (KFA) concentrated on building the Cooler Synchrotron COSY-Jülich in close cooperation with other groups in the KFA, and on preparation of experiments to be carried out at COSY. Especially the latter work is being done in very close collaboration with scientists of CANU and together with groups and researchers from laboratories in foreign countries. We also have very strong cooperation with groups from the Research Center Rossendorf near Dresden. The experimental activity is accompanied by theoretical studies in the field of medium energy nuclear physics.

The emphasis of our work necessarily aimed at the completion of COSY, which is expected to be available for experiments in 1993. Recommissioning of the injector cyclotron JULIC has started. The injection beam line to COSY has been set up. All COSY dipole magnets have been located and aligned. The quadrupole magnets have been delivered after a major delay. A large fraction of them has been tested and installed in the ring. Nearly all correction magnets have been delivered. The injection and extraction elements are available, also the magnets for the lines of the extracted beams. A section of the vacuum tube of the ring has been mounted and tested for ultra high vacuum. The electron cooler has been assembled for tests prior to installation in the ring. COSY control has aided in testing and commissioning of various components.

The preparation of COSY experiments concerns the realization of multipurpose facilities and other experiments. The large spectrometer BIG KARL for ejectiles with momenta up to 1.1 GeV/c has been moved into its new location. Upgrading of the entrance quadrupoles has been started. For other studies with extracted beams the large time of flight spectrometer TOF is under construction. Parts of its detectors are being built by CANU groups, also driftchambers for various experiments. Position sensitive solid state detector systems are being developed and built for experiments with BIG KARL and TOF. The main parts of a cluster target have been developed and assembled for experiments at the circulating COSY beam. The design of the 0^0 -facility for a number of internal target experiments and ejectiles with momenta up to the maximum in COSY has been optimized regarding the detectors, resolution, time of flight and ΔE particle discrimination. Research and development work was continued for the high-rate photon spectrometer and for the other COSY experiments. Detectors and equipment was acquired, built and tested. Studies have begun to make use of the COSY beam also for cancer therapy.

Work going on in the field of theoretical nuclear physics is strongly connected with research projects at COSY and processes induced in the proton-antiproton interaction.

Many-body methods developed in nuclear physics have been applied to understand the structure of the rho meson in free space as well as in the nuclear medium. The shape and the resonance energy of the rho meson is strongly modified in the medium which might be detected in dilepton experiments.

Scientists from the IKP have also taken part in research carried out in other laboratories. At LEAR/CERN the PS185 experiment $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ gave for the first time also information about $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\gamma$ and $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\pi^0$ channels. New evaluation techniques allow for very precise spin correlation data in $\bar{\Lambda}\Lambda$. A dedicated run for $\bar{p}p \rightarrow \Sigma^+\Sigma^+$ occurred in November. The PS202 $\bar{p}p \rightarrow 4K$ experiment accumulated data at various beam energies. A surprisingly large fraction of the 4K data was found to come from the $\bar{p}p \rightarrow \phi\phi$ two body channel. Preparations have been made for studies of meson production in pd reactions at CELSIUS/Uppsala.

Through nuclear spectroscopy two-phonon octupole excitation had been identified in ^{148}Gd . Now also the $2 \rightarrow 1$ phonon transition strength could be determined in a measurement with the NORDBALL detector. With this spectrometer also the deexcitation of the superdeformed band in ^{143}Eu has been investigated. With the OSIRIS spectrometer superdeformed bands were studied in ^{146}Gd and its neighbours.

The ISIS ECR source has been used for the production of ion beams for atomic physics research. A large part of our work is being carried out by our young colleagues, Diplomanden and Doktoranden, our engineers and technicians, the IKP service group and by members of service groups and other institutions of our center. I am very grateful for this help.

Our institute and our research center receive invaluable help through the advice from the members of the IKP Beirat, of the Program Advisory Committee, of the Committee "Experiment-Beschleuniger", and of the Machine Advisory Committee. As managing director of the IKP, I would like to express my sincere thanks to all members of these committees for their advice and to the board of management of the KFA for the support that is given our institute.

Alto Schult

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CONTENTS

I. EXPERIMENTAL NUCLEAR PHYSICS.....	1
1. MEDIUM ENERGY PHYSICS	
1.1. PREPARATION OF COSY EXPERIMENTS	
Progress on the Development of the Time-of-Flight Spectrometer.....	3
Technical development of the TOF-detector.....	4
A Start Detector System for the Investigation of the pp-Bremsstrahlung with the COSY Time-of-Flight Spectrometer (TOF).....	5
A Liquid Hydrogen Target with very Thin Windows.....	7
The Induction Drift Chamber (IDC) at High Rate.....	8
The Use of Induction Drift Chamber Signals as a Fast Trigger.....	10
Construction of a Vertex Detector Consisting of Independent Induction Drift chambers (IDC).....	12
Online Program for Data Acquisition with a Flash-ADC System.....	14
A Data Acquisition System for an Induction Drift Chamber.....	15
Current Controlled Delay Unit for Coincidences in the Scintillator Hodoscope.....	16
Monte Carlo Simulation of the COSY Time-of-Flight Detector (TOF).....	18
Feasibility Studies for a TOF- γ -Detector at COSY.....	20
Concept of the 0^0 Facility at TP2 in COSY.....	21
Optimization of the particle-flight paths in the 0^0 Facility.....	24
Monte Carlo Simulations for the 0^0 Facility.....	26
Subthreshold K^+ Production and the Fragmentation of the Residual Nucleus.....	28
Deuteron Momentum Spectra accompanied by Subthreshold K^+ Production.....	29
Model Predictions of K^+ Momentum Spectra and Angular Distributions from proton- ^{12}C Collisions at $E_p = 950$ and 1400 MeV.....	30
Preparation for the COSY Experiment to produce very heavy Λ -Hypernuclei and to measure their lifetime.....	32
Construction and Tests of Avalanche Detectors.....	34
Simulation of the Passage of Heavy ^{239}Th Hypernuclei through a thin ^{238}U Target.....	36
Time Structure of Scintillators for the COSY Two-Arm Photon spectrometer.....	39
COSY-11 Detector Arrangement.....	41
Reconstruction of the Shape of Meson Resonances Produced in the pp $\rightarrow$ ppM Reaction Near Threshold.....	43
A Fastbus System for Preparation of the Experiment COSY-11.....	45

The Experiment COSY 10 (Pionic Atoms): A Status Report.....	46
A Study of the $p+D \rightarrow \eta+^3\text{He}$ Reaction with a 4π Detector Geometry.....	48
Monte Carlo Simulations of Energy Dose Profiles of Protons in Water for Studies of Possible Medical Applications of the COSY Beam.....	50
1.2. EXPERIMENTS AT EXTERNAL FACILITIES	
Spin Correlation in $\bar{\Lambda}\Lambda$ -Production.....	52
Feasibility of CP-Invariance in $\bar{\Lambda}\Lambda$ Hyperon Decay.....	53
Missing Mass Criteria in the PS185 Data Analysis Program.....	54
A Trackfinding Procedure on the Basis of Event Rotation.....	55
Physics at JETSET.....	56
Performance of the Circular Scintillator Hodoscope in PS202.....	58
Determination of Luminosity with the Pixel Trigger of the Circular Scintillator Hodoscope.....	60
The JETSET Trigger System for the Circular Scintillator Hodoscope.....	61
Correction of the JETSET Experimental Data by Monte-Carlo Simulation.....	63
Heavy Hyperon-Antihyperon Production.....	64
Performance Study of Doubly Focussing Crystals for a Low Energy Bragg Spectrometer.....	66
γ -Decay of the Giant Dipole Resonance in Hot Rotating Rare Earth Nuclei.....	67
Radial Excitation of the Nucleon to the $P_{11}(1440 \text{ MeV})$ Resonance in Alpha-Proton Scattering.....	69
ZEUS Barrel Calorimeter Assembly and Testing.....	72
Development of a Vectorized Monte Carlo Code for the Simulation of Particle Transport in the HERA-ZEUS Calorimeter.....	74
Neutron Spectra Measurements and Monte Carlo Simulations of (p,xn) Reactions at 597 MeV and 800 MeV.....	75
Simulation Experiments for Planetary Gamma-Ray Spectroscopy by Means of Thick Target High-Energy Proton Irradiations.....	77
A Meson Production Experiment in Uppsala.....	79
2. NUCLEAR SPECTROSCOPY	
Is there Shape Coexistence in ^{99}Zr ?.....	81
Nilsson Parameters at $N=61$ and Particle-Rotor Calculations for ^{103}Mo	83
Mass-, Charge-, and Angular Momentum Transfer in $^{106}\text{Cd} + ^{54}\text{Fe}$ Collisions studied by $\gamma\gamma$ -Coincidences.....	85
Study of the Decay out of a Superdeformed Band in ^{143}Eu	88
Search for Superdeformation in $^{144} \text{ } ^{145}\text{Gd}$	90

Excited Superdeformed Band in ^{146}Gd	92
The Two— to One—Phonon E3 Transition Strength in ^{148}Gd	94
Energy Inversion of the $f_{7/2}$ and $h_{9/2}$ Neutrons in Yrast Configurations of $^{154}\text{Yb}_{84}$	96
Study of High—Spin Correlations in ^{180}Os	98
Search for Superdeformation in ^{190}Au	100
High Resolution Measurement of K X—rays of Lanthanum Following the Collisions with 403 MeV N^{6+} Ions.....	101
Development of a Composite Ge Detector for EUROBALL.....	103
EUROBALL Design Geometry for 72 Detectors.....	106
The Encapsulated Ge Detector.....	107
Digital Approach to High—Resolution Pulse Processing.....	108
New Approach to High Resolution Multiparameter Data Acquisition.....	112

3. NUCLEAR REACTIONS

Elastic Transfer in $^{11}\text{B} + ^{12}\text{C}$ System at 344.5 MeV (Lab).....	118
Inelastic Scattering to the Continuum in Collisions of ^{12}C ions with ^{12}C , ^{58}Ni and ^{90}Zr Nuclei at 344.4 MeV (LAB).....	119
Two Photon Correlations from $^{16}\text{O} + ^{24}\text{Mg}$ Interactions at 33 MeV $\cdot$ A.....	121
Fragment—Fragment Correlations at Fermi Velocity.....	122
Intermediate Mass Fragment Emission from $^{32}\text{S} + ^{197}\text{Au}$ at 33 MeV $\cdot$ A.....	123

II. THEORETICAL NUCLEAR PHYSICS.....125

4. MEDIUM AND HIGH ENERGY PHYSICS

The pion electromagnetic form factor.....	127
Form factor of the σ —term in a meson exchange model.....	129
Effects of final—state interactions in the $\bar{p}p \rightarrow \pi\pi$ annihilation process.....	131
Can NN data tell us something about the nature of the NN repulsion?.....	133
The Nucleon—Nucleon Potential in the Skyrme Model: Beyond the Product Approximation.....	135
Improved Boundary Condition, Coupled Channel Model: pp Predictions to 1 GeV.....	136
The $\Lambda\text{N} - \Sigma\text{N}$ Interaction with Isobar Coupling.....	139
$\pi\pi$ —Scattering in Nuclear Matter.....	142
Towards a consistent description of NN scattering and annihilation.....	145
$\text{NN} \rightarrow \pi\pi$ scattering amplitude in the pseudophysical region in the TOPT framework.....	147

Antinucleon nucleon scattering.....	149
New approach to vector mesons in the Sullivan processes.....	151
A schematic model for the selfenergy of a non-local Lagrangian.....	152
Renormalization of the ρ -meson mass due to tad poles.....	154
Geometrical picture of coherent γ - γ and γ -g processes in ultrarelativistic heavy ion collisions.....	157
Coherent γ - γ lepton pair production.....	158
Coherent electron positron production in π^+p collisions.....	160
The spin-longitudinal/spin-transverse response of ^{40}Ca to 500 MeV protons.....	162
Analysis of the $t_{NN,N\Delta}$ -amplitude by the $(\vec{p},\vec{n})$ -reaction.....	164
Comparison of different models for the ΔN^{-1} -residual interaction in the Δ -propagation in nuclei.....	166
Pionabsorption on ^{12}C in the Δ -resonance region.....	168
Photo-absorption on ^{12}C in the Δ -resonance region.....	169
Density dependent rescattering effects in the $(e,e'p)$ reaction.....	171
Coupling between $(e,e'p)$ and $(e,e'n)$ channels in the quasi-elastic domain.....	173
Δ -Excitations in Nuclei by the $(\vec{p},\vec{n})$ -reaction.....	175
Exclusive cross sections in (h,t) reactions to the Δ -excitations in nuclei.....	177
Rescattering contribution to proton-proton bremsstrahlung in a relativistic approach.....	180
Spin observables in pp bremsstrahlung and off-shell effects of the NN interaction.....	183
On the relativistic spin correction in proton-proton bremsstrahlung.....	186
A comparison of nucleon-nucleon potential models in proton-proton bremsstrahlung.....	189
Hard photon production in proton-deuteron reactions at intermediate energies.....	192
Consistent two-phase model of the nucleon.....	195
On the role of the vertex form factor for deep-inelastic lepton-nucleon scattering.....	197
Consequences of SU(2) flavor symmetry breaking in the nucleon sea for the Drell-Yan processes.....	200
Gottfried sum rule violation.....	201

5. NUCLEAR STRUCTURE AND FORCES

Calculations of E1 resonances in ^{40}Ca , ^{48}Ca and ^{208}Pb in the 1p1h + 2p2h + continuum approximation.....	203
Microscopic Calculation of the Giant Dipole Resonance in Hot Deformed Nuclei with Realistic Forces.....	207
Low-lying 1^+ levels in spherical open shell nuclei.....	209
E0 and M1 excitations in ^{56}Ni	211

Coulomb dissociation of loosely bound neutron rich nuclei and higher order electromagnetic effects.....	213
Coulomb dissociation of ^{11}Li and ^{11}Be in a direct fragmentation model.....	214
Coulomb dissociation as a tool of nuclear astrophysics.....	215
Semiclassical and Glauber theories of electromagnetic excitation with high energy heavy ions.....	216
Vacuum polarization in sub-Coulomb ^{12}C — ^{12}C Scattering.....	217

III. ACCELERATOR DIVISION.....219

6. COOLER SYNCHROTRON COSY

Alignment of the COSY—Ring.....	221
Assembling of the COSY—Vacuum System.....	227
Status of the COSY Electron Cooler.....	228
The Stochastic—Cooling System of COSY.....	226
Power Supply Status and First Tests on Original Load.....	229
Stripline Unit.....	232
Beam Position Monitors.....	233
Wall—Current Monitor.....	236
Schottky—Pickups.....	238
Schottky Beam Position Monitor.....	240
Phase—Switching of the COSY ($h=1$) Cavity for the γ —Transition.....	241
A Digital Amplitude Demodulation Circuit for Signals in Quadrature.....	244
An Analytical Approach to Transverse Tune Measurement by Continuous Wave Excitation.....	247
The COSY Control System.....	250
Polarisation in COSY.....	252
COSY Stripping Target.....	254
COSY Beams to the External Target Stations 2, 3 (TOF) and the BIG KARL Spectrograph.....	255

7. ISOCRONOUS CYCLOTRON AND ISIS SOURCE

Commissioning of the new RF—System at the Cyclotron.....	256
Measurement of the Magnetic Field of the Cyclotron for the H_2^+ —Beam (40 MeV/A).....	259
The Injection Beam Line.....	262
Beam Steering in the Injection Transfer Line.....	263
Measurement of the COSY Injection Energy.....	264

Beam Profile Measurement in the COSY—Extraction.....	264
Ion Source Status.....	265
8. SPECTROMETER BIG KARL	
Magnetic Spectrometer BIG KARL.....	266
Studying the Ion—Optical Bandwidth of the QQDDQ—System.....	268
9. RADIATION PROTECTION	
Radiation Protection.....	271
Computer Aided Shield Layout at COSY.....	272
<u>IV. TECHNICAL DEVELOPMENTS</u>	273
10. COMPUTER DEVELOPMENT	
Computer Development.....	275
COSY Experiment Network, IKPnet.....	275
COSY Experiment Data Processing.....	276
FERA System Controller, FSC.....	278
Testing EXABYTE Tape Drives.....	279
Feasibility Test to Use RISC Workstations for COSY Offline Data Analysis and Reconstruction.....	282
11. DETECTORS AND TARGETS	
Semiconductor Detectors.....	283
Targets.....	284
<u>V. SCIENTIFIC ADVISORY COUNCIL OF THE INSTITUTE OF NUCLEAR PHYSICS</u>	287
<u>VI. EXTERNAL ADVISORY COMMITTEES AT COSY</u>	287
<u>VII. PERSONNEL</u>	288
<u>VIII. PUBLICATIONS</u>	295
<u>IX. CONFERENCE CONTRIBUTIONS, TALKS</u>	302
<u>X. PATENTS</u>	317
<u>XI. INTERNAL REPORTS</u>	317
<u>XII. INTERNAL TALKS</u>	318
<u>XIII. INDEX OF AUTHORS</u>	319

I. Experimental Nuclear Physics

THE HISTORY OF THE

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BY

JOHN BURNET

OF

SCOTLAND

IN

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CHARLES THE FIRST

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SEVEN VOLUMES

Progress on the Development of the Time-of-Flight Spectrometer

H. Brand, S. Brand, H. Freiesleben, A. Fritsche, J. Krug, E. Kuhlmann, S. Lange,
H. Matthäy, P. Ringe, A. Röser, M. Steinke, R. Werding

for the TOF collaboration

The basic setup of the time-of-flight spectrometer (TOF) covering 4π (in the CMS) has been outlined in the previous annual report. In this contribution some comments are given with respect to only the various layers forming the stopdetector components; for information regarding the startdetector and target see the following articles.

The investigation of the timing characteristics of the long scintillator strips of BC408 has been continued. The effective light velocity within the scintillator was found to be 16.0 ± 0.3 cm/ns in good agreement with published data [1]. This value, which is an important input parameter for the Monte Carlo simulations, differs significantly from $c_{\text{medium}} = c_{\text{vac}}/n = 18.97$ cm/ns.

The outcome of these tests has led to some decisions regarding electronics and equipment: XP2020 PMT's, together with Philips S563 bases will be used; the first 100 units have been purchased. Likewise the first order for bent stripes of BC408 scintillators as part of the outermost layer of one of the barrels has been placed with the manufacturer. Before ordering a whole set, a critical inspection of geometrical tolerances and a rigorous performance test is planned; work on the test equipment is under way.

The scintillators in the final setup will be placed in vacuum of order 10^{-3} mbar, which puts additional weight on the selection of suitable wrapping material. The standard procedure — wrapping each strip in a layer of first crumpled aluminium and then blackened or aluminized mylar — cannot be applied since unavoidable air-bags will pop-up once the pumping starts. Instead it is planned to place between the various scintillator layers suitably formed sheets of aluminized polyethylene or polyester-foils, thus yielding sort of a sandwich structure. Here "suitably formed" means having sheets which are either pressed into a corrugated cardboard like manner or in the way of an "egg-carton". To keep the energy losses of the traversing particles within these "blind" layers as low as

possible, a thickness of order $30 \mu\text{m}$ is desirable. On the other hand to arrive at a stable and partially elastic structure — with amplitudes (difference of peak to valley) and wavelengths (distance of individual peaks or ridges) of order 4 and 12 mm, respectively — a thickness higher than $100 \mu\text{m}$ should be chosen. Tests in collaboration with outside manufacturers are under way to find an optimum solution to this problem.

Apart from the development of the main Monte Carlo simulation code a smaller version has been set up to yield further information on the performance of the detector system. Here the main findings are:

- i. with about a chance of 2 % the entrance of a given trajectory in a scintillator strip belongs to a fictitious pixelelement that differs from the one of the corresponding exit;
- ii. demanding for each track a signal from all three layers results in a loss of about 1 % (for $pp \rightarrow pp \gamma$ at 1.0 GeV/c);
- iii. the reconstruction efficiency of the detector consisting of the endcap and one 2.5 m long barrel was determined from just the "pixel"-information alone that is which scintillator elements have seen a hit; it was found to be larger than 97 % for multiplicities up to 4. Further improvement is expected by considering timing as well as dE/dx information.

Reference:

1. A. Peghaire et al., NIM A295 (1990) 315 and references therein.

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Technical development of the TOF-detector

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K. Kilian, R. Klein, U. Rindfleisch, H. Sachsenhausen*, P. Turek
for the COSY TOF Collaboration

The basic design of the TOF detector (Time-of-Flight) was already described in Ref. /1/. Since that time several further developments and optimization of the technical components have taken place. In cooperation with the KFA division ZAT a detailed design of the vacuum vessel and the support

structures for the scintillator hodoscopes was presented, ready for construction. Fig. 1 shows an overview drawing of the complete TOF detector design with vacuum vessel and the scintillator hodoscopes. The trigger system with target and start detectors is also suggested.

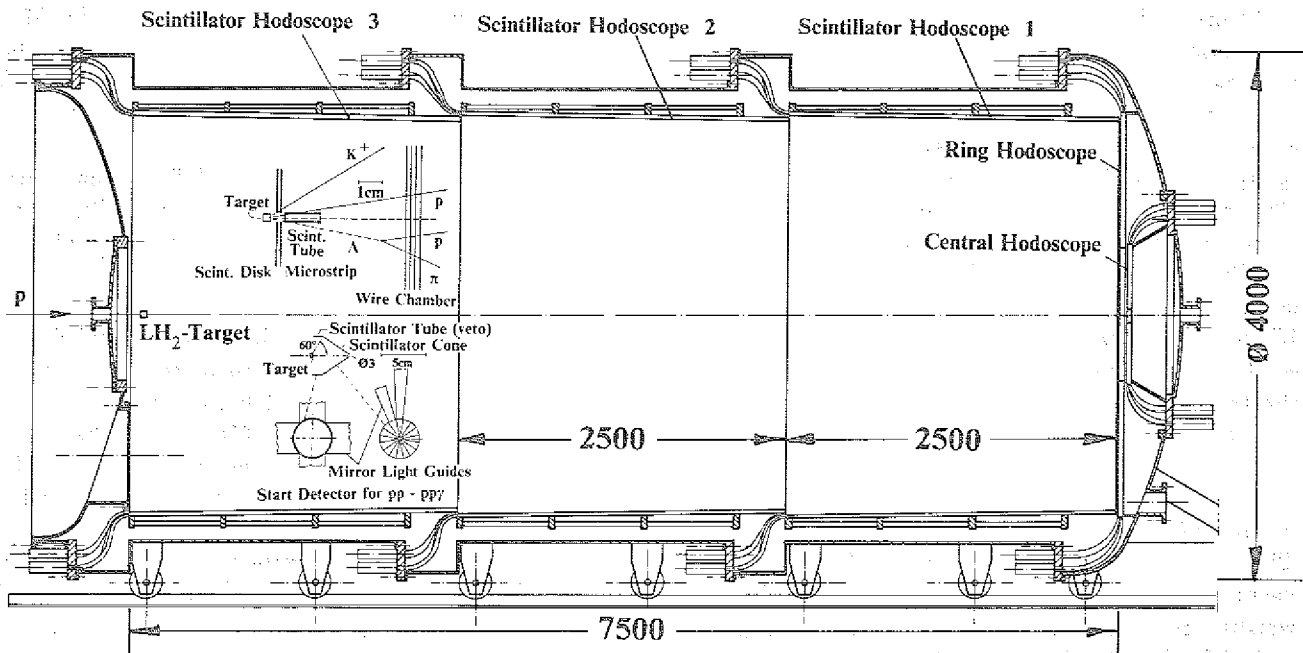


Fig. 1: Overview drawing of the complete TOF detector design

The TOF detector as shown in Fig. 1 consists of the following components.

1. The five three-layer scintillator hodoscope sections
 - Central hodoscope (Quirl) with an outer diameter of 1.08 m.
 - Ring hodoscope with an inner diameter of 1.0 m and an outer diameter of 3.08 m.
 - Three barrel hodoscopes 1-3, with length of 2.50 m each and a inner diameter of 3.15 m.
 - Five scintillator support structures
2. Vacuum vessel
 - Front cap with the 1 m diameter flange
 - End cap flange system
 - Three barrel vessel sections connected by the light flanges
3. Five light guide flanges, which support light guide and photomultiplier mounting

The main parts of effort, which were necessary to achieve a complete design are the following.

- The design of the front cap with a 1 m diameter flange was changed for bidirectional use.
- Stress and rigidity calculations were performed for the components of the vacuum vessel.
- Design of the support structure for the barrel sections.
- Design of a gantry for positioning of the barrel hodoscopes during mounting.
- Efforts to fix the precise dimensions of the barrel scintillators.
- Partial construction of the central hodoscope.
- Design and test of two different trigger systems (contributed by the University of Erlangen and ZfK Rossendorf).
- Construction and test of the liquid hydrogen target.

Reference:

/1/ Annual Report IKP 1990, Jül-2462

* KFA-ZAT

A Start Detector System for the Investigation of the pp-Bremsstrahlung with the COSY Time-of-Flight Spectrometer (TOF)

K. Kilian, P. Michel^a, B. Miethig^b, K. Möller^a, B. Naumann^a, L. Naumann^a, G. Schmidt^a
for the TOF-collaboration

For the measurement of the pp-Bremsstrahlung cross section by means of the time-of-flight spectrometer TOF [1] a start detector is needed which has to meet the following requirements:

- time resolution for minimum ionizing particles < 0.5 ns;
- segmentation (2-fold in θ -direction, 16-fold in φ -direction) to enable the registration of a sufficiently large pulse rate (10^6 /s) and to give a rough information on the start-off-direction of the ejectiles;
- detector mass as low as possible (< 50 mg/cm²) to avoid background and secondary reactions in the detector material.

We propose a start detector system (fig. 1) which is a reasonable compromise to meet these requirements. Two rings each with 16 trapez-like plastic scintillators are arranged in such a way to cover the angular range from $\theta = 5^\circ$ – 60° and 60° – 90° .

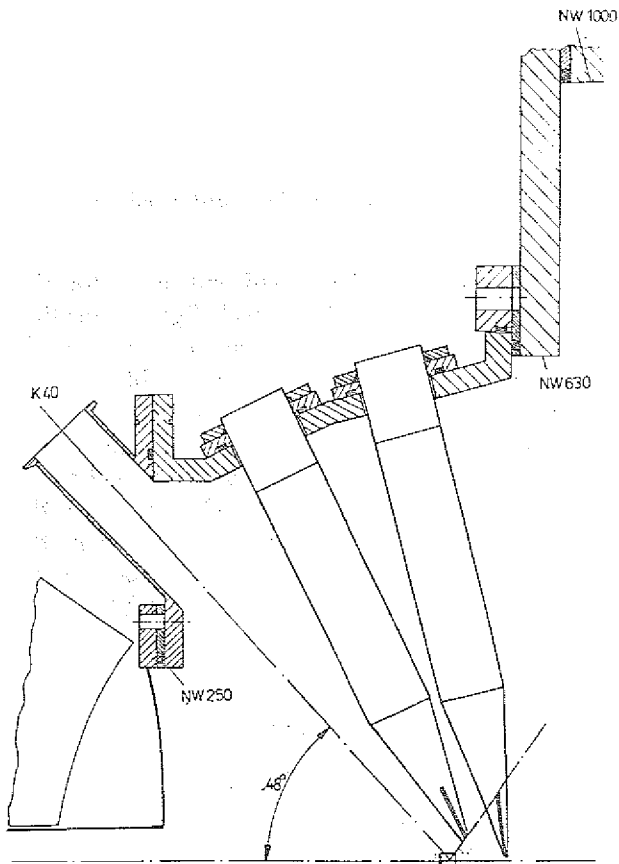


Figure 1: Start detector for the investigation of pp-Bremsstrahlung

To guarantee optimal transfer of light from the scintillators to the photomultipliers, to keep scintillator thickness small and to avoid Cerenkov-light emission the light read out from the scintillators is performed by air light guides with very thin walls. The whole system is placed within a vacuum vessel attached to the entrance flange of the TOF.

The test air light guides consist of nickel with a wall thickness of $6 \mu\text{m}$ made by a galvanoplastic procedure. The interior of the light guides is covered by a thin reflecting layer.

To study the performance of these light guides the light transfer coefficients were measured and compared with simulated ones.

The measurement of the light transfer was performed by two independent methods: either one used the light generated by α -particles in a sandwich scintillator (fig. 2a) or relativistic electrons were applied in a triple coincidence arrangement (fig. 2b). These tests were done with XP4222B PMT's.

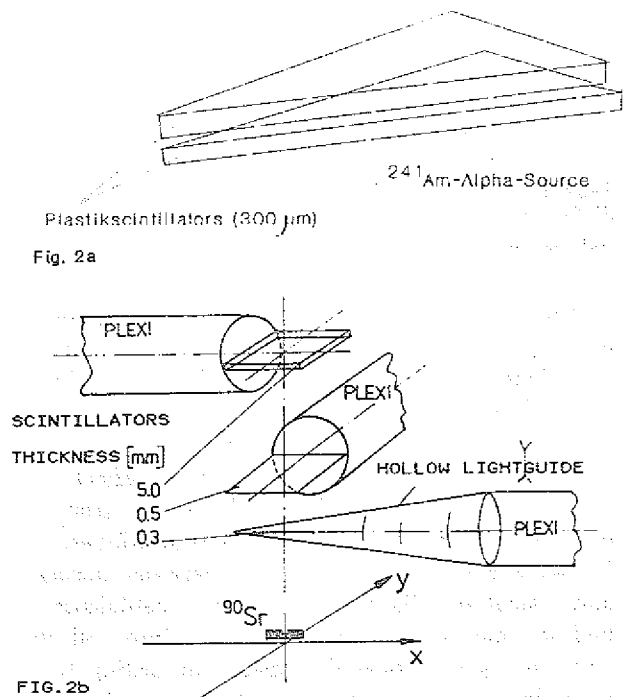


Figure 2: Arrangement for measuring light transmission in hollow light guides.

a) active sandwich scintillator with α -source

b) triple coincidence arrangement with β -source

The results of these measurements for different reflector materials are summarized in Table 1.

Table 1: Measured relative light transmission coefficients for different reflector materials.

Al/mylar		Ag		Al(foil)		Ni		white paper		black paper	
α	β	α	β	α	β	α	β	α	β	α	β
100		84	95	80	78	26		18		4	

Besides the investigation of different reflector materials the light transmission coefficients were measured also as functions of the length of the light guide tubes. The results were compared with simulation calculations (fig. 3).

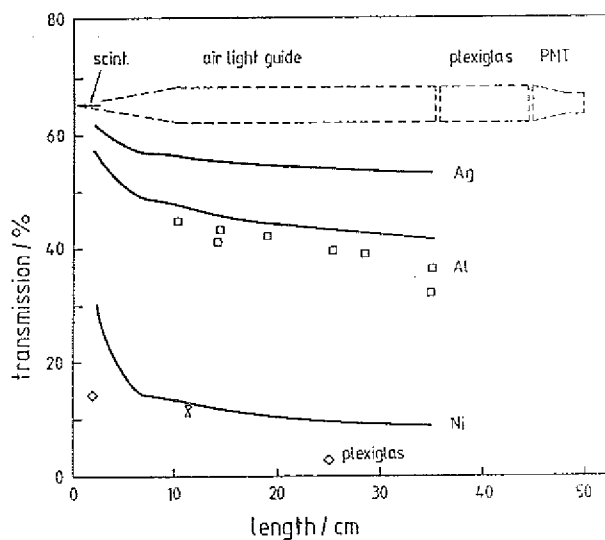


Figure 3: Light transmission for air lightguides with different reflecting materials (solid lines: calculated values, open symbols: measured data).

In the simulation for each event the origin of the light pulse in the scintillator as well as the directions of propagation are chosen randomly. After decomposing each ray into its two polarization components the intensity loss of the light rays by deflection (scintillator) and reflection is traced through the hollow light guide from the origin of the light pulse to the photomultiplier.

In fig. 3 a comparison is given between measured and calculated light transmission coefficients for hollow light guides with an aluminium reflecting layer and geometrical dimensions according to the start detector arrangement (indicated in the top of the figure). The comparison shows a good agreement between measured and calculated results.

A comparison was made of the light transmission efficiency of the hollow light guides with light guides of conventional plexi type. For this purpose the same scintillator as used in the hollow light guides was optically coupled by silicon grease to a plexi light guide of 1 mm thickness and 250 mm length (same as for the hollow light guide) as shown in fig. 4.

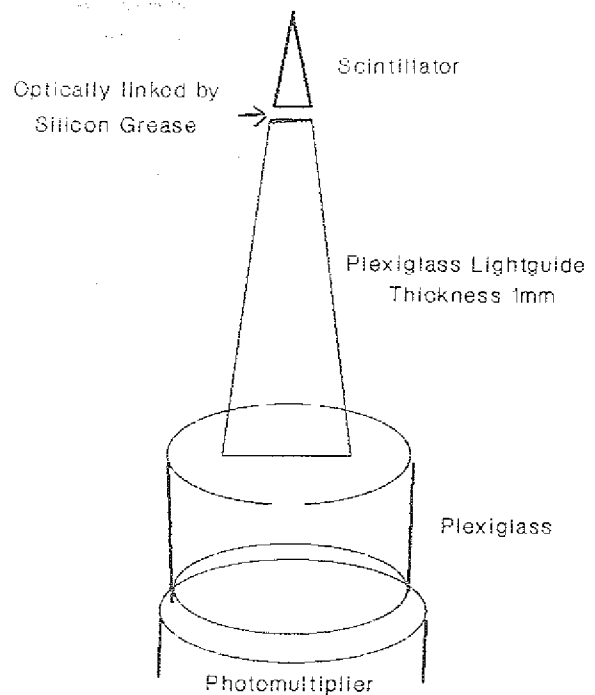


Figure 4: Plexi light guides for light read out of a scintillator element.

Assuming a transmission coefficient of 1 for the plexi light guide the theoretical light transition from the scintillator should amount to 19 % approximately. Due to mechanical problems with light coupling the measurements gave only 4 % transmission whereas for the hollow light guides an average light transmission of 40 % was measured. By means of measuring the one electron peak the average number of photoelectrons generated in the photomultiplier by one relativistic electron passing through a scintillator of 500 μm thickness was determined to be 60.

- a Zentralinstitut für Kernforschung Rossendorf, Dresden
- b Technische Universität Dresden

Reference:

1. Annual Report 1990, KFA, p.206–211.

A Liquid Hydrogen Target with very Thin Windows

V. Jaeckle, K. Kilian, R. Klein, H. Machner, W. Oelert, P. Turek

for the TOF collaboration

For several experiments at COSY we develop a small liquid hydrogen target with very thin windows based on the positive experience with similar targets built at CERN [1]. Small target volumes are of considerable interest, e.g.

- in measurements of excitation functions, because of energy loss and straggling;
- in all experiments to reduce systematic errors coming from secondary interactions;
- in experiments calling for a good definition of the interaction vertex;
- they reduce all safety problems.

In order to make the effects of the target window much smaller than the target effects a good liquid hydrogen target (~ 1 mm) should have Hostaphan windows with a thickness of $1 \mu\text{m}$ or less. An additional advantage of reduced window thickness is the possibility to effectively combine the liquid hydrogen target with a Cerenkov light readout. Thin walled containers have then at least two further important advantages:

- Unwanted Cerenkov light emission from the foils is reduced.
- The absorption of Cerenkov light crossing the foils on the way to the photomultipliers is drastically reduced. Especially blue and violet transmission is improved.

Lowest possible window thickness is achieved by stabilizing the pressure difference between the inside and outside of the target cell to a value just above the triple point of hydrogen. This has been achieved by using a very flexible cylindrical metal bellow as a hydrogen reservoir with a mass on top. This reservoir is connected to the target cell and is put inside the same vacuum as the target cell. We have chosen the mass to give a pressure difference of 0.2 atm between inside and outside of the windows independent of the operational mode (including cooling and warming up). The pressure difference of 0.2 atm allows the use of other elements like deuterium, helium, nitrogen and oxygen. Fig. 1 gives a schematic view. If we start in warm condition and pump down the isolation vacuum, then the hydrogen volume expands by a factor of six (going at the same time from an absolute pressure of 1.2 atm to 0.2 atm). It is just the volume of the bellow (it's elongation) which changes. Cooling down the target cell and

condensing hydrogen there, reduces the gaseous hydrogen content in the reservoir at constant pressure, and correspondingly the reservoir shrinks (Fig. 2). The target cell is made entirely of Hostaphan and is attached to a cooled copper rod where the hydrogen is liquified. No bubbles occur in the region where the beam should traverse the target. Their occurrence is suppressed by free convection of the liquid hydrogen in the cell. This is achieved by dividing the cell into two halves with a thin foil separating an upstream and downstream part.

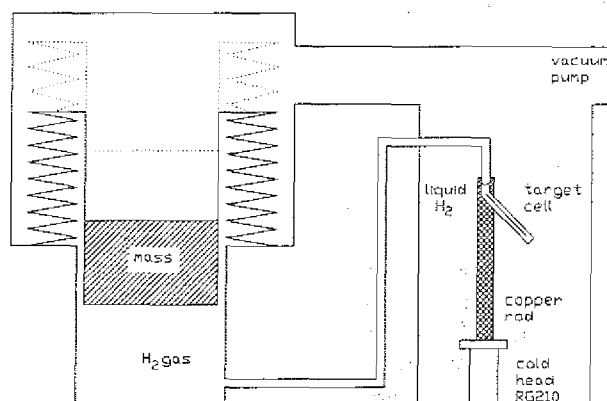


Figure 1: Schematic view of a liquid hydrogen target with low pressure difference.

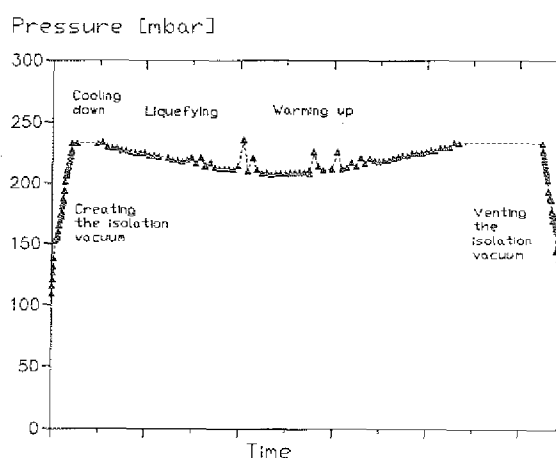


Figure 2: Measured pressure difference between inside and outside of the windows.

Reference:

1. K. Kilian, L. Mazzone and G. Novellini, A hydrogen target with very thin windows, CERN EP, Internal Report 85-02 (1985).

The Induction Drift Chamber (IDC) at High Rate

R. Bröders, M. Dahmen, G. Decker, K. Kilian, A. Kurtenbach, C. Lippert, W. Oelert, E. Roderburg,
K. Röhrich, T. Seifick, O. Steinkamp, R. Stratmann, S. Walsh, and M. Ziolkowski

The high rate capability of the Induction Drift Chamber was tested with 10 MeV protons of the Tandem accelerator at the Erlangen University. The mechanical parameters of the chamber are given in [1].

We tested whether the intense beam produces high voltage sparks or damages, and measured the rate dependence of the gas amplification.

The rate was measured with a Tektronix 2467 oscilloscope which has the option to measure the trigger rate. An indirect measure for the rate was the current between one cathode and ground. For zero rate this current was in order of 1 nA.

The beam had a symmetric shape with FWHM = 3–4 mm. The fluctuation of the intensity of the beam during one measurement was less than a factor of 2.

Results

1. The maximum rate measured was $2 \cdot 10^7$ protons per second and mm^2 . No damages and sparks have been observed. The electrical field is lowered by very high rates due to space charge effects and to voltage loss on the feed resistor. When the beam was switched off, the chamber recovered in less than a second.
2. Figure 1 shows signals of one wire for a rate of $2 \cdot 10^7$. This plot was taken with a 500 MHz digital scope (Tektronix 2440). The single events can be still separated.

CH1 20mV0 A 1x0 -10.6mV CH1

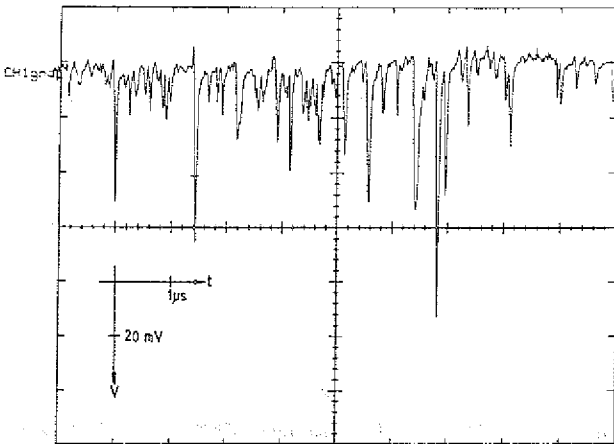


Figure 1: Signal of one wire for 10^7 protons per second and mm^2 .

In order to examine the dependence of the gas amplification on the rate, the signals of the anodes were recorded with FERA-ADC's (the gate was derived from the signals and had a width of 65 ns). The variation of the mean-value of the Landau-distribution with the rate for two different settings of the high voltage is shown in figure 2.

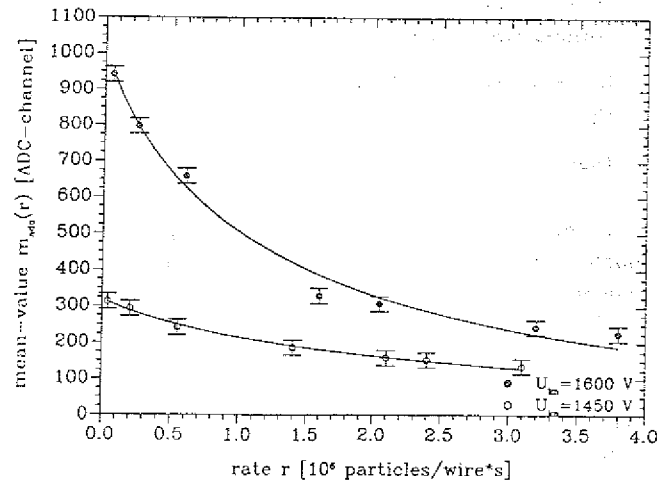


Figure 2: Mean-value as a function of the rate for two different voltage settings at the nearer cathode.

The calibration of the gas amplification is given in figure 3. The curves show the calculated gas amplification [2] normalized to the amplification with rate = 0.

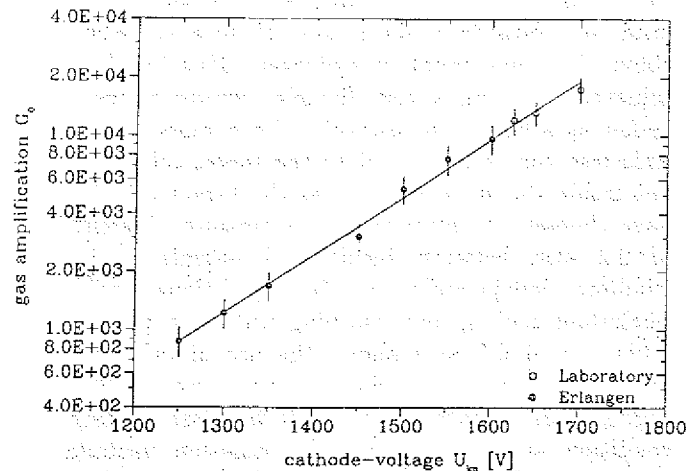


Figure 3: Gas amplification as a function of the voltage on the nearer cathode for rate 0.

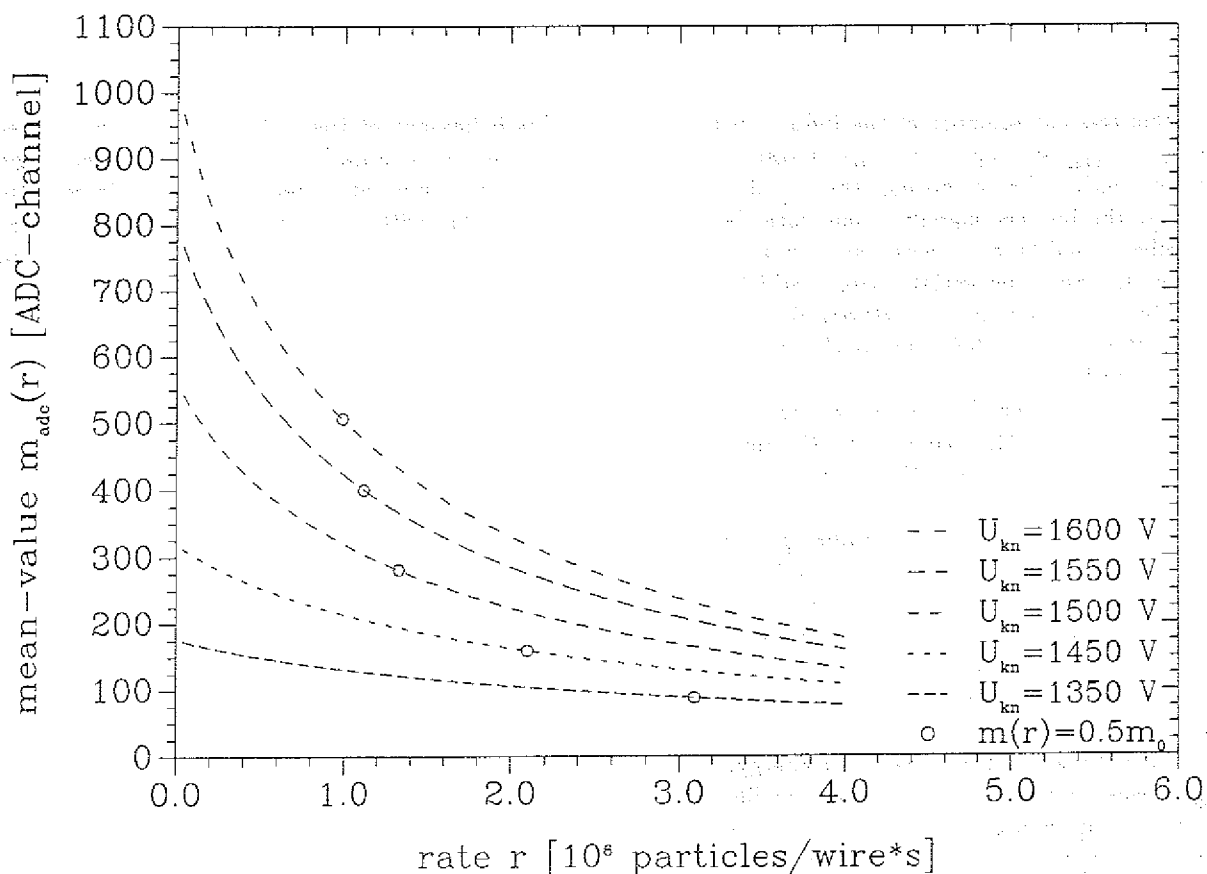


Figure 4: Calculated mean-values as a function of the rate for different voltage settings on the nearer cathode.

Measurements with minimum ionizing 2.5 GeV electrons at the synchrotron in Bonn have shown that the spatial resolution for a gas amplification of $1.5 \cdot 10^4$ is $30 \mu\text{m}$ and for an amplification of $8 \cdot 10^3$ is $40 \mu\text{m}$ (0 degree track, evaluated with ADC's). Compared to minimum ionizing particles the 10 MeV protons yield a factor 20 more ionization. Therefore, the rate dependence of the spatial resolution for minimum ionizing particles by a gas amplification of $1-2 \cdot 10^4$ is given by the measurement with 10 MeV protons at the lowest gas amplification (same ADC entry) (fig. 4). In fig. 4

the rate, where the gas amplification is reduced to a factor 2, is marked.

The chamber can stand a rate of more than $3 \cdot 10^6$ protons per second and mm^2 , without serious deterioration of spatial resolution.

References:

1. R. Bröders et al, Construction of a Vertex Detector consisting of independent Induction Drift Chambers, this Annual Report.
2. R. Bröders, Jül-Bericht, 1992.

The Use of Induction Drift Chamber Signals as a Fast Trigger

R. Bröders, M. Dahmen, G. Decker, K. Kilian, A. Kurtenbach, C. Lippert, W. Oelert, E. Roderburg,
K. Röhrich, T. Seifick, O. Steinkamp, R. Stratmann, S. Walsh, and M. Ziolkowski

Two possible readout schemes of the induction drift chamber (IDC [1]) are reported to get a fast trigger information: either by recording the anode wire signals or by the induced signals of one cathode.

The chamber consists of a wire plane with alternating anode and potential wires which are separated by 0.3 mm and two cathode foils with a distance of 0.5 mm and 3.0 mm to the wire plane (see [2] for details).

The cathode foils which are produced in the "Targetlabor" of the IKP consist of 25 μm thick capton foils and are coated with a layer of 3000 Å of gold.

Isobutane was used as the drift chamber gas. It has an electron drift velocity in the order of 50 $\mu\text{m}/\text{ns}$ at a typical electric field of about 8 kV/cm in the drift region. Isobutane is a dense gas with a high primary ionization density (40–50 ion pairs/cm for minimum ionizing particles).

The presented TDC drift time spectra were recorded [3] with 2.28 MeV maximum energy. The stop signal was obtained from a plastic scintillator. The drift time spectra of 4 adjacent anodes of one plane are shown in figure 1. All TDC spectra have a sharp peak of 6–9 ns FWHM and a tail of longer drift times.

The events with longer drift time correspond to inclined tracks, which have in general short response times at another wire in the plane.

This behaviour of the TDC spectra is confirmed by a computer simulation of the drift cells. Figure 2 shows the simulated spectra of 3 adjacent anodes for tracks with an angle of 15 degrees to the normal. All tracks start in cell 1 and the primary ionization clusters are distributed with a Poisson statistics along the track and the shape of the measured spectra is well described.

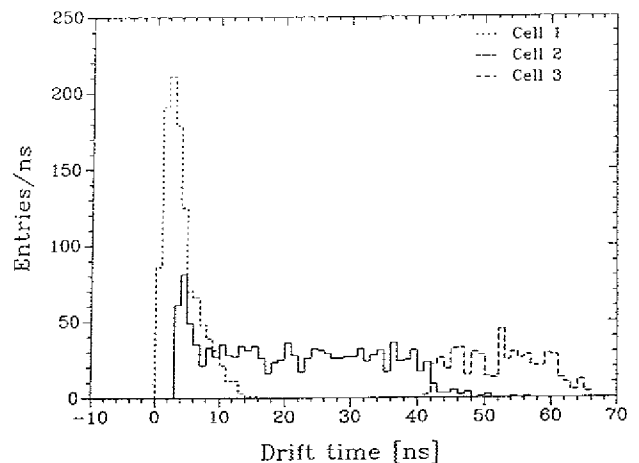


Figure 2: Computer simulation of the drift time in 3 adjacent cells.

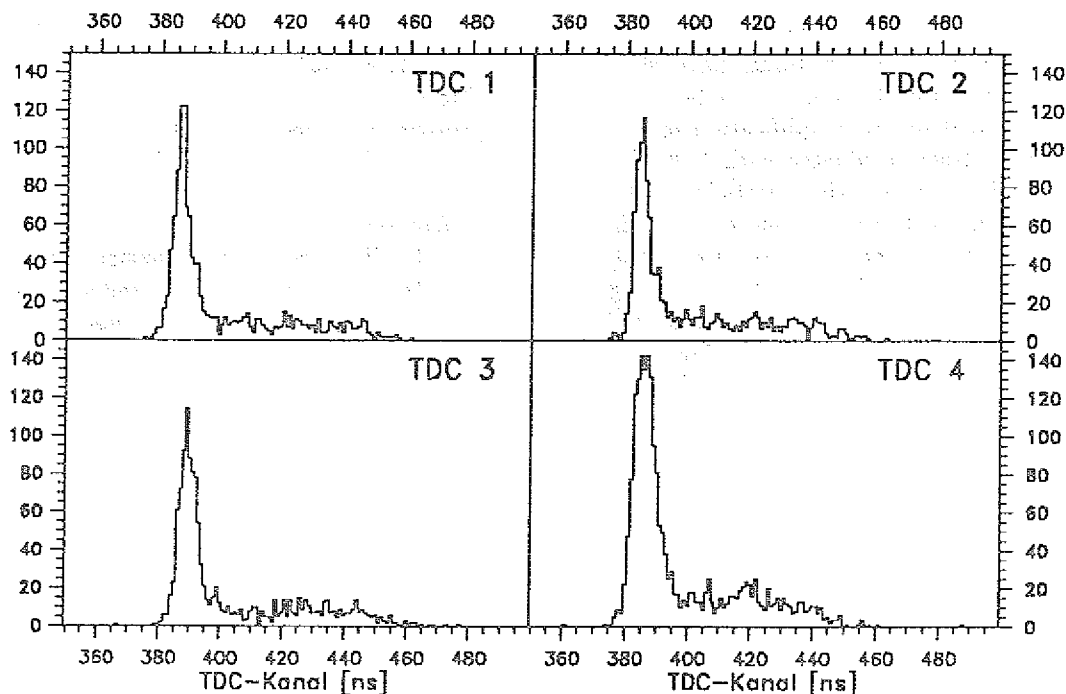


Figure 1: TDC spectra of 4 adjacent anodes.

Another possible readout of the fast trigger signal can be realized by recording the induced signals on the near cathode (distance 0.5 mm to the wire plane). The time behaviour is of course the same as for the anode signals. The use of the cathode signals allows different subdivisions of the foils to be made.

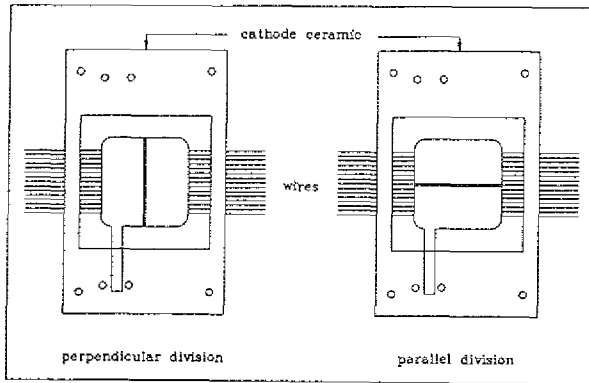


Figure 3: Realized subdivisions of one cathode.

For a parallel division of the cathode (fig. 3) the performance of the track separation between the two parts is examined. Therefore the normalized difference of the induced signals is plotted versus the track position, which is measured from the wire signals (fig. 4). A clear separation of the two parts for normal incident tracks can be seen.

The combination of very small drift cells together with a high ionization density and a high drift velocity allows the use of the IDC signals for a fast trigger (minimal coincidence length about 10 ns). This increases, for example, the efficiency of a vertex detector using the chamber signals in a first level trigger.

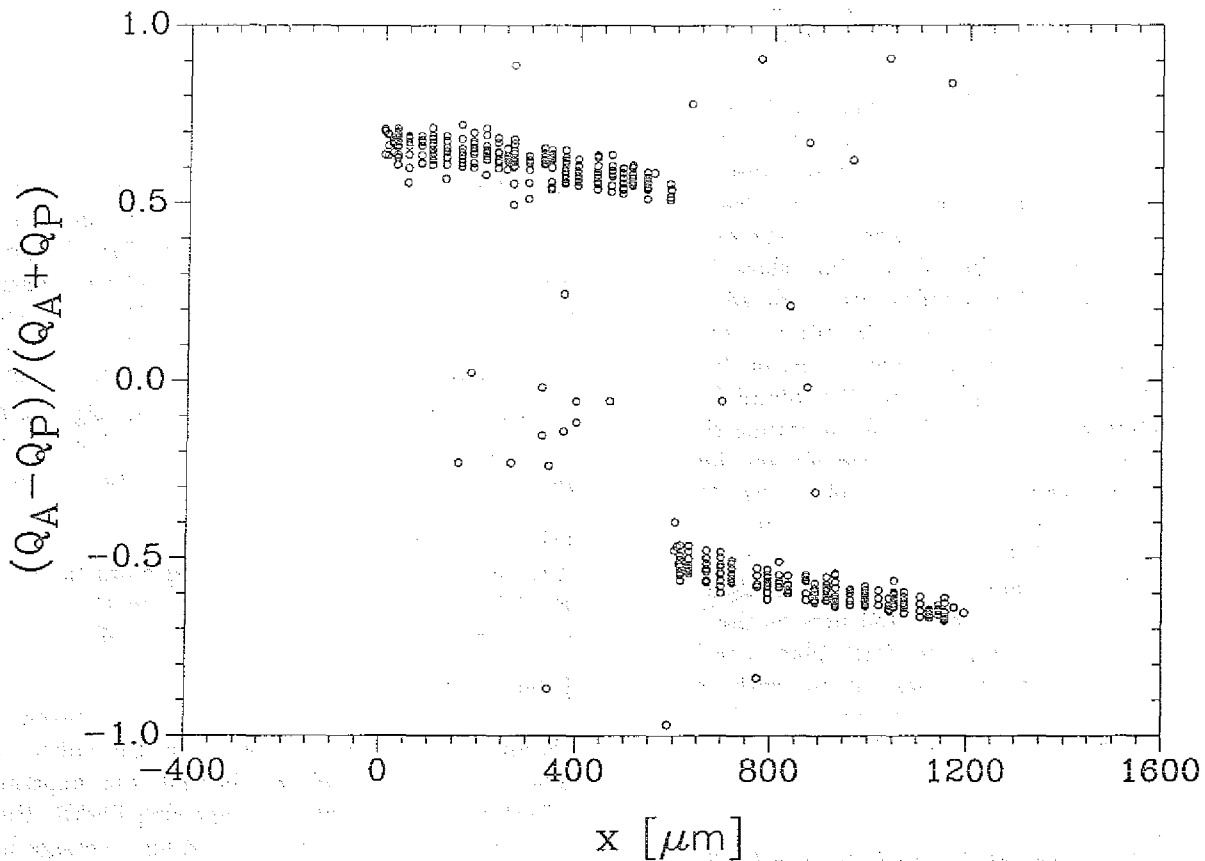


Figure 4: Scatter plot of the normalized difference of induced signals on two halves versus x-position of track for normal incidence.

References:

1. A.H. Walenta et al., Nucl. Instr. and Meth. A265 (1988) 69.
2. R. Bröders et al., Construction of a Vertex Detector consisting of independent Induction Drift Chambers, KFA Annual Report 1991.
3. M. Dahmen, Diploma-thesis, Jül-Bericht, 1992.

Construction of a Vertex Detector Consisting of Independent Induction Drift Chambers (IDC)

R. Bröders, M. Dahmen, G. Decker, K. Kilian, A. Kurtenbach, C. Lippert, W. Oelert, E. Roderburg, K. Röhrich, T. Seifick, O. Steinkamp, R. Stratmann, S. Walsh, and M. Ziolkowski

Strange particles which decay due to the weak interaction have a mean life time in the order of 10^{-10} seconds. This corresponds for momenta which occur in the COSY region to a decay length in the centimeter range. To evaluate the kinematics of strange particles in a reaction it is very helpful to reconstruct the decay vertex. Therefore it is necessary to measure tracks near the target region with high spatial resolution.

As a detector we chose the IDC [1], as this chamber is able to measure with high precision in a high rate background, which is found near the target.

Mechanics

We built a detector which is a modular stack of drift chambers which are only by 3.5 mm apart from another. Each chamber consists of 36 drift cells, which are defined by one anode and two potential wires placed between a near (.5 mm) and far (3 mm) cathode (Fig. 1).

The active region of one plane is about 20-20 mm². In order to form a homogeneous electrical field at each side 10 anodes and 11 potential wires are added. The wires are glued on ceramic frames which have burned in grooves to provide the exact position of the wires. On alternating sides the potential wires and the anodes are soldered on printed circuit boards. The circuit boards are made by flexible capton foils with copper strips on both sides [2]. The foils are glued in the aluminium frame providing a smooth surface for a sealing ring of the next chamber. All electrical signals and high voltages of the chamber are supplied by this printed circuit. To each wire plane a ceramic frame with two cathode foils is attached. One foil is directing to the wires and has a distance of .5 mm to the wire plane. The opposite foil acts as the far cathode for the next wire plane. Each plane can be gas tight connected to the next plane with one sealing ring. The orientation of the wire direction is changing by 120° degrees. The full detector layout is shown in Fig. 2.

High Voltage

As the chamber is operated with isobutane (+ 5 % ethanol) very high voltages are needed to get a gas amplification of more than 10^4 . Each chamber is supplied with one voltage source, the voltages are divided with a resistor chain and the currents are measured with digital voltmeters with floating power supplies (Fig. 3). The high voltage supplies are built modular so it is possible to adjust the gas amplification for each chamber independently.

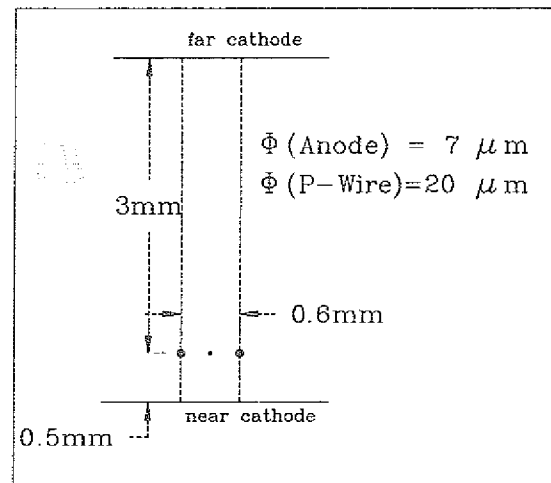


Figure 1: Mechanical parameters of the IDC.

Electronics

400 Channels are provided for the read out of the anodes and potential wires. The preamplifiers are mounted directly at the end of the flexible capton printed circuits. The amplifiers are charge sensitive with a rise time of about 5 nsec. The differential outputs are fed out from the experiment area by 15 m twisted pair cables. To avoid pile up effects from the long fall time of the chamber signals, they are symmetrized by pole zero shaper. The preamplifiers and shaper modules are developed and fabricated at the University of Siegen with thickfilm technology.

The performance and the programming of the FADC-system is described in [3] and the computer system for data taking is described in [4].

Measurement

A first measurement with a stack of seven wire chambers was done at CERN in December 1991. The detector was placed behind the experiment PS185 at the antiproton storage ring LEAR. During the run three wire chambers had high voltage break downs and could't be operated. With the remaining four chambers more than 1 million antiproton tracks were recorded.

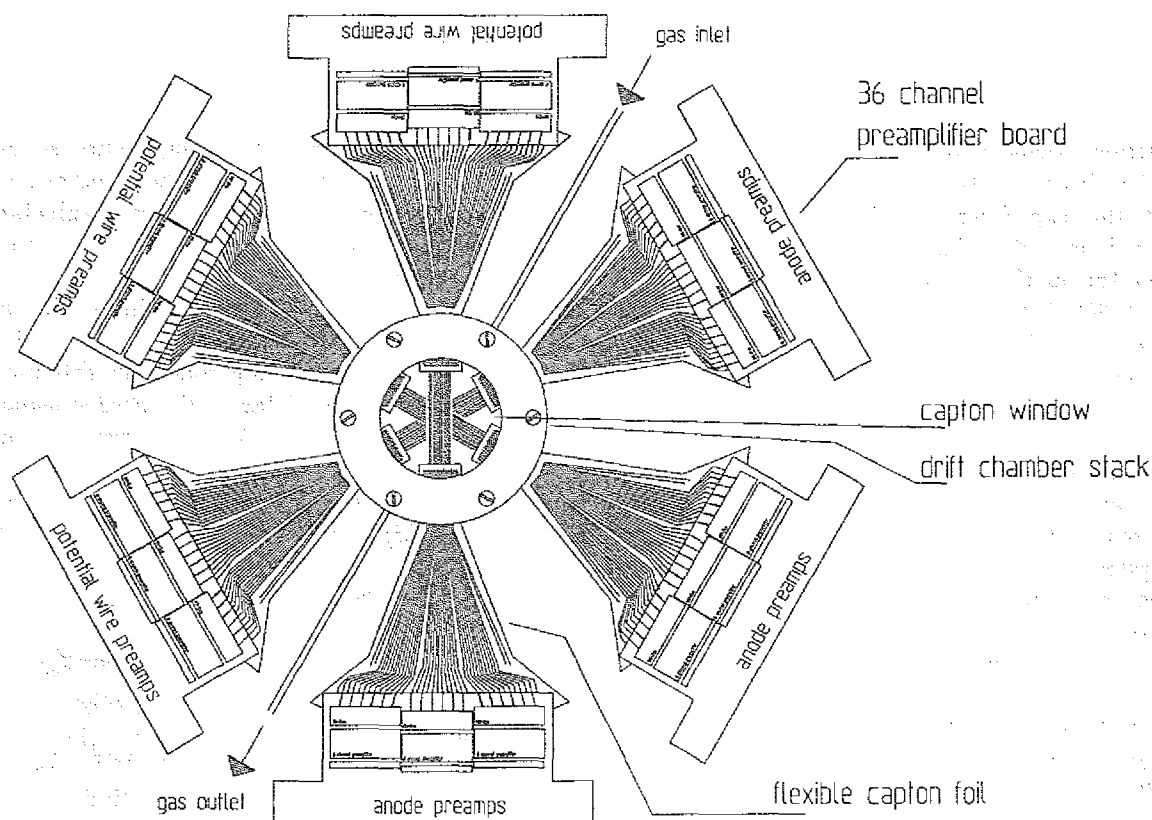


Figure 2: Schematic view of a stack with three drift chambers.

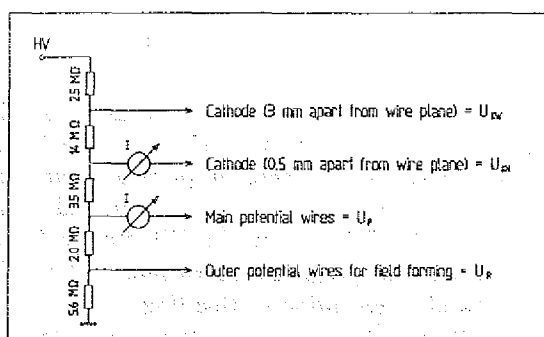


Figure 3: High Voltage Distribution

References

1. A.H. Walenta et al., Study of the IDC as a High Rate Vertex Detector for the ZEUS Experiment, NIM A265 (1988) 69;
E. Roderburg et al., The Induction Drift Chamber, NIM A252 (1986) 285.
2. The printed circuit capton foils were fabricated in the workshop of Mr. Gandi, CERN.
3. C. Lippert and G. Decker, this annual report.
4. G. Decker and C. Lippert, this annual report.

C. Lippert, R. Bröders, M. Dahmen, G. Decker, A. Kurtenbach, W. Oelert, E. Roderburg,
K. Röhrich, Th. Seifick, R. Stratmann, S. Walsh, M. Ziolkowski

An Induction Drift Chamber (IDC) with 400 channels has been constructed to measure charged sigmas at the experiment PS185 at LEAR. The read out is done with a Flash-ADC system DL350¹ to improve the spatial resolution of inclined tracks. We use DL362 FADC modules² with a digitalisation rate of 100 MHz. The sampling length is 256 words per channel (2.56 usec), and the resolution is 8 bit with a non linear response which corresponds to a dynamic range of 10 bit [1]. Besides the FADC modules each DL350 system houses two or more controller and interface modules (SIM, SAM), which provide the micro programming of the FADC-system.

The SAM (sampler module DL352) is the central timing controller, which asserts the clock, the address and the control lines for the FADC modules.

The two tasks of the SIM (Scanner and Interface module DL357) are: interface to a Front-end-processor (FEP) with a VSB-bus port and autonomous scanner and hit detector with zero suppression.

As FEP we used a VME-processor FIC 8232 under OS9. The DL350-system is handled by the FEP as a local I/O system.

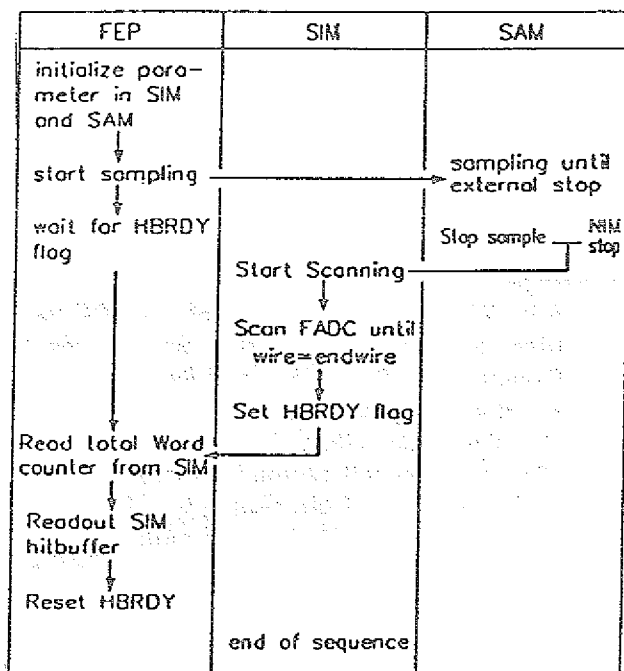


Figure 1: Flow Chart of FEP, SIM, SAM

¹Dr. B. Struck, Bäckerbald, D-2000 Tangstedt

²The DL362 modules were rent from the Obelix group at LEAR

To handle these hardware components an online-program is written in "C". The program can be used interactively. From the inputs it calculates the base addresses of the crate and the modules for later access to them.

The flow chart (see. Fig. 1) shows the general interaction between the FEP, SIM and SAM operating in common stop mode. In this mode the system provides the delay of the analog signal with respect to the trigger. After programmable initialization of SIM and SAM the FEP starts the continuous sampling process by software. The external trigger signal of the experiment stops the SAM and is also used to start the scan process of the SIM (see Fig. 2).

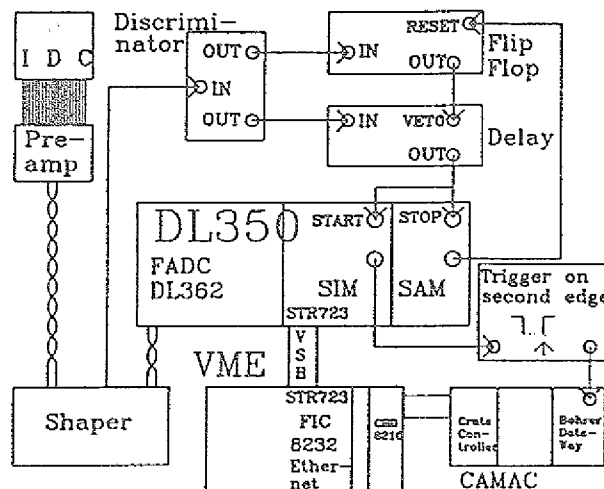


Figure 2: Trigger plan of experiment

The SIM scans through all specified FADC memories and stores only the relevant data in the local hit buffer. For that zero suppression a variable threshold can be set by software. At the end of the scan process the SIM raises the Hit Buffer Ready Flag. The FEP is programmed to read out the hit buffer data after reception of this flag.

Using the Scan Offset and Scan Range registers the scanner operation can be optimized to the actual drift time and the trigger delay conditions (see Fig. 3). Because 16 bits are scanned in parallel both the anode and the potential wire must be read out even if data are seen only in one wire due to zero suppression. Therefore a special routine has been written to distinguish between anode and potential wire.

The Scan Offset, the Scan range, the threshold for the zero suppression, the thresholds and offsets for the FADCs, the Sampling rate and mode, etc. are

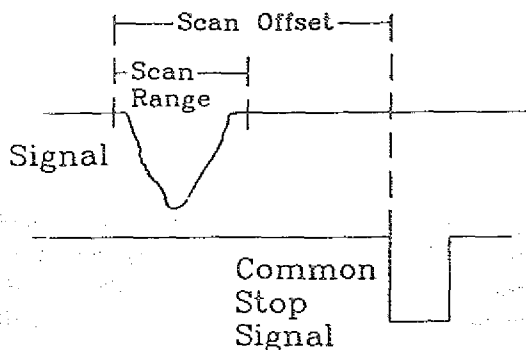


Figure 3: Scan Offset, Scan Range

accessible by read/write registers which is managed by special bit operations from the program. An automatic run is programmed to adjust the baseline for all channels. A semi graphic is implemented to

check signal height, overflow, timing, etc. Data acquisition to the VME—Hard disc is possible, even with the programmed DMA—channel, but with a slow rate.

It should be mentioned that the event size is very high (≥ 1 kByte) and by the high rate of 200 events/sec an efficient read out system and a fast tape drive is necessary. For that read out system (SPIDER, see [2,3], all registers have to be loaded by the online program and in addition, a file with all relevant hardware addresses is generated, which is an include file for the SPIDER—program.

References:

1. STRUCK: Technical Manual, DL350, Flash ADC Crate, DL257 SIM
2. G. Decker, this annual report: A real time data acquisition system
3. Y. Perrin, SPIDER, User's Guide, CERN

A Data Acquisition System for an Induction Drift Chamber

R. Bröders, M. Dahmen, G. Decker, A. Kurtenbach, Ch. Lippert, W. Oelert, E. Roderburg, K. Röhrich, Th. Sefzick, R. Stratmann, S. Walsh, M. Ziolkowski

For measuring charged Σ 's with the Induction Drift Chamber (IDC) at the experiment PS185 at LEAR a data acquisition system has been built.

The needed hardware (a CPU in a VME environment running under the real time operating system OS9, an interrupt source, tape drive and interfaces to datataking devices) was installed.

From the CERN data handling division a software package called SPIDER [1] is available. This software consists of a circular memory buffer manager and handels several frames which have to be filled and adapted to the users requirements by the user himself. Each frame creates its own process. In addition there are interprocess communication and debugging tools available.

For the use in the PS185 experiment we programed in the "producer" frame the Flash—ADC readout [2], the CAMAC readout and an event by event synchronisation to the already existing data acquisition system of PS185. The taken data are written to the circular buffer. This process is started by an interrupt signal whenever required by the connected devices. This happens without respect to other processes on the CPU.

The "recorder" frame reads out the data from the buffer and writes them to a STK4280 SUMMIT drive which is able to store up to 1 Mbyte per second.

In the "analyser" frame we implemented a transfer of a data sample out of the circular buffer via cheapernet (under TCP/IP). The time intensive analysis of data is done by an other computer. In addition some graphics for fast checks of data is built in.

Finally the process of the fourth frame controlles the other three processes and is simultaneously the user interface.

Time measurements for the whole setup show an event overhead for SPIDER lower than 1 ms per event. The total used time with read out of devices and writing to tape is 2 ms per event with an eventlength of 2 kbytes.

The average event rate of the taken data at LEAR was about 200 Hz which is below the possible limit.

References:

1. Y. Perrin, SPIDER, User's Guide, CERN.
2. Ch. Lippert, ONLINE Program for data acquisition with a Flash—ADC system, this Annual Report.

The construction and performance of a three layer circular hodoscope has been described in recent annual reports of the IKP-KFA. Employing this hodoscope it is possible to determine pixels defining the spatial area of a charged ejectile of a reaction. This fast pixel determination is very useful, but it is not straight forward because of the existence of "ghosts" — i.e. the number of pixels one gets by counting the mathematically and geometrically possible combinations of "fired" segments from the three layers is often larger than the real number of pixels (fig. 1a) if there are two or more coincident hits. This is due to the fact that counting the pixels under combinatorial view can not prevent from combining segments which are hit by different particles. One can get rid of "ghosts" by increasing the granularity. Another way to get rid of these ghosts is to look at the timing of the scintillator signals. Due to symmetry there has to be the same light collection length and collection time for a true hit in both left and right twisted szintillator components. If the time difference between two combined twisted segments is larger than the equivalent time difference due to the pixel size these segments must have been hit by different particles. This then allows to omit this combination as ghost and to reduce the number of possible pixels (fig. 1b). Only a "well timed" coincidence of the three segments (one pie-shaped, two twisted) gives confidence that a true hit has been found.

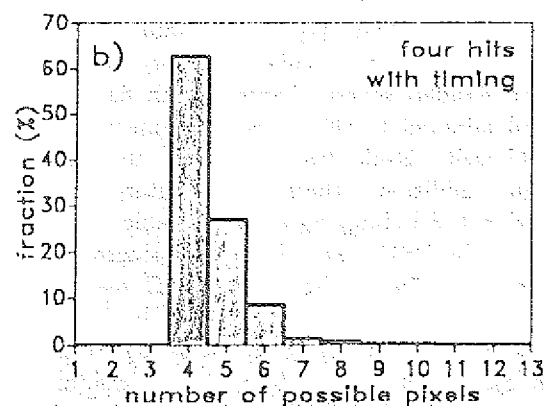
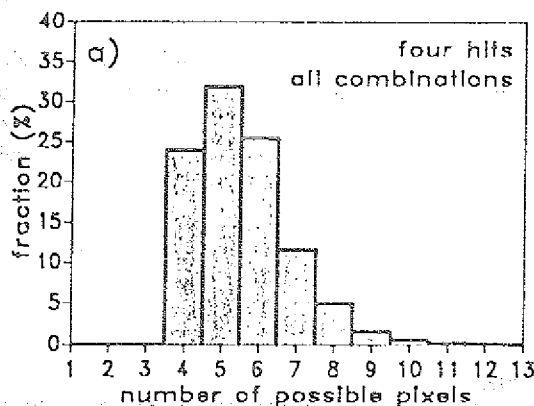


Figure 1: Monte-Carlo calculations: number of possible pixels for four incident particles, without (a) and with a coincidence (1 nsec. gate) between the two twisted szintillators (b).

During our works on the superfast coincidence system with the time resolution less then 1 ns, we found it necessary to develop a delay unit for delaying the leading edge of pulses. Such a unit has been designed as a "hybrid" in the SMD technology. It allows for a continuous regulation of the delay by changing the control current. The block diagram is shown in Fig. 2.

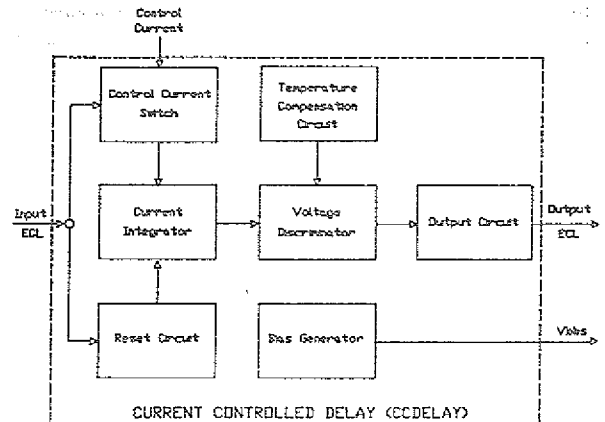


Figure 2: Current controlled delay block diagram.

A prototype unit has been tested using the Fastbus TDC converter LeCroy 1875. The results are shown in figs. 3, 4 and 5. The delay-current relation is shown in fig. 3a, 3b. The time-difference-spectra

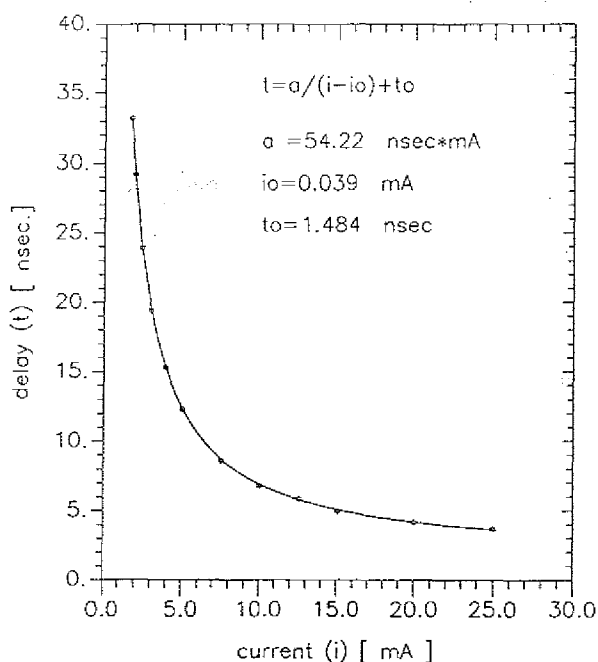


Figure 3a: Propagation time in the delay circuit via current control in the range of small delays.

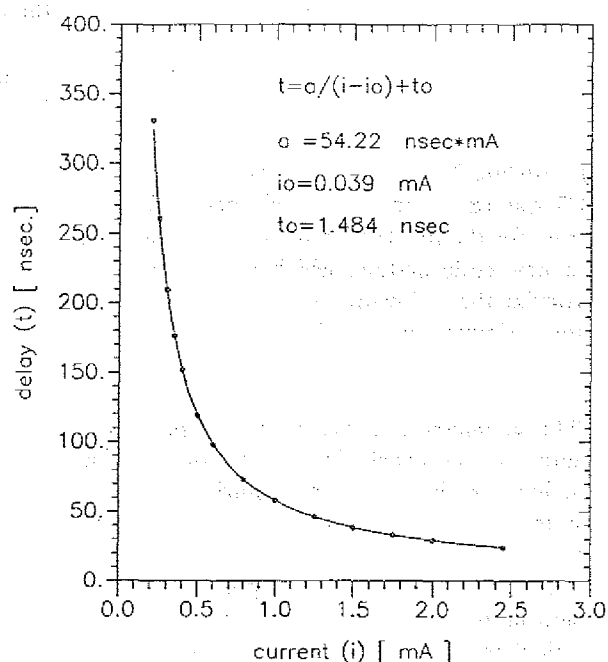


Figure 3b: Propagation time in the delay circuit via current control in the range of large delays.

containing the coincidence peaks obtained with various delays within the range of about 200 ns are shown in fig. 4. The delays were obtained using 50 Ω coaxial cable as well as with the built delay unit CCDELAY (see fig. 4). No relevant differences in the delay jitter for these two cases has been observed. The time resolution (FWHM) of the measuring system was 50 ps. The self-coincidence peak obtained at the delay of 97 ns with the CCDELAY is shown in fig. 5. Its width confirms the good long-time stability of the delay obtained with the CCDELAY unit. After the successful test of the prototype a coincidence unit for a whole scintillator hodoscope will be developed.

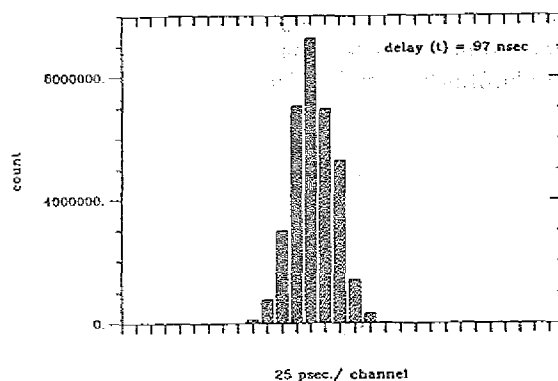


Figure 5: TDC spectrum using the developed delay circuit with a delay of 97 nsec, measuring time 12h.

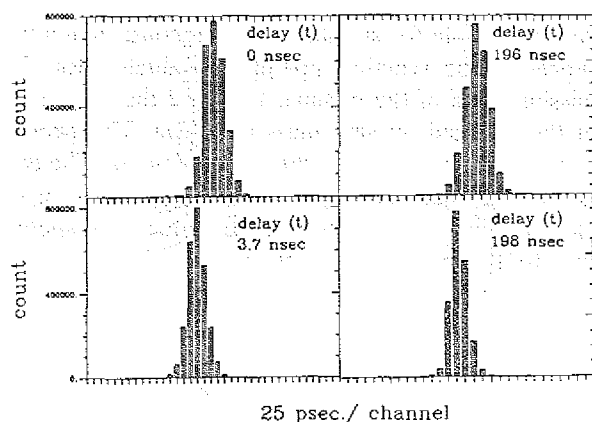


Figure 4: TDC spectra using 50 Ω cable delay (upper part) using the developed delay circuit (lower part).

^a Jagellonian University, Krakau/POLEN

Monte Carlo Simulation of the COSY Time-of-Flight Detector (TOF)

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for the COSY TOF Collaboration

The existing Monte Carlo detector simulation code MC3 was extended by several special features, necessary for the TOF detector simulation. This led to a new code system, which was named MC4. In particular the following extensions to end up with a realistic simulation of TOF response were developed [1].

- The geometry package is now capable of taken into account a realistic structure of the scintillator stripes for the barrel- and quirl-type parts of the detector. It calculates the sectioning of particle tracks by the geometrical structure of the detector. The tracks may start at any point of the geometry and in any direction. This feature is necessary to calculate particle tracking starting at e.g. decay vertices or from nuclear interaction points in the detector material.
- An algorithm for dE/dx -energy loss determination was developed for the special TOF requirements. The code uses for long tracks an inverse table lookup in an energy range table generated by the Bethe-Bloch-formula. For short tracks a numerical integration scheme on dE/dx is used.
- The scintillation light is calculated taking into account the quenching effect by Birk's law.

$$dL/dx = \frac{S}{1 + k_B \cdot dE/dx}$$

L = scintillation light

S = material dependent parameter

k_B = material dependent parameter describing the quenching effect.

A typical used value for k_B is 8.5×10^{-6} [cm/GeV].

The light transition times of a scintillator and the measured fluctuations thereof are also computed.

The scheme of performance of the MC4 TOF simulation code is shown in Fig. 1. The simulated detector response can be written on file for later and repeated evaluation. The data are written in ZEBRA format.

In the MC4 version it is now possible to simulate the response of the TOF detector on particle hits as scintillator hits and the related time and pulse height signals including realistic simulation of the fluctuations of these quantities.

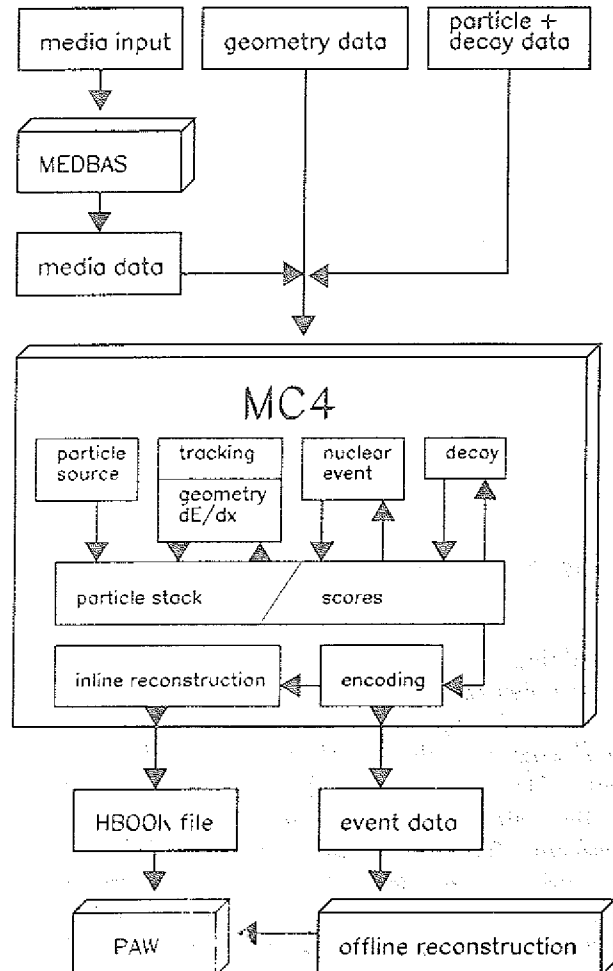


Fig. 1: Scheme of performance of the MC4 TOF simulation code

As an example for simulation and reconstruction we represent the reaction $pp \rightarrow pp\gamma$ evaluated for the missing mass of the gamma. In Fig. 2 the frequency of the squared missing mass is drawn. The prominent peak around zero is due to the gamma. The two additional peaks are the missing masses of one or two protons due to incomplete detector response due to the proton stopping in the first or second scintillator layer.

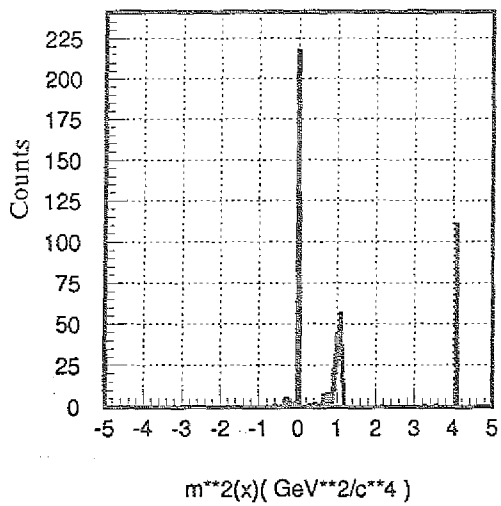


Fig. 2: Missing gamma mass of the $pp\gamma$ reaction

To demonstrate the simulation of fluctuations Fig. 3a-c may serve. They show scatter plots of the TDC signal and the corresponding precise time-of-flight of the protons from the $pp\gamma$ reaction. The figures are generated for different scintillator layers and different detector modules. The detector resolution is obviously better in the quirl type module than in the barrels. Also there are losses of hits in the outer layer of barrel 3 as compared to the inner layer. This is due to stopping of protons in the first and second layer.

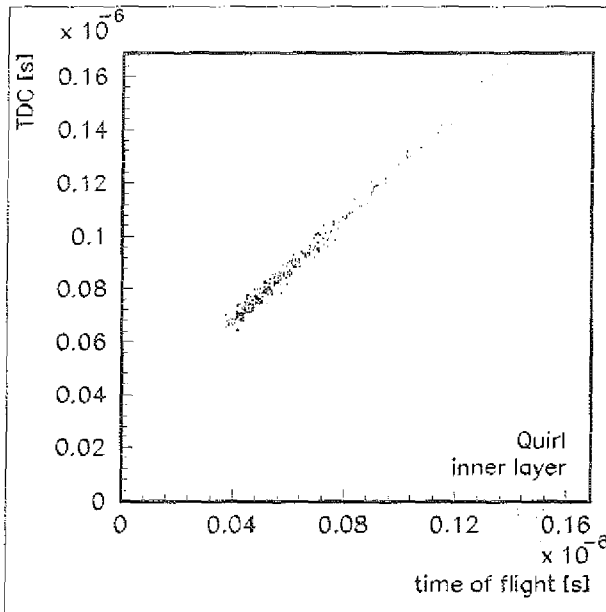


Fig. 3a: Quirl inner layer

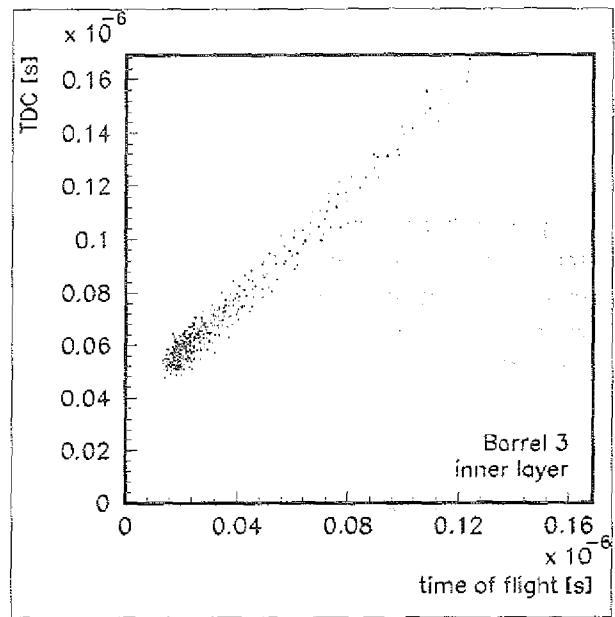


Fig. 3b: Barrel 3 inner layer

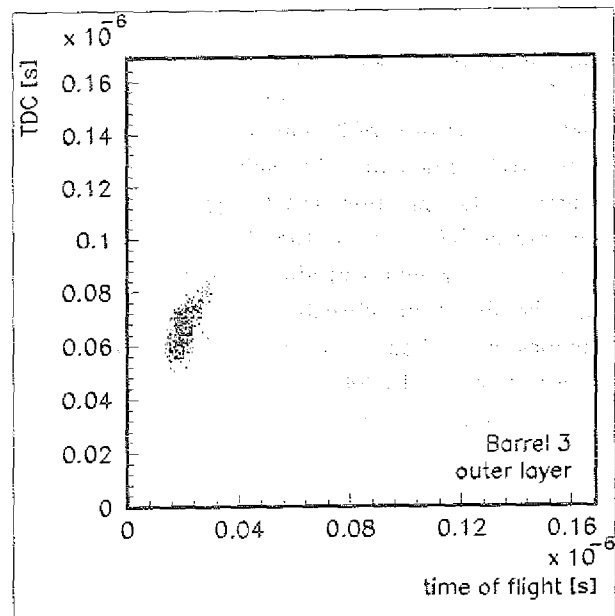


Fig. 3c: Barrel 3 outer layer

Reference:

/1/ Annual Report IKP 1990, Jül-2462

Feasibility Studies for a TOF- γ -Detector at COSY

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If a proton beam hits a LH-target in the TOF detector, in case of a reaction leading to $pp\gamma$ the protons can be measured inside the detector. But the gammas, which normally fly in the backward half-space with respect to the proton beam, can only be calculated as the missing mass. So it was suggested to use the Bochum- γ -detector to measure those gammas. This detector consists of 54 modules, each one having the shape of a truncated triangular pyramid and being made of polycrystalline pieces of NaI (Tl). The 54 modules cover an angular regime of about one steradian. Detailed information about the detector can be found in ref. 1.

We performed a computer simulation to test the accuracy of the γ -detector with respect to direction and energy of the gammas. In our calculations we put the detector around the proton beam axis with a centered beam tube of 2.5 cm radius (see fig. 1). The code MC4 calculated by solving the kinematic equations the direction and energy of gammas for the case of 300 MeV protons hitting an LH-target. 2.5 % of the gammas produced by code MC4 went into the direction towards the γ -detector. Those gammas were handed over to the Monte Carlo transport code EGS4 [2], which calculated the electromagnetic shower in the NaI(Tl)-detector originating from the incoming gamma. The result of the EGS4 calculation were the deposited energies in the modules of the detector for each incoming gamma.

For comparison, which amount of the energy of the incoming gammas is found in the γ -detector, the deposited energies of all modules were added up and compared to the starting γ -energy. The following results were found for 886 events:

energy loss in EGS4 (%)	frequency of events (%)
0 - 5	59.3
5 - 10	18.2
10 - 20	12.8
20 - 99	9.7

So it can be stated that $\sim 90\%$ of the events show energy losses less than 20%.

To reconstruct the flight direction of the gamma the deposited energy of each module was multiplied by the direction vector of the midpoint of this module and then all those products were added up. The resulting direction was then compared to the flight direction of the incoming gamma. We got the following results for 886 events:

angular difference between reconstruction and incoming gamma (degrees)	frequency of events (%)
0 - 2	33.0
2 - 5	61.9
5 - 12	5.1

So it can be stated, that $\sim 90\%$ of the events show angular differences less than 5 degrees.

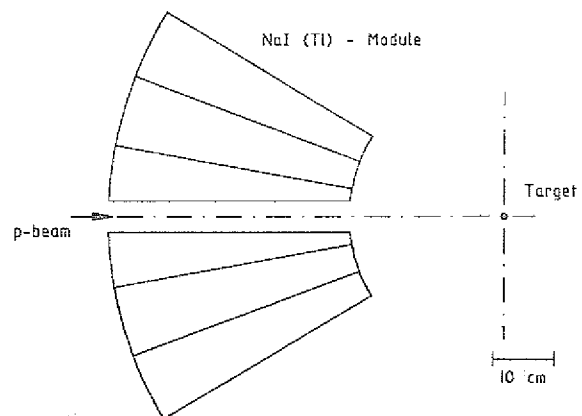


Fig. 1: Cross section of detector geometry for EGS4 calculation

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Concept of the 0^0 Facility at TP2 in COSY

W. Borgs, M. Büscher, D. Gotta, D. Grzonka, H.R. Koch, W. Oelert, H. Ohm, O.W.B. Schult, H. Seyfarth, K. Sistemich, K.H. Watzlawik, J. Ernst^a, F. Hinterberger^a, L. Jarczyk^b, B. Kamys^b, St. Kistryn^b, A.Kozela^b, J. Smyrski^b, A. Strzalkowski^b, S. Dienel^c, K.W. Leege^c, H. Müller^c, Chr. Schneidereit^c, S.V. Dshemuchadze^d, V.I. Komarov^d, V. Lyachenko^d, F. Nichitiu^d, V. Travkin^d, V. Zalyhanov^d, N.I. Zhuravlev^d, V. Koptev^e, V. Abaev^e, P. Birien^f, W. Cassing^g, H. Dombrowski^h, R. Santo^h

The 0^0 Facility will be a multi-purpose device for the investigation of ejectiles from proton-induced reactions at internal targets in COSY. It is planned to install a magnet system at the target place TP2 in the straight section of the ring which imposes a bump on the closed orbit of the proton beam, see Fig. 1a. The magnets D1 and D3 deflect the beam out of its original direction and back into it. The central magnet D2 will serve simultaneously for the separation of ejectiles from the circulating beam and as a momentum spectrometer. Its properties mainly determine the experimental possibilities at the 0^0 Facility. In view of the envisaged studies and in accordance with the recommendations¹⁾ by the COSY Program Advisory Committee and with the results of discussions during a "brain storming" session in Febr. 91, D2 is designed as a C-shaped magnet, see Fig. 1a and b, with rectangular pole shoes of 140 cm length and a gap height of 20 cm. With such a magnet the 0^0 Facility has the following characteristic features:

- Possibility to study simultaneously positively and negatively charged ejectiles, the latter in the large window of the yoke of D2.
- Large solid angle (> 50 msr) which is important especially for studies of correlated ejectiles. The large gap height also enables an efficient discrimination against particles scattered from the pole pieces.
- Momentum focussing for ejectiles which leave D2 through the sides. This is shown in Fig. 1a for particles which are emitted under 0^0 , $\pm 4^0$ with respect to the beam and with momenta $p = p_0/n$, $p_0 =$ projectile momentum. The position of the focal surface and the momentum resolution depend on the distance between the target and D2 and can be optimized for each experiment. A momentum resolution as good as $\Delta p/p = 2 \cdot 10^{-3}$ can be achieved²⁾. If the scattering in vacuum windows or detectors is included the expected resolution²⁾ is about 1 – 2%.
- Possibility to measure angular distributions up to $\sim 20^0$ in the setup shown in Fig. 1a. Larger angles are accessible with D2 turned around the target. Then, D1 and D3 are not used and the beam passes through the window in the yoke.
- Space of about 120 cm between D1–D2 and D2–D3 for the installation of targets and detectors. This allows to optimize the target–D2 distance, to have close geometry³⁾ for gas–jet or cluster⁴⁾ targets with maximum overlap between jet and beam (high luminosity, $L \sim 3 \cdot 10^{30}/\text{cm}^2 \cdot \text{s}$) and to place detectors around the target as well as

between D2 and D3. Also polarized–gas targets of the Filtext type^{5,6)} can be used. These properties can only be achieved at TP2 since this is the place with optimum beam characteristics and the largest free space for experiments in the COSY ring.

- Possibility to study correlated ejectiles which are emitted under 0^0 and 180^0 as in the deuteron breakup. In addition to D2, D1 is then used as a spectrometer magnet.
- High flexibility in the selection of the momentum range of the ejectiles through a power supply for the magnets which is independent of that of the COSY ring. Although the 0^0 Facility magnets are ramped with the ring, their field strength can then be chosen separately.

The following studies are planned at the 0^0 Facility:

- Study of the Subthreshold K^+ Production with a 0^0 Facility at TP2 in COSY (Proposal 18, Ref. 7). Inclusive double differential K^+ cross sections^{7,8)} $d\sigma/d\Omega dp$ and correlated $K^+ + X$ ($X = p, d, \dots$) spectra will be measured at different projectile energies below the NN threshold at 1.58 GeV for K^+ production. The aim is to study the influence of the nuclear medium on elementary reactions. Information on high internal momenta, cluster phenomena and of two-step processes with intermediate pions is expected.
- Subthreshold Production of K^- Mesons (Letter of Intent 21, Ref. 9). Here even higher Fermi momenta can be probed than in the subthreshold production of K^+ mesons since higher masses are involved (associated production of K^- and K^+). Moreover, the creation of K^- mesons implies information about sea quarks and about the question whether the K^+K^- pairs are formed directly or via resonances. The K^- and K^+ momentum spectra will be measured below the NN threshold at 2.5 GeV.
- Exclusive Deuteron Break-up Study with Polarized Protons and Deuterons at COSY (Proposal 20, Ref. 10). The vector- and tensor-analyzing powers and the spin correlation asymmetry will be measured exclusively with a slow p under 180^0 and the fast projectile under 0^0 . The deuteron S- and D-state wave functions will be investigated at high internal momenta.
- Study of Neutral Mesons in Nuclei with a Two-Arm Photon Spectrometer (Letter of Intent 8, Ref. 11).
- Investigation of the $A(p, \eta)X$ Reaction in COSY (Proposal 23, Ref. 12). An interesting topic is a

possible production of η -mesons via the $N^*(1535)$ resonance in $p + C$ collisions with recoiling ^{12}C nuclei.

Additionally, the studies "Threshold Meson Production at the Internal COSY Beam in the Range of Scalar Mesons involving Strangeness" (Proposal 11, Ref. 13) which is planned to be performed in an early stage of the experimental program at COSY with the use of one of the COSY-ring dipoles might be continued under improved experimental conditions at TP2.

The envisaged experiments make full use of the possibilities offered by the 0^0 Facility. They all profit from the large solid angle. The K meson studies require small target-detector distance. Simultaneous spectroscopy of positive and negative ejectiles is necessary in the experiments on K^- production. The study of K^+ mesons requires the full flexibility for the selection of the momentum range of the ejectiles. The deuteron break-up measurements exploit the possibility to deflect ejectiles with D1 and D2 simultaneously.

There will be a basic set of detectors, but optimized additional setups will be necessary for the individual experiments. The layout for the measurement of correlated K^+d is shown in Fig. 2. The expected K^+ and d momentum spectra have been calculated¹⁴⁾ in the frame of the Modified Phase Space Model.

Fig. 2 outlines the expected trajectories of the K^+ mesons¹⁴⁾ and deuterons. Wire chambers will be used for the determination of the momenta and emission angles of the ejectiles. The measurement of the time-of-flight between the start counters (STC) and stop scintillators (SPC) and ΔE measurements either in the STC's, SPC's or in separate ΔE detectors will provide the necessary fast triggers. Time-of-flight and energy loss will also enable the separation of K^+ mesons from pions which are several orders of magnitude more frequent. The energy loss of the K^+ mesons will be optimized through the use of degraders. A prerequisite for the triggers is a momentum resolution in the order of a few percent of D2. The C and veto counters serve for background reduction. Since the detectors are tailored for K^+ mesons with a given momentum it is important that these momenta are focussed to the same position at the focal surface independently of the projectile energy. This can be achieved with the power supply which is independent of that of the ring magnets¹⁵⁾.

The details of the design of the 0^0 Facility, especially of the magnets and the detectors are presently under study.

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- h Institut für Kernphysik, Universität Münster, Germany

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Optimization of the particle-flight paths in the θ^0 Facility

M. Büscher, O.W.B. Schult, K. Sistemich

The θ^0 Facility¹⁾ at TP2 in the COSY ring will offer a high flexibility in the selection of the ejectile momenta and will transmit the COSY-beam in a wide momentum range. In this article we want to point out that a power supply independent from that of the COSY dipoles will offer additional advantages for the optimization of the detectors used — for example — for the measurement of K^+ mesons²⁾.

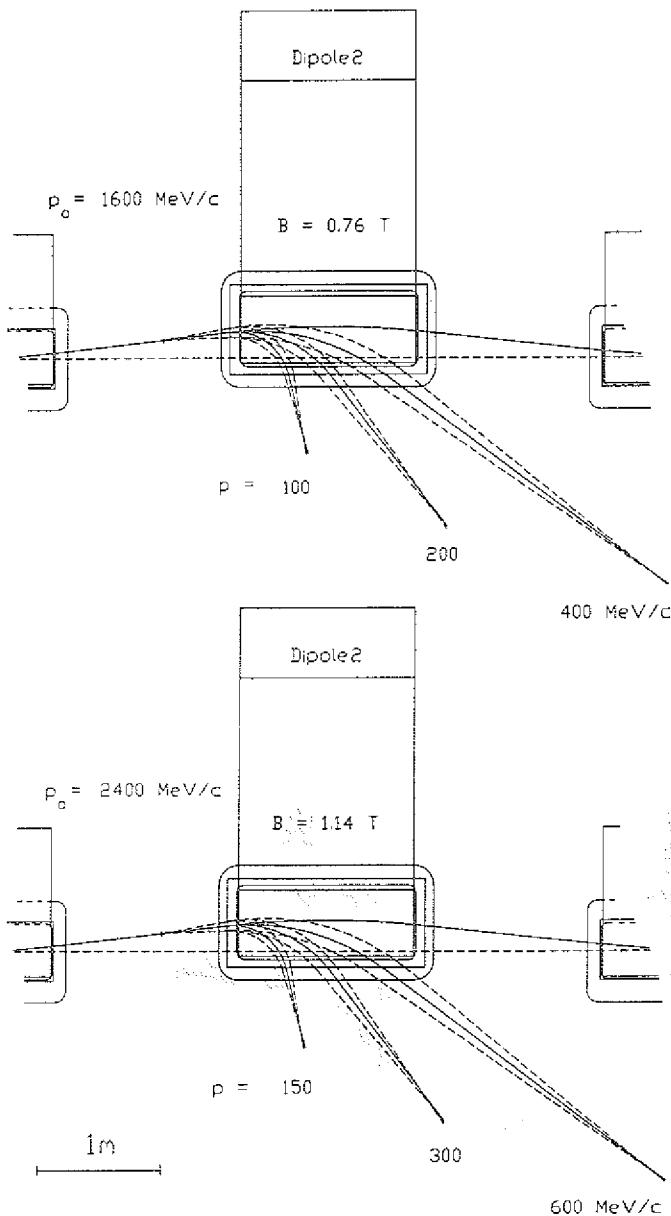


Fig. 1: Mode of operation of the θ^0 Facility with the field strength equal to that in the COSY dipoles. The radius of curvature of the circulating beam is constant ($\rho_0 = 700 \text{ cm}$) independent of the beam momentum. The ejectile momenta vary in correspondence with the beam momentum.

Fig. 1 shows the flight paths of the COSY beam and of ejectiles — emitted into forward angles between $\pm 4^\circ$ — through the facility. The field strength in the dipoles is chosen equal to that in the ring magnets and the radius of curvature ρ_0 of the circulating beam in the facility, therefore, is equal to that in the COSY dipoles (i.e. 700 cm). One can see that a wide momentum band of positively charged ejectiles is covered for both projectile momenta. The position of dipole 2 has been optimized in the sense that the COSY beam traverses the magnetic field around the magnetic symmetry axis in order to minimize influences from field inhomogeneities. This mode of operation could be realized if the facility is powered by the COSY-dipole power supply, but it has a serious drawback.

As depicted in Fig. 1 ejectiles which leave dipole 2 at the same position and under the same angle have different momenta for different beam momenta. This well known fact does not permit the use of an optimized system of detectors for measurements at different projectile momenta. However, such a detector is desirable as fast on-line trigger e.g. in the proposed 'subthreshold K^+ production' experiment²⁾. This trigger is based on TOF and ΔE measurements³⁾ and each detector element has to be optimized for one particle momentum.

Fig. 2 shows the mode of operation of the facility in a way which permits optimization of the detector, if an independent power supply is used. The field strength in the facility is now kept constant at e.g. 1.0 T whereas the COSY-beam momentum varies between e.g. 1.6 GeV/c and 2.4 GeV/c. The radius of curvature ρ_0 in the facility correspondingly varies between 533 cm and 800 cm. If the position of dipole 2 is adjusted such that in all three cases ejectiles with equal emission angles relative to the laboratory enter the magnet at the same position then ejectiles with equal momenta will leave the spectrometer at the same position and under the same angle, independent of the projectile momentum. This is possible since dipole 2 will be mounted on a support which allows horizontal displacement perpendicular to the direction of the unperturbed COSY-beam (indicated as a dashed line in Figures 1 and 2). It should be noted that the displacements compared to the mode of operation where the COSY beam lies around the magnet's symmetry axis (see Fig. 1) are relatively small (3.0 cm at $p_0 = 2.4 \text{ GeV/c}$, 1.8 cm at 2.0 GeV/c, 0 cm at 1.6 GeV/c) and it is ensured that

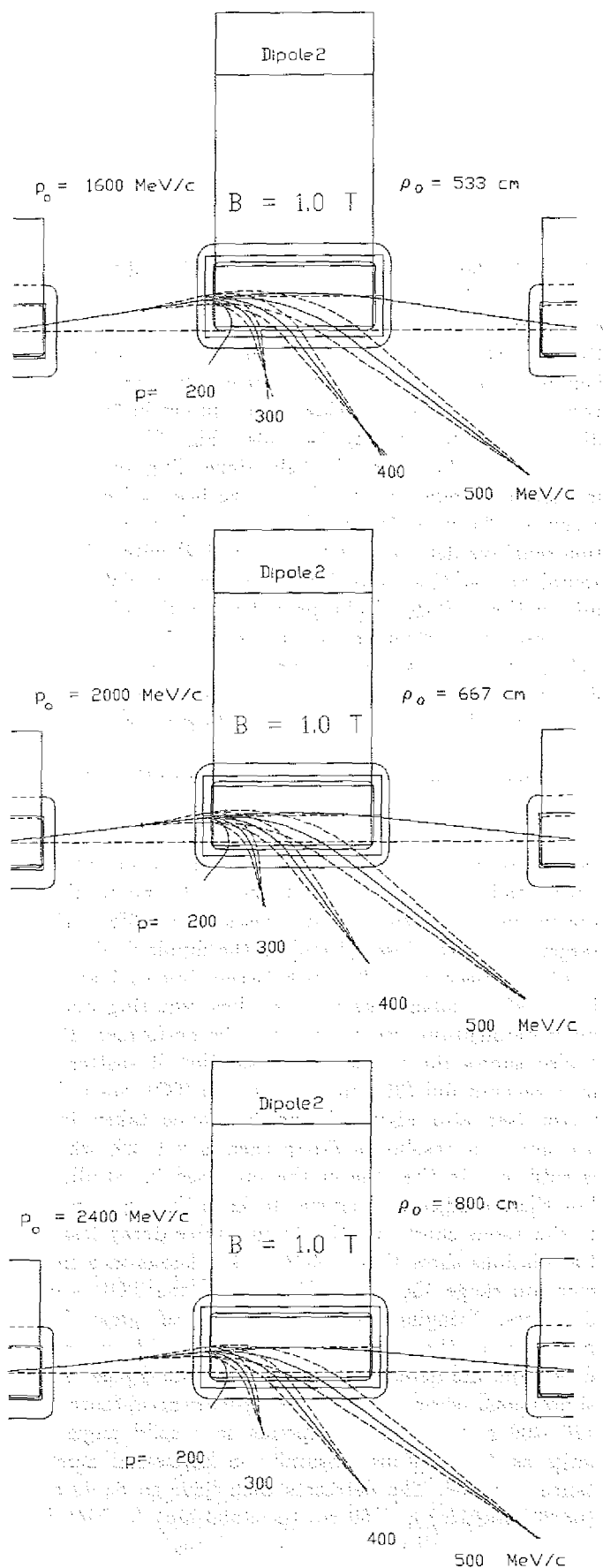


Fig. 2: Possible mode of operation using an independent power supply. The field strength is kept at 1.0T whereas the beam momentum is varied between 1.6 GeV/c and 2.4 GeV/c. Ejectiles with the same momentum leave dipole 2 at the same position and under the same angle in all three cases.

the circulating beam is not influenced by field inhomogenities.

Summarizing, an independent power supply for the 0^0 Facility permits higher flexibility of the facility and is indispensable for the optimization of the trigger detectors.

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Monte Carlo Simulations for the 0° Facility

M. Büscher

Monte Carlo simulations have been performed in order to determine the spectroscopic properties of the 0° Facility which is described in an accompanying article¹⁾ as well as the performance of various detector concepts for an on-line π^+ - K^+ discrimination. For the magnet simulations the raytracing program ZGOUBI²⁾ has been used³⁾ and subroutines for the detector responses have been added by the author.

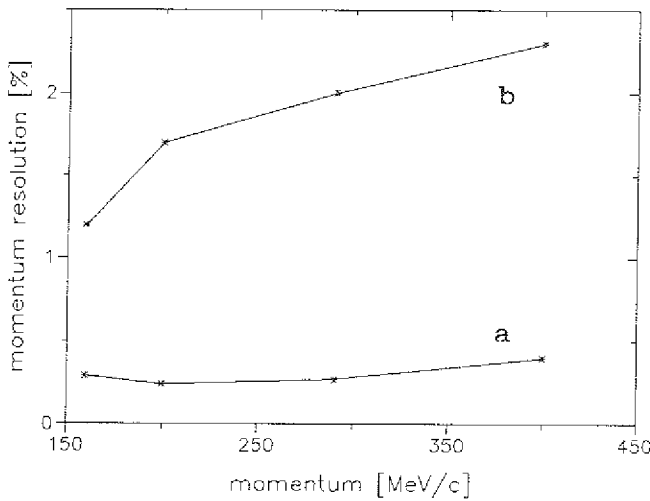


Fig. 1: Resolution of the 0° Facility for various particle momenta a) theoretically achievable resolution without scattering in vacuum foils or detectors, b) scattering taken into account.

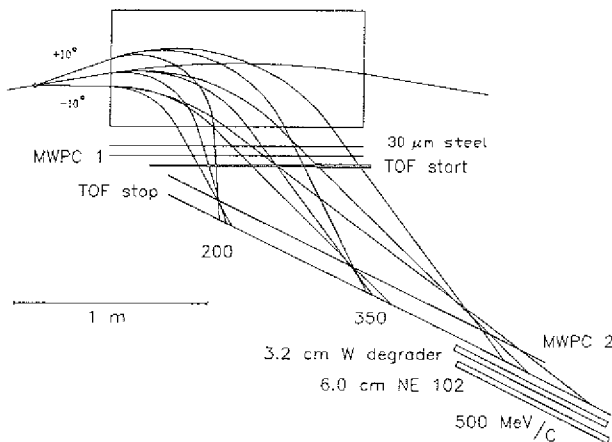


Fig. 2: Schematic view of the detector arrangement used for the simulations. The ejectiles pass through a 30 μ m steel window, a first position sensitive detector MWPC 1, the TOF-start detector and subsequently through the second position sensitive detector MWPC 2 and a TOF-stop detector. For high momentum particles also a W degrader and detection in a ΔE counter are simulated.

Fig. 1a) shows the theoretically achievable (i.e. without multiple scattering in vacuum foils or detectors) momentum resolution of the 0° Facility. During the simulation particles are emitted homogeneously under $\pm 10^\circ$ horizontally and $\pm 3.5^\circ$ vertically from a 2 mm wide target 40 cm in front of dipole 2 of the Facility (see also Fig. 2) and are deflected in the magnetic field depending on their momenta. Fringing field effects have been taken into account. The particles are then detected in two position sensitive detectors (MWPC 1 and 2) with a horizontal and vertical resolution of 1 mm FWHM. The information on the flight path behind the dipole is then used for backward raytracing treating the particle momenta as a free parameter. The momentum is determined by demanding that the reconstructed flight paths have their origin at the target region.

As shown in Fig. 1a) a momentum resolution $\Delta p/p$ as good as $2 \cdot 10^{-3}$ can be achieved. Similar calculations show that the emission angles from the target can be reconstructed with an uncertainty of $\sim 0.5^\circ$ horizontally and vertically. It should be noted that the momentum resolution decreases to $\sim 0.5\%$ if the target is located closer (7 cm) to the dipole field. The available space of ~ 120 cm between dipoles 1 and 2 is, therefore, advantageous if studies requiring maximum momentum resolution are to be performed. Fig. 1 also shows the momentum resolution if scattering in a vacuum foil (30 μ m steel) and a TOF start detector (see also Figure 2 and below) is taken into account. The resolution $\Delta p/p$ then is $\sim 1-2\%$ which is sufficient in the case of the proposed K^+ studies¹. For these studies it is crucial to keep the flight path of the kaons short in order to minimize decay losses. Calculations show that $\sim 40\%$ of the kaons in a momentum range 150–700 MeV/c reach the TOF-stop detectors. Fringing field effects are of great importance for the vertical acceptance of D2 and also determine the detector height behind the dipole. The simulations show that the vertical acceptance is 120–200 mrad. This corresponds to a solid angle as large as 40 – 70 msr assuming a horizontal acceptance of $\pm 10^\circ$. The detectors then have to be 40 cm ($p=200$ MeV/c) and 60 cm ($p=500$ MeV/c) high in order to cover all vertically accepted angles.

For the proposed K^+ studies an extremely efficient on-line trigger for π^+ - K^+ discrimination will be necessary since the π^+ rate is expected to be 10^5 – 10^7 times higher than the K^+ rate. The suppression of protons usually is less difficult; they can, for example, be stopped in degraders. An on-line π^+ suppression

factor of $\sim 10^4$ has to be achieved. There are two concepts which have been studied in Monte Carlo simulations: time-of-flight and ΔE measurements.

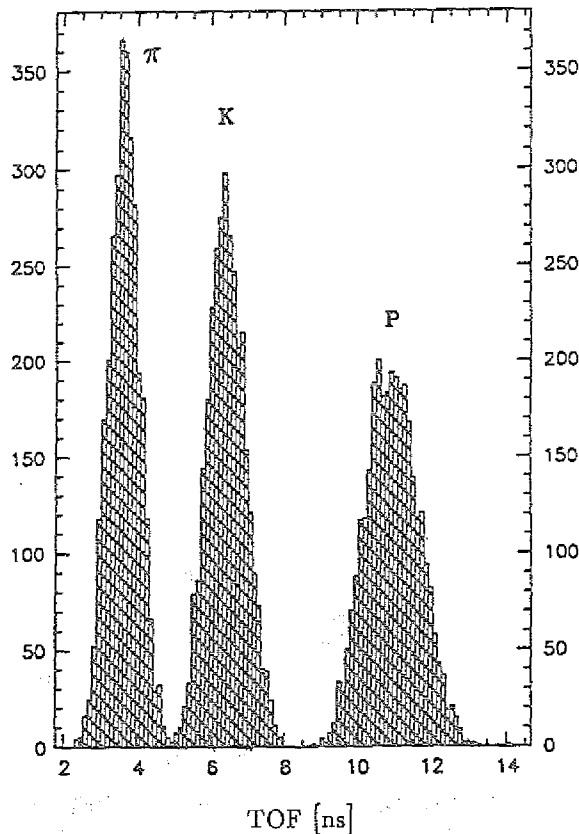


Fig. 3: Simulated TOF spectrum for pions, kaons and protons.

Fig. 3 shows a simulated TOF spectrum for pions, kaons and protons. The start and stop detectors are located as depicted in Fig. 2. The time resolution of each counter is 600 ps FWHM, the detector widths are 5 cm (start) and 12 cm (stop). It can be seen that the different particles are clearly separated. The distance between the pion and kaon peaks is $\sim 6\sigma$ where σ is the average peak width. This should allow a pion suppression by about two orders of magnitude. Clearly, the detector widths will influence the quality of the TOF trigger. With 10 cm wide start detectors the π^+-K^+ separation is only $\sim 5\sigma$, whereas 2 cm wide detectors improve the separation only marginally. Therefore, a start-detector width of ~ 5 cm seems to be the optimum choice.

For the simulation of the resolution of the ΔE counters a subroutine calculating Vavilov distributed^{4,5)} energy losses has been written. The result is the calculations for particles with $p=480$ MeV/c shown in Fig. 4. Included in the simulation is scattering and energy loss in the following detector components: 30 μ m steel foil, 0.3 cm TOF start detector, 5 cm H₂O Cerenkov detector (not shown in Fig. 2), 1.0 cm TOF stop detector, 3.2 cm W degrader and a 6.0 cm plastic ΔE counter. The thickness of the last two components has been optimized

in order to maximize the π^+-K^+ separation in the ΔE spectrum. A degrader is needed since pions and kaons with 480 MeV/c are both minimum ionizing which would result in similar ΔE spectra. Only after the particles are slowed down in the degrader the specific energy loss of pions and kaons is significantly different. It can be seen from Fig. 4 that pions and kaons are clearly separated but it should be noted that background from hadronic processes induced by π 's and p's is not included in the simulation. A pion suppression of at least two orders of magnitude, nevertheless, seems to be realistic. Calculations including the effects of hadronic processes as well as the influence of the background from scattered particles (e.g. from the pole shoes) using GEANT are under preparation.

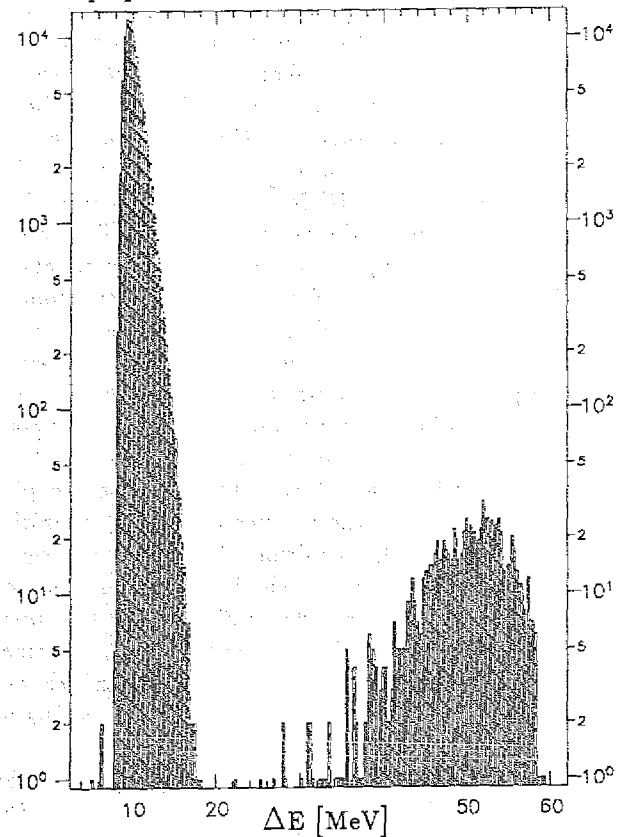


Fig. 4: ΔE spectrum of pions and kaons in a 6.0 cm thick NE102 detector. The ratio $n_{\pi}:n_K$ is 100:1.

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Subthreshold K^+ Production and the Fragmentation of the Residual Nucleus

H. Müller*, K. Sistemich

Subthreshold K^+ production is investigated in the framework of the modified phase-space (MOPS) model [1] implemented as a Monte-Carlo code, which samples complete events. An analysis of different reaction channels can be carried out simultaneously. As an especially interesting example we discuss here the interplay between subthreshold K^+ production and the fragmentation of the residual nucleus.

A proton-nucleus (pA) reaction is considered as a superposition of nucleon-nucleon interactions. In accordance with the well-known spectator-participant picture the events are classified by the number l of target nucleons involved in the reaction, and the corresponding cross sections are calculated using the Monte-Carlo code DIAGEN [2] based on a probabilistic interpretation of the Glauber theory. The interaction of the incoming proton with the participants is described by a two-step picture with the participants treated as a single unit called cluster. In the first step translational energy is converted into internal excitation energy of the projectile, the cluster and the residual nucleus, where the amount of excitation energy is determined by a "temperature" parameter Θ . The excited projectile- and cluster-like systems are pictured as consisting of the valence quarks of the initial nucleons. During the collision further quark-anti-quark pairs may be created which recombine randomly with the valence quarks building-up the final hadrons. The multiplicity of particles in the final state is governed by a radius parameter, which defines the volume occupied by the excited systems at the "break-up" point. In the second stage of the reaction the excited subsystems decay into the allowed channels with probabilities given by the corresponding statistical weights.

A K^+ meson may be produced if the excitation energy of the projectile or the cluster becomes high enough. The necessary energy of the projectile-cluster collision is reached due to the "weight" of the cluster consisting of several nucleons and its Fermi motion. According to the independent particle model the momentum distribution of the residual nucleus is approximated by a Gaussian with an l dependent width $s_l = \sqrt{l(A-l)/5(A-l)} \cdot p_F$, where for the Fermi momentum a value of $p_F=260\text{MeV}/c$ is used.

The de-excitation of the residual nucleus proceeds via break-up into several fragments. The parameters of the model defining the properties of the projectile- and cluster-like systems are taken from an investigation of the deuteron break-up at nuclei [1], while the parameters describing the residual nucleus are

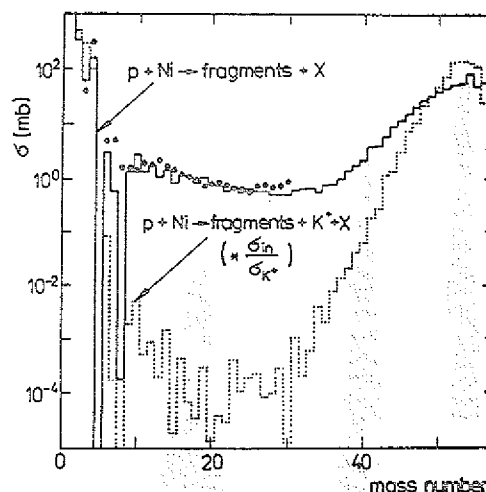


Fig. 1: Experimental mass spectrum [3] from the reaction $p+\text{Ni}$ at 1 GeV compared to MOPS [1] calculations. The full line represents the inclusive spectrum, the dotted one the mass spectrum in coincidence with K^+ (multiplied by $\sigma_{\text{inel}}/\sigma_{K^+}=2.3 \cdot 10^6$).

adopted such that for the interaction of 1 GeV protons with nickel the mass spectrum [3] (see Fig. 1) and the integrated cross section for K^+ production, $\sigma_{K^+}=(350 \pm 25)\text{nb}$, derived by interpolating the data of [4], are reasonably well reproduced.

Comparing the inclusive mass spectrum in Fig. 1 to the mass spectrum coincident with K^+ mesons we see that the shapes of the spectra drastically differ from each other. The relative yield of medium-mass fragments in the coincidence spectrum is decreased by about four orders of magnitude while the output of heavy and light nuclei is correspondingly increased. This result is plausible if we consider the distribution of the available energy between the various subsystems. The production of a K^+ meson requires a large amount of energy. Consequently, the energy left is not sufficient for breaking the residual nucleus into many pieces, it remains instead essentially intact with only few very light fragments evaporated.

It would be interesting to experimentally verify whether this strong effect predicted mainly on the basis of exact energy-momentum conservation within all the final particles is modified by dynamical peculiarities of subthreshold K^+ production.

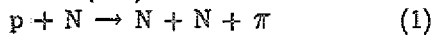
* Forschungsz. Rossendorf, O-8051 Dresden, Germany
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Deuteron Momentum Spectra accompanied by Subthreshold K^+ Production

M. Büscher, H. Müller*

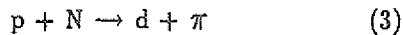
Recently Koptev et al. [1] measured the cross sections for subthreshold K^+ production at energies between 0.8 and 1.0 GeV for several target nucleons. These data were successfully reproduced by Cassing et al. [2] in the framework of a two-step model. In a first nucleon-nucleon (NN) interaction



an intermediate pion is created, and in a subsequent pion-nucleon (πN) reaction



the final kaon is produced. Ingredients of the model are parametrizations of the elementary NN and πN cross sections and a Hartree-Fock distribution for the momenta of the target nucleons. In the calculations of [2] the channel



has not been taken into account. Although the cross section $\sigma_{d\pi}$ of channel (3) is much smaller than the cross section $\sigma_{NN\pi}$ of channel (1), its importance for

the whole two-step process must be considered very carefully. Due to kinematics the energy of the pions arising from reaction (3) is larger than the energy of the pions produced via channel (1). This fact is decisive for the second step. In the vicinity of the threshold the cross section $\sigma_{K\Lambda}$ of process (2) strongly increases with increasing pion energy. Thus, the smaller cross section $\sigma_{d\pi}$ in the first step may be compensated by larger values of $\sigma_{K\Lambda}$ attainable in the second step. In order to verify this assumption the K^+ cross section has been calculated using the same parametrizations for the cross sections as in [2], but now we have channel (3) included. The momentum distribution of the target nucleons is parametrized by a Gaussian with a width of $\sqrt{1/5} \cdot 230$ MeV/c. Exact energy-momentum conservation is ensured by taking the four-momentum of the residual nucleus into account. We find that for the reaction $p+^{12}\text{C}$ at 1 GeV 66% of all K^+ mesons are produced via channel (3). Since the ratio $\sigma_{d\pi}/\sigma_{NN\pi}$ becomes larger with decreasing energy, the importance of channel (3) still increases at lower projectile energies.

The deuteron accompanying the K^+ is created in the quasi two-body process (3). It has a well-defined momentum, the value of which is smeared-out only by the momentum distribution of the target nucleons. Consequently, if subthreshold K^+ production proceeds as a two-step process then the deuteron momentum spectrum coincident with K^+ mesons should exhibit a distinct peak.

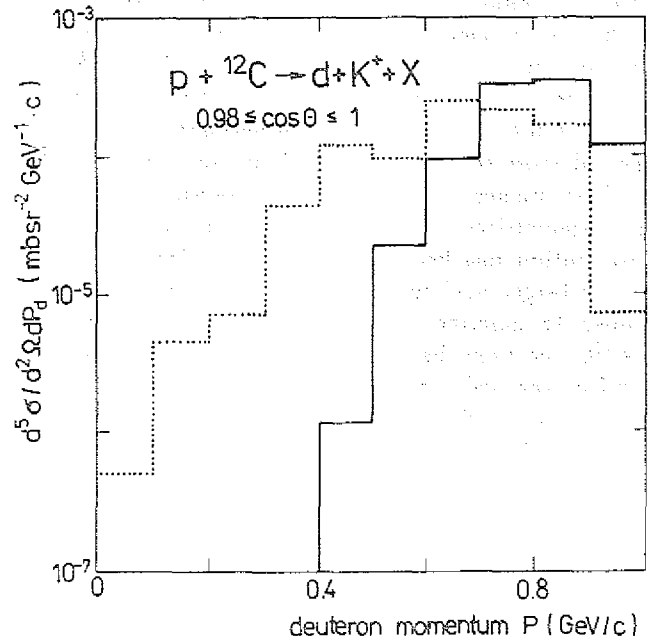


Fig. 1: Calculated deuteron spectra for the reaction $p+^{12}\text{C}$ at 1 GeV. Solid histogram: two-step process, dotted histogram: model [3]. The result from the two-step process is normalized to the integral over the curve [3].

In Fig. 1 model predictions for such a coincidence measurement are shown. From the two-step process we get the expected peak. This curve is compared to the result from the modified phase-space model [3], where subthreshold K^+ production is described by assuming that the projectile interacts with several target nucleons treated as a single unit called cluster. Among the decay products of a cluster there may be deuterons and other light fragments. Deuterons can be produced also via final-state interactions between nucleons, which are "near" enough in momentum space (coalescence model). A third source of deuterons is the fragmentation of the residual nucleus excited during the interaction process. Due to the more complex picture of deuteron production the resulting spectrum is much broader. So the measurement of a deuteron spectrum in coincidence with K^+ mesons would be helpful in clarifying the mechanism of subthreshold K^+ and of deuteron production.

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Model Predictions of K^+ Momentum Spectra and Angular Distributions from proton- ^{12}C Collisions at $E_p = 950$ and 1400 MeV

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and A. Strzalkowski^a

The K^+ production probability in collisions of protons with ^{12}C nuclei has been calculated within the folding model [1]. One step ($p+N \rightarrow N+K+\Lambda$) as well as two step processes ($p+N \rightarrow N+N+\pi$; $N+\pi \rightarrow K+\Lambda$) were taken into account. The experimental cross sections for the elementary processes in nucleon-nucleon or pion-nucleon collisions were used in parametrized form [2,3] while a Hartree-Fock distribution has been assumed for nucleon momenta in the target nucleus. The model has been previously shown to describe the inclusive K^+ cross section in $p+^{12}\text{C}$ reactions below 1 GeV [1]. Here we have extended the calculations with respect to an exact treatment of relativistic kinematics in all elementary channels. Moreover, another representation of results has been applied since it is more convenient to pre-

sent the differential cross section $d^2\sigma/dp d\Omega$ as function of emission angle and modulus of kaon momentum rather than Lorentz invariant differential cross sections ($E d^3\sigma/dp^3$) as function of longitudinal and transversal components of kaon momenta.

A realistic estimate of differential cross sections for the K^+ production is crucial for a program of experiments devoted to the study of subthreshold production of kaons in nuclei. It is very desirable to predict in which ranges of angles and momenta the various possible mechanisms (e.g. one-step or two-step) of the production should play a dominant role. Therefore, the angular distributions as well as spectra have been calculated separately for one-step and two-step processes. The calculations have been performed for

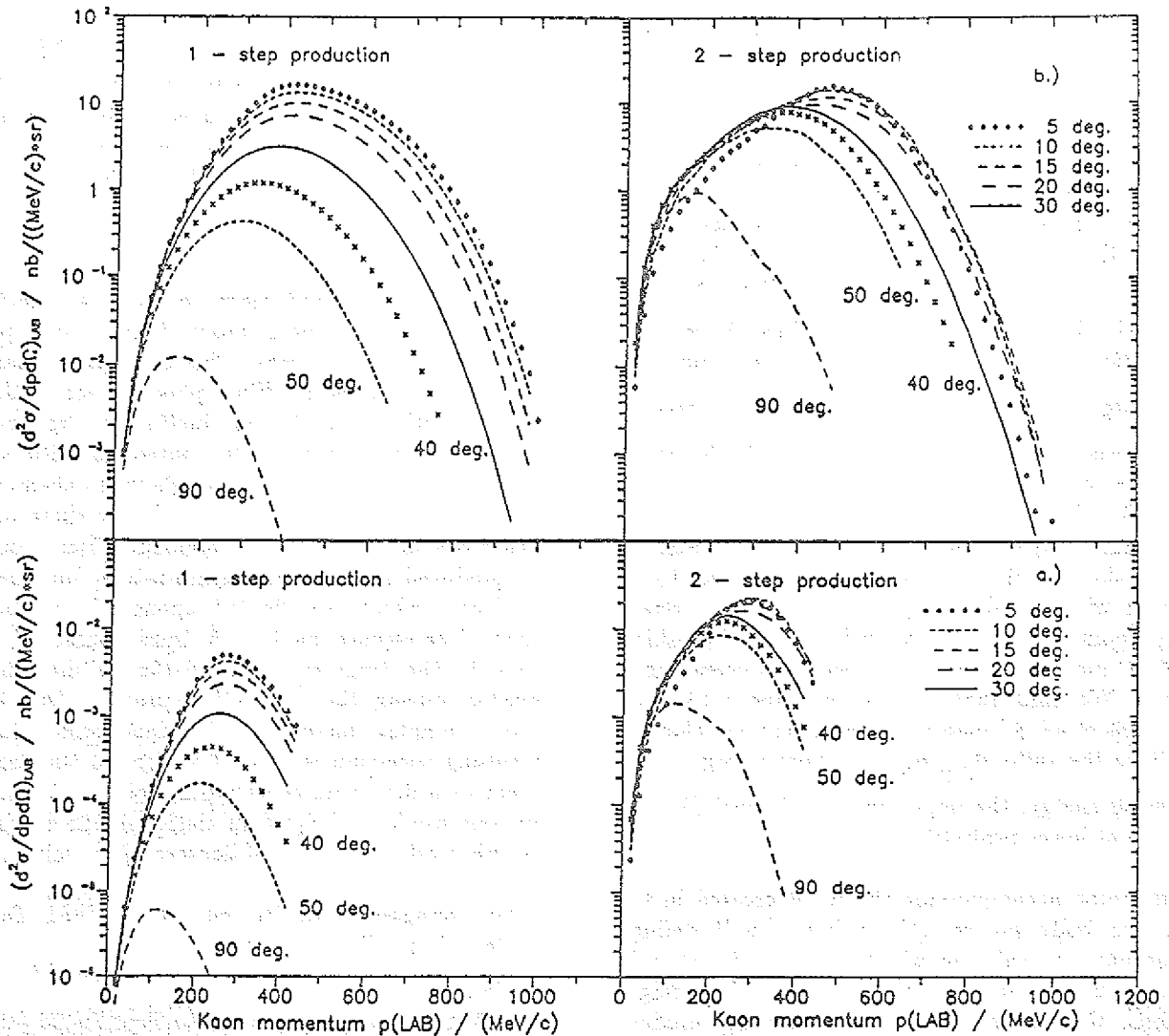


Fig. 1: Spectra of kaons evaluated for various LAB angles; the lower part of the figure (a) corresponds to lower beam energy ($E_p = 950$ MeV) the upper part (b) to $E_p = 1400$ MeV.

four different proton beam energies (950, 1100, 1250 and 1400 MeV). Since the results change systematically with the energy, only those for lowest and highest energy are shown in figs. 1 and 2.

It is evident from the spectra in fig. 1 that the cross sections are very small at low momenta (smaller than 100 MeV/c) as well as for the highest momenta achievable for a given beam energy. Furthermore, the spectra vary smoothly with emission angle. On the other hand the ratio of the cross section at kaon momentum corresponding to the maximum of the spectrum to that at a momentum smaller e.g. by 100 MeV/c is significantly different for one-step and two-step processes. Summarizing, one has to plan an experimental set up which is optimized for kaon momenta in the range of about 100 – 500 MeV/c (for beam energies of 0.9 GeV).

An inspection of fig. 2, which presents angular distributions of kaons for several kaon momenta, leads to the conclusion that the largest differences

between one-step and two-step mechanisms may be expected at very small angles (around 0 degree). At these angles the one-step mechanism has maximal cross sections while the two-step mechanism favours angles larger than about 20 degrees. This suggests that the detection system of kaons has to be able to measure in a rather broad range of emission angles in order to disentangle the mechanism of the kaon production.

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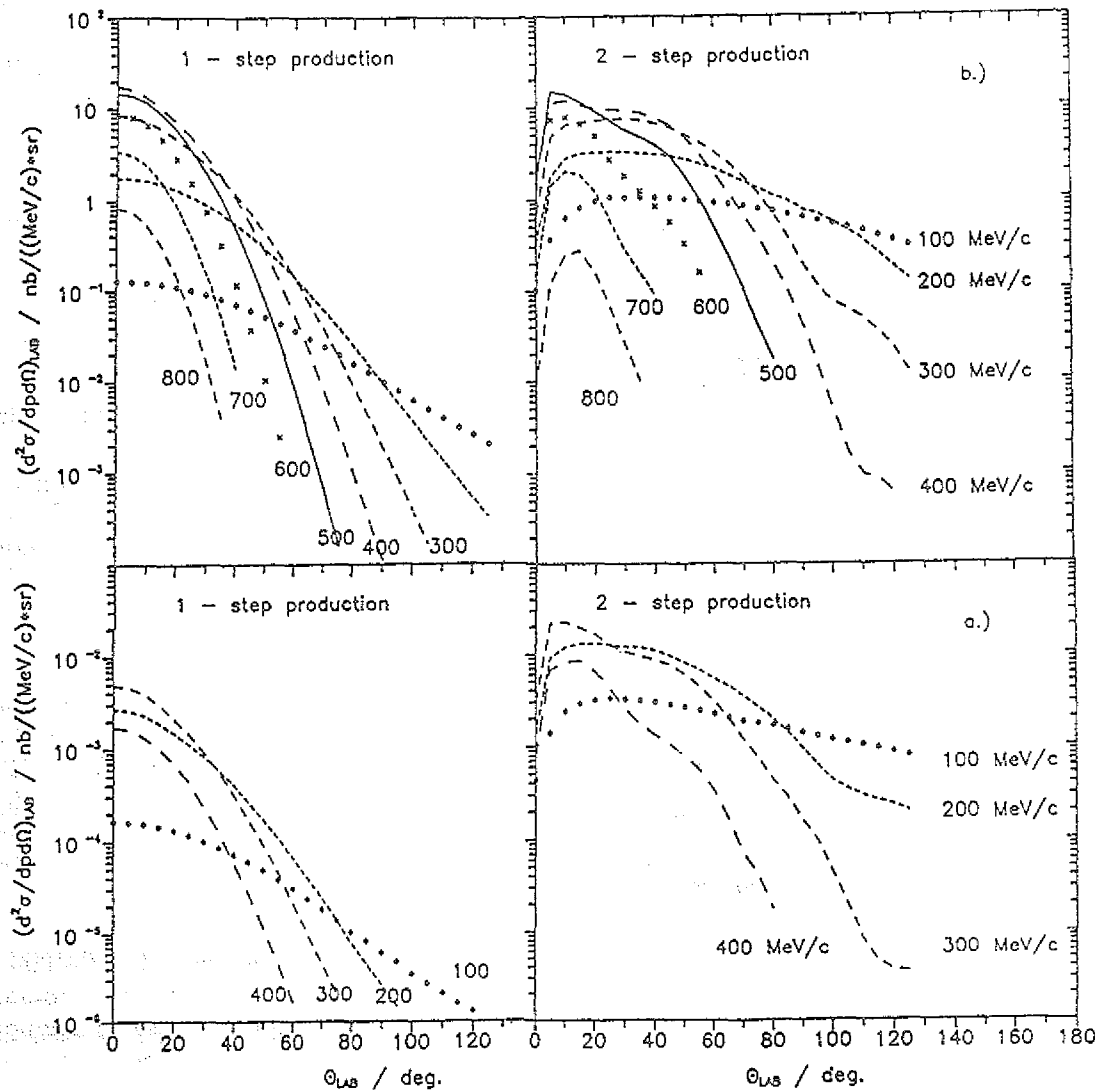


Fig. 2: Angular distributions of kaons evaluated for selected kaon momenta (in LAB); the lower part (a) corresponds to a proton beam energy $E_p = 950$ MeV, the upper part (b) to $E_p = 1400$ MeV.

Preparation for the COSY Experiment to produce very heavy Λ -Hypernuclei and to measure their lifetime

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J. Smyski^a, A. Strzalkowski^a, B. Styczen^a, P. von Brentano^b

The aim of the experiment is to measure the (p, K^+) production cross section of very heavy Λ hypernuclei, to obtain information on the probability for capture of the produced Λ by heavy and very heavy nuclei and on the probability of very heavy nuclei to survive the process of associated strangeness production. Finally we want to measure the momentum transfer in the production process and the lifetime of very heavy Λ hypernuclei, where $\tau(\Lambda)$ is determined essentially¹ by the four baryon weak interaction $N + \Lambda \rightarrow N + n$.

Details of the method which we want to employ and a sketch of the apparatus have been given in a contribution to the preceding annual report². The status of the work by May 1991 has been described elsewhere³. Figure 1 shows in a simplified manner a cut through the apparatus. The manipulator OMNIAX has been acquired from VG Instruments. It permits moving (up to 600 mm $\pm$ 0.02 mm) the target horizontally into the circulating beam (from the right to the left in Fig. 1) and it allows precise adjustment (up to $\pm$ 25 mm $\pm$ 5 μ m) of the target in vertical direction, so that the lower target edge coincides with the axis of the beam, and longitudinally, in the direction of the beam, so that the shadow can be adjusted in such a way that prompt fission fragments are suppressed as desirable by moving the shadow edge with the target relative to the shadow counter. The manipulator was set up in the laboratory and tested. It can be operated manually or under computer control. After baking at 150°C a vacuum of 10^{-9} mbar was obtained. Also most of the parts were purchased needed for changing targets under UHV conditions.

The detector system to be used for the measurement of the fission of heavy Λ hypernuclei has been set up in the laboratory and is presently being tested. Since the detectors will be operated inside the COSY vacuum for the tests a vacuum system including a 50 l vessel for the detectors has been set up where a pressure of $6 \cdot 10^{-5}$ mbar can be maintained with the detectors in operation.

For the detection of Λ delayed fission fragments two position sensitive low-pressure multiwire proportional chambers have been constructed by H. Stelzer⁴. The active areas of the detectors are 34x17 cm² and 20x10 cm², respectively. The cathode consisting of a 1.5 μ m thick mylar foil with 20 μ g/cm² Ag layers on each side is connected to high voltage. At a distance of 3 mm it is sandwiched between planes of anode wires (20 μ m diameter W wires with Au coating) which have a 1mm wire spacing, one plane with wires parallel to the short sides, the other plane with wires parallel to the long sides of each detector and perpendicular to the circulating beam. The detector windows (each at 12 mm distance from the cathode) are made of 2 μ m (300 μ m/cm²) Toray (Polyamid, $\rho = 1.5$ g/cm³) film.

Position determination is obtained through delay line read out. In the bigger counter the anode wires are connected in groups of 5 or 2, along the longer and shorter sides, respectively. Analogous grouping for the smaller detector is 2 and 1 wire. Each group is connected with a delay-line tap limiting the position resolution to the spacing of these wire groups. The timing signal is taken from the cathode. Delay-line chips and fast preamplifiers are mounted directly on the detector housings.

The detectors are filled with isobutane and the optimum pressure of about 4 mbar is kept constant by an electronic gas pressure control unit GPCU Mod. 8402. This unit measures the pressure with a Honeywell Micro Switch 164PC pressure sensor and adjusts it by operating an electromagnetic valve Brooks model 5835N.

Data acquisition is done with the FATIMA system⁵ which is based on the LeCroy Fera system. It is run on a Microvax computer which is equipped with a Exabyte tape drive for the storage of mass data.

Detector tests are carried out with fragments from the spontaneous fission of ²⁵²Cf. Fig. 2 displays the position distribution for the 10 cm coordinate of the

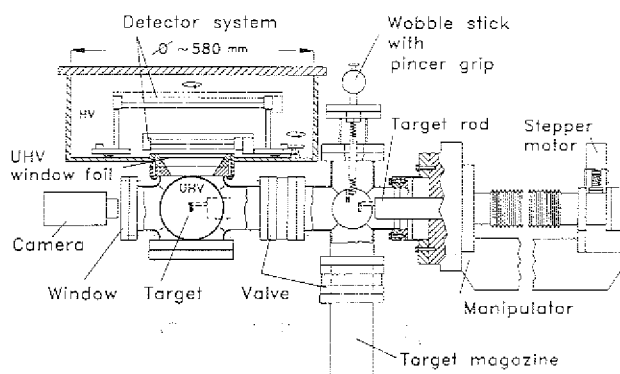


Fig. 1: Simplified drawing of the apparatus for the fission study of heavy hypernuclei.

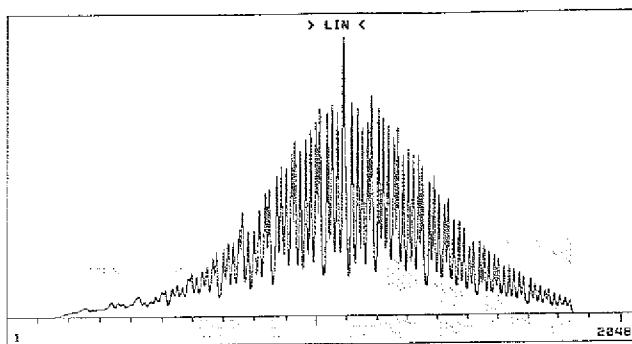


Fig. 2: One-dimensional position spectrum of ^{252}Cf fission fragments.

small detector illuminated with the source placed at a distance of 5 cm above the center of the detector. The overall shape of the distribution is given by the geometry whereas the fine structure reflects individual anode wires with 1 mm spacing. This distribution is measured with the TDC start signal from the cathode and the stop signal taken from one end of the delay line. The position resolution is expected to be further improved when timing signals from both ends of the delay line are used and the position information is deduced from the difference of the delay times from both ends of the delay lines.

The possibility to measure two-dimensional position spectra is demonstrated with a scatter plot of ^{252}Cf fission fragments which passed through a Λ -shaped mask before entering the detector, see fig. 3. The few

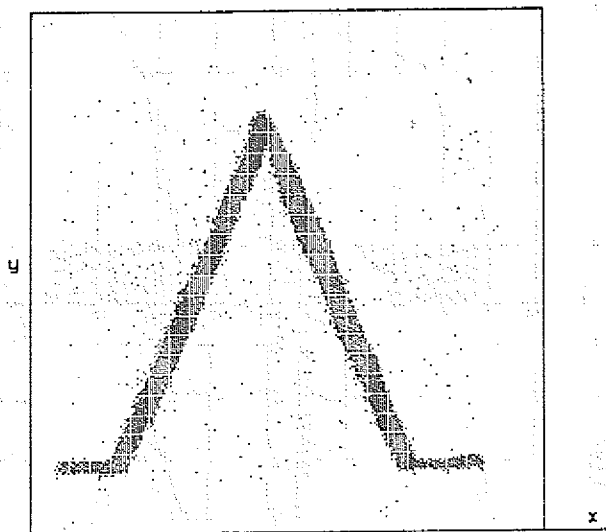


Fig. 3: Two-dimensional distribution of ^{252}Cf fission fragments passing through a Λ -shaped opening of a mask. The pattern of white lines appearing across the symbol stems from the structure of the screen from which this hardcopy was made.

events outside the Λ -symbol will probably disappear when proper gates are set on pulse-height distributions and on the sum of the delay times from both ends of each delay line which has to be constant for single-hit events.

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Construction and Tests of Avalanche Detectors

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As was described in details earlier [1], the experiment COSY #13 aims at the study of the Λ -decay induced fission of uranium or other very heavy nuclei. The so called recoil shadow method will be used to detect the fragments, eliminating the background coming from prompt fission and other reactions.

The prompt fission fragments should not reach the MWPC's located in the shadow. Monitor detectors placed at the front edge of the first MWPC will enable a precise monitoring of the target position relative to the shadow edge and allow to align the MWPC's in respect to the target-shadow edge system. These detectors will be position sensitive (in one direction) with a position resolution better than 1 mm. The active area of this detector should be 8 mm $\times$ 16 mm.

In the later stage of the experiment the fragment coincident to the one detected in the shadow counter will be measured in a large coincidence detector [2]. This detector should also be position sensitive in two directions with position resolution of a few millimeters and with good timing resolution (about 1 ns). Other required properties of this detector have been described elsewhere [2], [3].

Considerations concerning the type of these detectors led to the conclusion that both, the small monitor detectors and the coincidence detector should be constructed as position sensitive, low pressure avalanche counters. Detectors of this kind are well suited for experiments with heavy ions and fulfill essentially all special requirements of the COSY #13 experiment.

In 1991 the work on prototypes of avalanche detectors has been undertaken in the Detector Laboratory of the Nuclear Physics Department of the Institute of Physics of the Jagellonian University. Two detectors have been built and tested.

1. Parallel Plate Avalanche Counter (Fig. 1) with electrodes made of metalized mylar foil spaced by 2 mm and with windows out of the same foils. The active area of the detector has a diameter of 30 mm. The detector was filled with 90% Ar + 10% CH₄ or with isobutan. It was found that the Ar + CH₄ mixture was not suitable for the avalanche detector. The detector was tested at different gas pressures and voltages with 6.05 and 8.77 MeV α particles from ²¹²Bi.

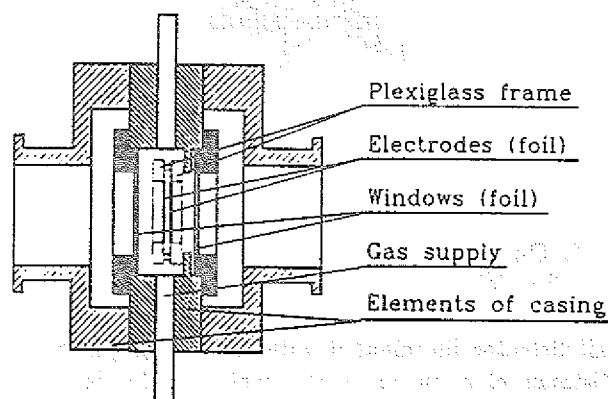


Fig. 1: Schematic design drawing of the prototype PPAC.

The aim was to find the optimal working parameters of the detector i.e. the pressure and voltage which yield 100% efficiency and the best timing characteristics. Figure 2 shows the obtained efficiency of the PPAC as a function of voltage for several gas pressure values. The pressure of 12 Torr and the voltage of 750 V were found to be optimal. Under these conditions the PPAC has a small effective thickness while the signal amplification is the best without breakdown risk. The fast output signal contained only the electron part. The slow ion part of the

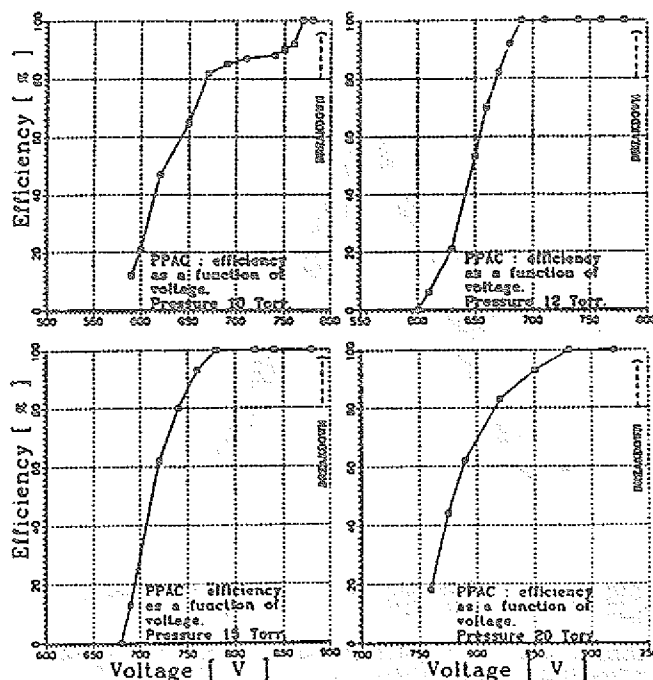


Fig. 2: Efficiency curves of the prototype PPAC.

signal was suppressed by proper choice of the time constant $\tau = R \cdot C$ of the detector.

2. A Double Grid Avalanche Chamber has been built as a prototype of a position sensitive detector (Fig. 3). Two wire planes (acting as anodes) give the position information in X and Y direction. Between them is located a common cathode made of a metalized mylar foil, delivering the time signal. The distance between the cathode and each anode plane is 4.5 mm and the distance between anode wires is 1.5 mm. The anode wires are connected in pairs. The detector volume is limited by two mylar foil windows. As in the case of the PPAC the gas used to fill the DGAC was isobutan.

Results obtained from the DGAC efficiency tests and its output signal are similar as in the PPAC case. A pressure of 16 Torr and the voltage of 680 V are the best working parameters of the DGAC. Differences to the PPAC in performance parameters are due to the different geometry of the detectors. The position resolution of the DGAC is found (as expected) to be equal to the distance between pairs of connected anode wires (3 mm).

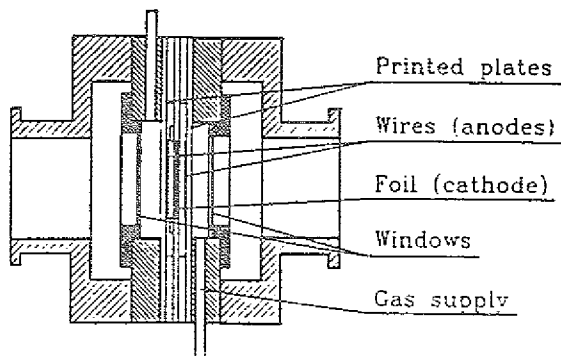


Fig. 3: Schematic design drawing of the prototype DGAC.

On the basis of the gained experience a Parallel Plate Avalanche Counter was constructed to be used for testing the properties of the shadow detectors [4]. This PPAC consists of two electrodes made of metalized mylar foil acting simultaneously as windows limiting the detector volume. The distance between these electrodes is 1.5 mm. With the isobutan pressure of 4 mbar and the voltage of 430 V the rise time of the output signal is 1 ns.

We now plan to prepare a prototype of a DGAC with individual wire readout, finding or developing the optimal preamplifier for the signal extraction and finishing on that basis the final design of the monitor detectors.

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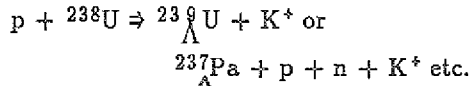
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Simulation of the Passage of Heavy ^{239}U Hypernuclei through a thin ^{238}U Target

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The measurement of the lifetime of the Λ hyperon within nuclear matter is one of the aims in the planned experiment COSY #13 [1]. Heavy hypernuclei will be produced in the reaction



A signature of the hypernucleus formation will be the detection of fragments of fission caused by the Λ decay in the region not accessible for fragments from prompt fission ("in shadow", cf. [2]). This experimental procedure allows to measure a fission fragment distribution which reflects the length of the flight path of the hypernuclei outside the target.

The distance which a hypernucleus can cover before it decays is determined by the hyperon lifetime and the momentum the hypernucleus possesses while leaving the target. Since the $^{239}\Lambda\text{U}$ (or $^{237}\Lambda\text{Pa}$) hypernucleus is produced within the target volume its momentum and direction of motion at the moment of emerging from the target can differ from the initial values due to interaction with the target material (energy losses and small angle scattering). In order to unfold from the measured distribution the value of the Λ lifetime it is necessary to study the influence of these effects.

Data concerning the energy losses of heavy ions in a heavy target in the low energy range (the initial energy of $^{239}\Lambda\text{U}$ is < 2.5 MeV) are scarce and the information cannot be simply taken from literature. Thus in order to find the magnitude of the straggling

effects one has to perform a computer simulation of the process of propagation of the heavy hypernuclei through the ^{238}U target.

The main topic of this work was to find a distribution $N(p, \Theta)$ of momentum, p , and angle, Θ , of hypernuclei leaving the target. We used the method of Ziegler, Biersack and Littmark [3] to calculate the energy loss and straggling of the hypernuclei. It was assumed that the hypernucleus is produced at $\Theta_{\text{LAB}}=0^\circ$ with equal probability at any point along the one-dimensional beam within the thin target. Following these assumptions a computer code was written which simulates the passage of hypernuclei through the target.

The target thickness was chosen to be 20, 30, 40, 50 and $100 \mu\text{g}/\text{cm}^2$. For each one two runs were performed assuming two different initial momenta of the heavy hypernuclei: $p_{\text{LAB}}=1040 \text{ MeV}/c$ and $630 \text{ MeV}/c$. Production of 10^5 hypernuclei was simulated in each run. Table I presents the fraction (in %) of the

Table I: Fraction of hypernuclei leaving the target at $|\Theta| \leq 64^\circ$.

Target thickness [$\mu\text{g}/\text{cm}^2$]	p_{LAB} 1040 MeV/c	p_{LAB} 630 MeV/c
20	97.0%	93.0%
30	95.5%	88.9%
40	93.7%	84.2%
50	91.8%	79.4%
100	79.6%	54.5%

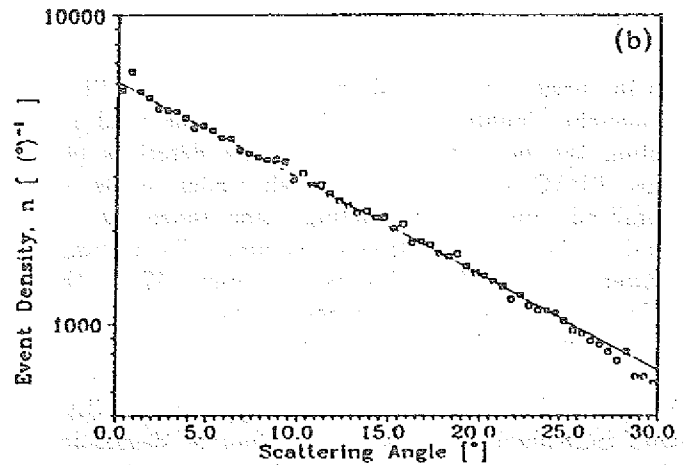
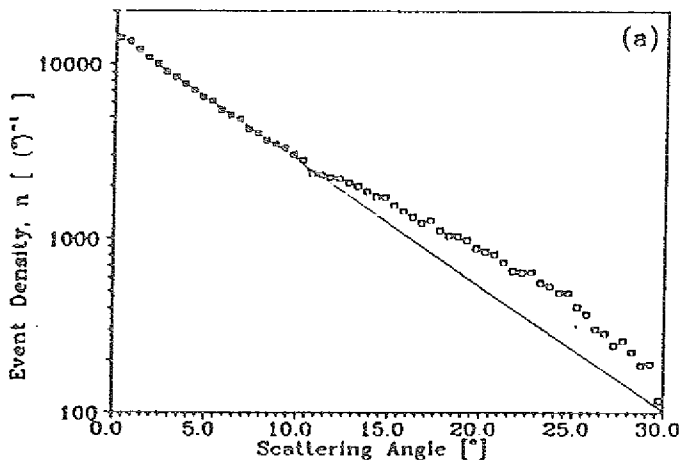


Fig. 1: Projections of simulated distributions, $N(p, \Theta)$ on the Θ -axis for $50 \mu\text{g}/\text{cm}^2$ for different assumed initial momenta, p_{LAB} : 1040 MeV/c (a), 630 MeV/c (b).

hypernuclei emerging from the target within the angular range $0^\circ \leq \Theta \leq 64^\circ$, having momentum larger than 920 MeV/c (for $p = 1040$ MeV/c or 520 MeV/c (for $p_{LAB}=630$ MeV/c). Figs. 1a and 1b present projections of the simulated distributions $N(p, \Theta)$ for $p_{LAB}=1040$ MeV/c and $p_{LAB}=630$ MeV/c, respectively, ($50 \mu\text{g}/\text{cm}^2$ target) on the Θ -axis.

Table II: Parameter α .

Target thickness [$\mu\text{g}/\text{cm}^2$]	p_{LAB} 1040 MeV/c $\alpha[^\circ]$	p_{LAB} 630 MeV/c $\alpha[^\circ]$
20	2.6	5.7
30	3.8	8.1
40	5.0	11.2
50	6.1	13.9
100	11.7	28.5

These projections can be represented by a curve of the form:

$$n = \frac{dN}{d\Theta} = (N_0/\alpha) \exp(-\Theta/\alpha).$$

A fit was made for points with $\Theta \leq 10^\circ$ (full squares) in order to reproduce as well as possible the simulated distribution in the region of small angles. The α parameter yields information about the spread of the direction of the hypernucleus momentum vector. Values of this parameter for different target thicknesses and two initial hypernuclei momenta are given in Table II. Additionally, plots of the momentum of the hypernuclei within the angular range $0^\circ \leq \Theta \leq 5^\circ$ are presented in Figs. 2a and 2b. The ratio of the number of events within $0^\circ \leq \Theta \leq 5^\circ$ and with momentum larger than 920 MeV/c or 520 MeV/c, respectively, to the total number of events is given in Table III. These results form a good basis for choosing the target thickness.

Table III: Fraction of hypernuclei with $|\Theta| \leq 5^\circ$ and $p > 920$ MeV/c for $p_{LAB}=1040$ MeV/c and $p > 520$ MeV/c for $p_{LAB}=630$ MeV/c.

Target thickness [$\mu\text{g}/\text{cm}^2$]	p_{LAB} 1040 MeV/c	p_{LAB} 630 MeV/c
20	74.4%	50.8%
30	65.6%	40.3%
40	58.3%	35.3%
50	52.0%	27.0%
100	33.1%	15.0%

Two additional runs were made for a $50 \mu\text{g}/\text{cm}^2$ target to obtain better statistics. Production of $5 \cdot 10^5$ hypernuclei was simulated in these runs. It was then possible to make projections of two-dimensional distributions $N(p, \Theta)$ on the p -axis within strips of 0.5° width ($0.0^\circ-0.5^\circ$, $1.0^\circ-1.5^\circ$, $2.0^\circ-2.5^\circ$, and so on). Results for 0° and 5° are shown in Fig. 3a (1040 MeV/c) and Fig. 3b (630 MeV/c). Average values, $\langle p \rangle$, and standard deviations, $\sigma(p)$, are given in Table IV.

Table IV: Average values of p and $\sigma(p)$.

$\Theta_{LAB}[^\circ]$	p_{LAB} 1040 MeV/c		p_{LAB} 630 MeV/c	
	$\langle p \rangle$ [MeV/c]	$\sigma(p)$ [MeV/c]	$\langle p \rangle$ [MeV/c]	$\sigma(p)$ [MeV/c]
0.0-0.5	1037	2.2	628	1.4
5.0-5.5	1020	11.1	615	11.0
10.0-10.5	1011	18.8	605	17.4

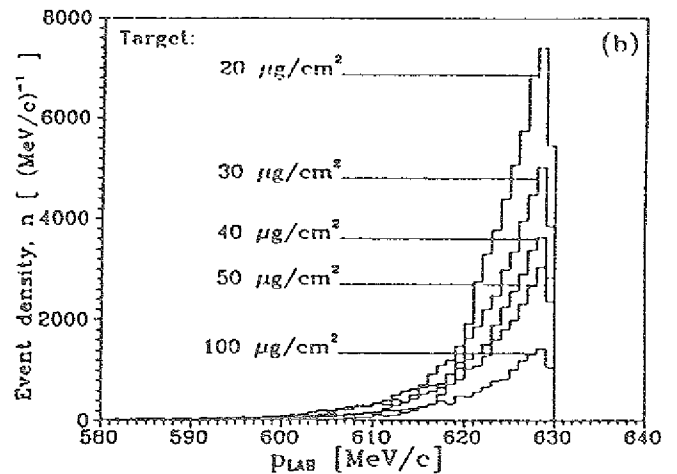
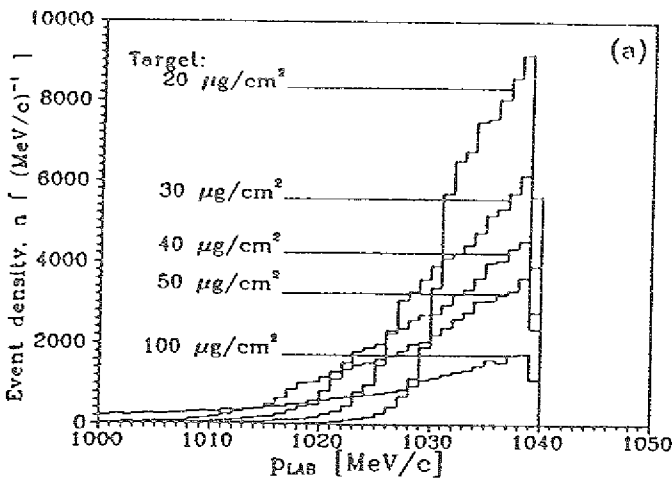


Fig. 2: Projections of simulated distributions, $N(p, \Theta)$, on the p -axis for events within $0^\circ \leq \Theta \leq 5^\circ$ for different target thicknesses and different initial assumed momenta, p_{LAB} : 1040 MeV/c (a), 630 MeV/c (b).

It can be concluded that for thin targets (thickness $\leq 50 \mu\text{g}/\text{cm}^2$) the momentum of the hypernuclei emerging from the target deviate only very little from the value with which they are produced.

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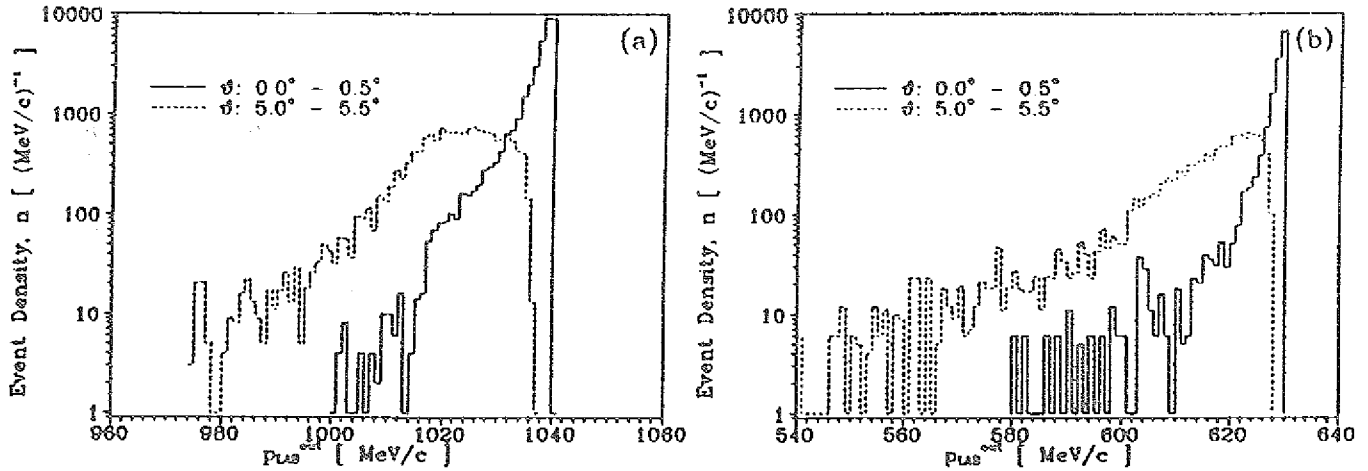


Fig. 3: Projections of simulated distributions, $N(p, \Theta)$, on the p -axis for events within $0.0^\circ \leq \Theta \leq 0.5^\circ$ and $5.0^\circ \leq \Theta \leq 5.5^\circ$ for a target thickness of $50 \mu\text{g}/\text{cm}^2$ and different initial assumed momenta, $p_{\text{LAB}}: 1040 \text{ MeV/c}$ (a).

Time Structure of Scintillators for the COSY Two-Arm Photon spectrometer

J. Fabian, H.R. Koch, D. Kopyto and I. Strakovsky*

The maximum rate of particles hitting a detector is of course a very important detector design parameter. If a thick target, as 1 g/cm^2 of carbon, is exposed to the external COSY beam of $4 \cdot 10^9$ protons/s, a luminosity of $2 \cdot 10^{32}/\text{s} \cdot \text{cm}^2$ is achieved. At an internal irradiation position luminosities of this size can easily be realized even with very thin targets. Then $\approx 4 \cdot 10^7/\text{s}$ reactions occur in the target and — for proton beam energies around 1 GeV — about $2 \cdot 10^8/\text{s}$ secondary particles are emitted into 4π .

The photon detector, which is proposed for the study of eta production in nuclei at COSY [1], is composed of scintillator elements with an acceptance of $\approx 60^\circ$, corresponding with a detection efficiency of $0.87 \cdot 10^{-3}$. Thus each element registers $\approx 10^5$ signals per second, corresponding with a mean signal distance of $5 \text{ } \mu\text{s}$. If one has to determine pulse heights precisely, the main decay time constant of the scintillator should be in the order of $<1\%$ of the mean signal interval to make pile up effects and baseline shifts induced through preceding large signals sufficiently small.

The use of the photon spectrometer for neutral meson detection through the measurement of 2 photons, emitted in meson decay, requires precise energy determination. Therefore, the scintillation time constants should be $<50 \text{ ns}$, if one wants to make use of the full luminosity, obtainable with COSY. Weak additional, longer time constants can be corrected in principle electronically by pulse shaping amplifiers with appropriate pole zero cancellation circuits, or through the multichannel measurement of the time structure of every scintillation signal. However, the costs of the electronic equipment of the spectrometer increase strongly, if one compares with simple registration with one or two charge sensitive ADC-s per scintillation element. In any case, the presence of slow scintillation light will decrease the dynamic range of the energy measurement. The ratio of maximum to minimum measurable pulse heights should be in the order of 400 for a good photon spectrometer. About 400 MeV may be deposited in a single detector element in the center of a shower caused by a photon from the decay of an eta meson, which was produced at the nucleon nucleon threshold energy $T(p)=1258 \text{ MeV}$. Typical energy depositions in the outer region of an electromagnetic shower are only in the order of 1 MeV per detector element.

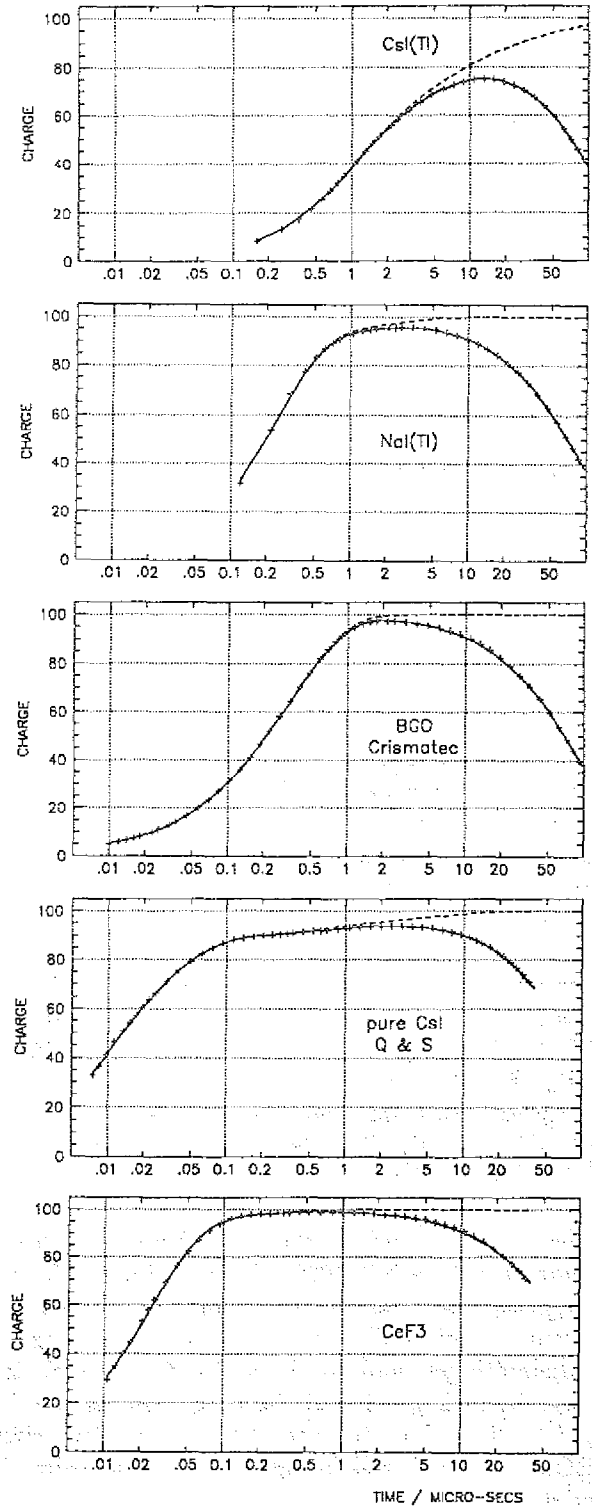


Fig. 1: Time structure of RC-integrated photo-multiplier currents for different scintillators. +++: measured values; —: result of fit of exp-functions; - - -: corrected for RC-charge loss, normalized to 100 for $t=\infty$.

Crystal	Integration Period / μs								
	0.1	0.2	0.5	1	2	5	10	20	50
CsI(Tl)	5	11	24	39	55	72	81	88	94
NaI(Tl)	27	51	82	94	97	99	100	100	100
BGO Crismatec	31	49	77	94	100	100	100	100	100
CsI Bicron Z	68	75	81	83	85	88	92	95	99
CsI Bicron E	85	88	91	93	96	98	99	100	100
CsI Q & S 57	87	90	92	94	96	97	99	100	100
CeF-3	95	98	100	100	100	100	100	100	100

Table 1: Charge at photomultiplier anode after different time intervals, normalized to 100 at $t = \infty$.

The fast scintillators of high Z and sufficient size, which are actually available, are BaF_2 and pure CsI [2]. BaF_2 is specially interesting for experiments, where fast timing is needed as in the TAPS spectrometer [3] at the SIS ring at the GSI. However, this material has also a 625 ns component, which contains 80% of the light intensity [4]. Pure CsI has only components of ≈ 10 ns and ≈ 40 ns. Longer time constants are also observed. They are due to impurities in the scintillation material. At present we are building a small test detector, consisting of 25 pure CsI crystals, representing a subset of the detector proposed in [1]. Examination of several elements, which have already been delivered, shows, that the intensities of the slow scintillation components, measured in different crystals, vary considerably. We have, therefore, made a series of measurements, to determine the time structure of scintillation pulses of pure CsI crystals [5]. For comparison also the slower scintillators CsI(Tl), NaI(Tl), BGO, and the new fast scintillator CeF_3 [6] were included. The pulse shapes were measured with a fast digitizing oscilloscope (sampling rate ≤ 500 MHz). The relevant time constants in the region from ≈ 10 ns to $\approx 100 \mu\text{s}$ were deduced.

Figure 1 shows some of the pulse shapes, which were measured with a XP2020Q photomultiplier. Averages are given of signals, which were induced by minimum ionizing cosmic radiation. In Table 1 the integrated photomultiplier current is given for several time intervals, each starting at the time of signal rise. The charge loss due to the RC integrating circuit ($RC = 104 \mu\text{s}$) was compensated numerically. We see strong very slow components in the time structure

of the CsI(Tl) signals. A fit of exponential functions to the measured line shapes yields the following time constants and intensities (μs ; %): 0.90;47, 3.8;26, 14;16, 77;11. This material is certainly inadequate for countrates of $\geq 1 \cdot 10^4/\text{s}$. The pure CsI crystal Bicron Z seems to be one of the early products of this company. The crystals (Bicron E) and (Quartz&Silice 57) exhibit a very good fast/total ratio. Material of this type is suited for our purposes. Also CeF_3 shows excellent scintillation properties. For CeF_3 we deduce the following signal composition (μs ; %): 0.025;95, 0.20;5. At the moment only relatively small crystals of this type are available. We tested a crystal of $15\phi \cdot 25 \text{ mm}^2$, obtained from the St. Petersburg Institute of Nuclear Physics.

* Institute of Nuclear Research St. Petersburg

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COSY-11 Detector Arrangement

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COSY-11 is an internal target experiment at COSY for the threshold meson production in the mass range up to $1 \text{ GeV}/c^2$.

The planned setup of the experiment is shown in fig. 1 together with the tracks of the ejectiles for the $pp \rightarrow ppK^+K^-$ reaction at a beam momentum of $3.302 \text{ GeV}/c$.

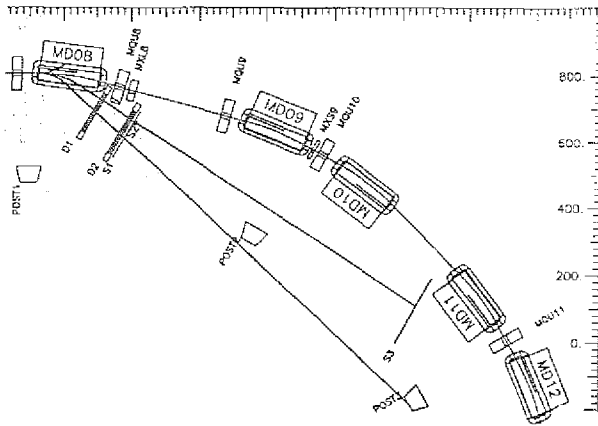


Figure 1: Setup of the COSY-11 experiment.

The target will be installed in front of COSY-dipole no. 8 which is used as a magnetic separator for the ejectiles and allows 0° measurements. The various detector components are drift chambers D1 and D2, scintillation counters S1, S2 and S3 and a silicon strip detector Si. The position and the dimensions of the detectors result from tracking studies of the ejectiles at maximum possible momentum.

The target will be a hydrogen cluster target with a density of $\sim 10^{14} \text{ atoms/cm}^3$ and a variable dimension [1].

The essential analysis of the various meson production experiments will be a missing mass determination by measuring the two proton momenta accurately. This will be done by the two drift chambers with an active area of $40 \text{ cm height} \times 130 \text{ cm width}$. Each of them consists of 6 layers, 2 with vertical and 4 with inclined ($+31^\circ$ and -31°) wires [2]. The position resolution will be $\sim 200 \mu\text{m rms}$ in x - and $\sim 500 \mu\text{m rms}$ in y -direction. The resulting invariant mass resolution for a $1 \text{ GeV}/c^2$ meson calculated from the reconstructed proton momenta is 0.2 MeV rms for a target dimension of 3 mm .

The trigger signal for our experiment comes from the scintillation counters. An event is accepted if two protons are seen in the scintillators. S1 will cover the whole drift chamber by 16 scintillation

counters of 10 cm width . The region where the protons pass the scintillator for reactions with a low relative momentum will be covered by S2, 16 scintillation counters with a width of 1 cm . All scintillator strips will have a lengths of 45 cm with photomultiplier readout on both sides. For the third scintillator S3 a scintillator wall of 2.2 m width and 1 m height , similar to the Amadeus detector [3], is planned. The main feature of this wall is the time of flight measurement of the protons for a further momentum determination. In addition it will be very helpful as a trigger element and for particle identification. The output of S3 will be the time, the position of a particle hit and its energy loss.

The silicon strip detector is installed to detect negatively charged particles especially kaons for the background reduction in a selective analysis of the charged kaon decay channel of the produced strange mesons. The detector consists of silicon pads with the dimension $19 \text{ mm} \times 24 \text{ mm}$ divided into 4 strips of 4.6 mm width . They will be installed in the dipole gap in three overlapping rows with a total height of about 7 cm over a length of $\sim 1.4 \text{ m}$. To reduce the costs of this detector system the readout of the Si-pads is performed by AMPLEX chips and a DRAMS module which were developed at CERN [4]. AMPLEX is an integrated circuit with 16 preamplifier, filter amplifier and sample and hold stages, a multiplexer and an output driver. DRAMS is a CAMAC module which manages the readout, digitizing and storage of up to 1024 AMPLEX channels.

The vacuum chamber of our experiment will be a setup like the schematic drawing in fig. 2.

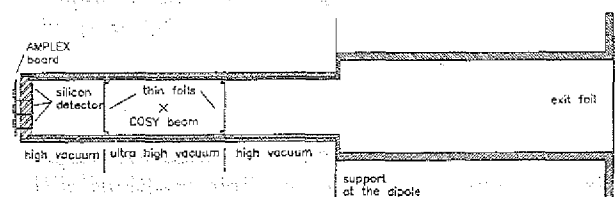


Figure 2: Schematic sectional drawing of the COSY-11 vacuum chamber.

The idea is to separate the ultra high vacuum in the COSY-ring ($p \cdot 10^{-10} \text{ mbar}$) from only high vacuum regions ($p > 10^{-6} \text{ mbar}$) with thin foils. Thus we avoid the UHV compatibility for the detector installation and the exit foil. The material of the foils, polyamide or metal, is still under

discussion. Only the si-pads will be mounted in the vacuum chamber whereas the electronic parts are located outside because of the heat production of the chips. For this reason we will use a plate of plastic material with electric contact pins which will be mounted on the 1.4 m long window on the left side of the chamber. On the right side the exit foil should cover the whole length of the chamber. To provide a sufficient stability a suspension at the dipole and a further massive suspension outside of the dipole is discussed.

Apart from the work on the detector setup extensive Monte Carlo calculations concerning the detector efficiencies and the trigger rates have been done. In fig. 3 the detection efficiency for the ejectiles of the $pp \rightarrow pp\phi$ at a beam momentum of 3.416 GeV/c is shown.

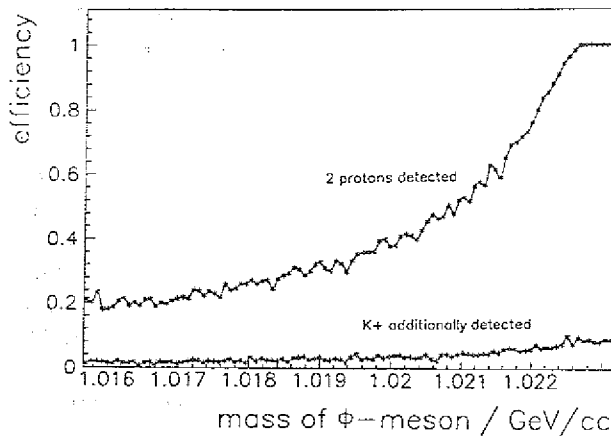


Figure 3: Geometrical efficiency for the $pp \rightarrow pp\phi$ reaction.

The upper curve shows the efficiency for the detection of the two protons in the drift chambers as a function of the invariant mass calculated from the momentum components of the two protons. In a range of 0.5 MeV/c² the efficiency is 1 and drops to 0.2 at 7 MeV/c² below the maximum invariant mass. Therefore the setup is an excellent 4π detector for narrow meson resonances if the, for the physics most interesting, threshold production will be considered. If the detection of a K^+ is required in addition the efficiency will drop as shown by the lower curve in fig. 3.

To check the background and the requirements on our data acquisition system we looked at the sensitivity of our setup to all possible reactions with a sufficient high cross section. As an example in table 1 the trigger rates for the $pp \rightarrow pp\phi$ experiment at a beam momentum of 3.422 GeV/c are summarized. A trigger is defined here by the detection of two protons in the scintillation detectors.

For the ϕ -production we get 0.26 events/s if the cross section is 1 μ barn. The trigger rate for all considered background reactions is expected to be as low as 13.5 events/s.

Table 1: Trigger probabilities, cross sections and trigger rates for all possible reaction channels in the case of the $pp \rightarrow pp\phi$ experiment at a beam momentum of 3.422 GeV/c. The cross sections are from [5] and the trigger rates were calculated with a luminosity of $10^{30} \text{ cm}^{-2}\text{s}^{-1}$.

reaction channel	trigger prob. [%]	cross section [mb]	trigger rate [1/s]
$pp \rightarrow pp\phi$	28.01	?	-
$pp \rightarrow ppK^+K^-$	6.46	?	-
$pp \rightarrow ppK_L^0 K_S^0$	8.43	0.003	0.25
$pp \rightarrow pp\pi^0$	0.02	2.90	0.64
$pp \rightarrow pp\pi^0\pi^0$	0.11	0.92	0.99
$pp \rightarrow pp\pi^+\pi^-$	0.11	2.67	2.96
$pp \rightarrow pp\pi^+\pi^-\pi^0$	0.26	0.74	1.95
$pp \rightarrow pp$	<0.001	15.00	<0.15
$pp \rightarrow pn\pi^+$	0.02	11.44	2.52
$pp \rightarrow pn\pi^+\pi^0$	0.04	4.07	1.79
$pp \rightarrow pn\pi^+\pi^0\pi^0$	0.09	1.80	1.58
$pp \rightarrow pn\pi^+\pi^+\pi^-$	0.08	1.15	0.95

The status of the different COSY-11 components is given in the following list:

- All mechanical parts of the cluster target are ready and first test measurements of single components were performed. A test setup to produce a hydrogen cluster beam is now installed at the IKP in Münster.
 - A first prototype of the drift chamber is built and will be tested next. The production of all drift chambers will start during 1992 in Cracow.
 - For test measurements of the scintillator S3, a small representative unit is being built.
 - A prototype of the Si-detector is under construction at the ZEL to test the detector specification and the readout.
 - The vacuum chamber will be constructed at the ZAT. At the end of 1991 first design studies have started.
- a ZEL, KFA Jülich
b IKP, University of Münster
c IKP, Jagellonian University of Cracow, Poland
d Institute of nuclear Physics, Cracow, Poland

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Reconstruction of the Shape of Meson Resonances Produced in the $pp \rightarrow ppM$ Reaction Near Threshold

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The aim of the COSY-11 experiment is to study the structure of resonances in the meson production reaction $p+p \rightarrow p+p+M$ in the threshold range. The four momentum vector of M should be determined in a missing mass method by analysing the proton momenta spectra in a COSY dipole magnet¹.

The mass determination depends on reconstruction of particle tracks taking into account the resolution of the detector system. In a simulation program it was assumed that a proton beam hits a parallelepiped target and the outgoing protons are momentum analysed in the COSY dipole magnet with a field distribution approximated by a box field (Fig. 1). They are measured in a detector system consisting of two packets of six drift chambers each displaced by a distance of 60 cm.

corresponding to these resonances, sampled according to the phase space distribution, were transported through the magnet; their track parameters were smeared out according to the above detector performance parameters and were retraced back to the target through the magnetic field. From the proton momenta found in this way the invariant mass was calculated for each event and a reconstructed mass spectrum was determined. The quality of the reconstructed mass spectrum (see Fig. 2 for an example) is described by the half widths $\Gamma_1, \Gamma_2, \Gamma_3$ and their valley-to-peak ratios r_1, r_2 .

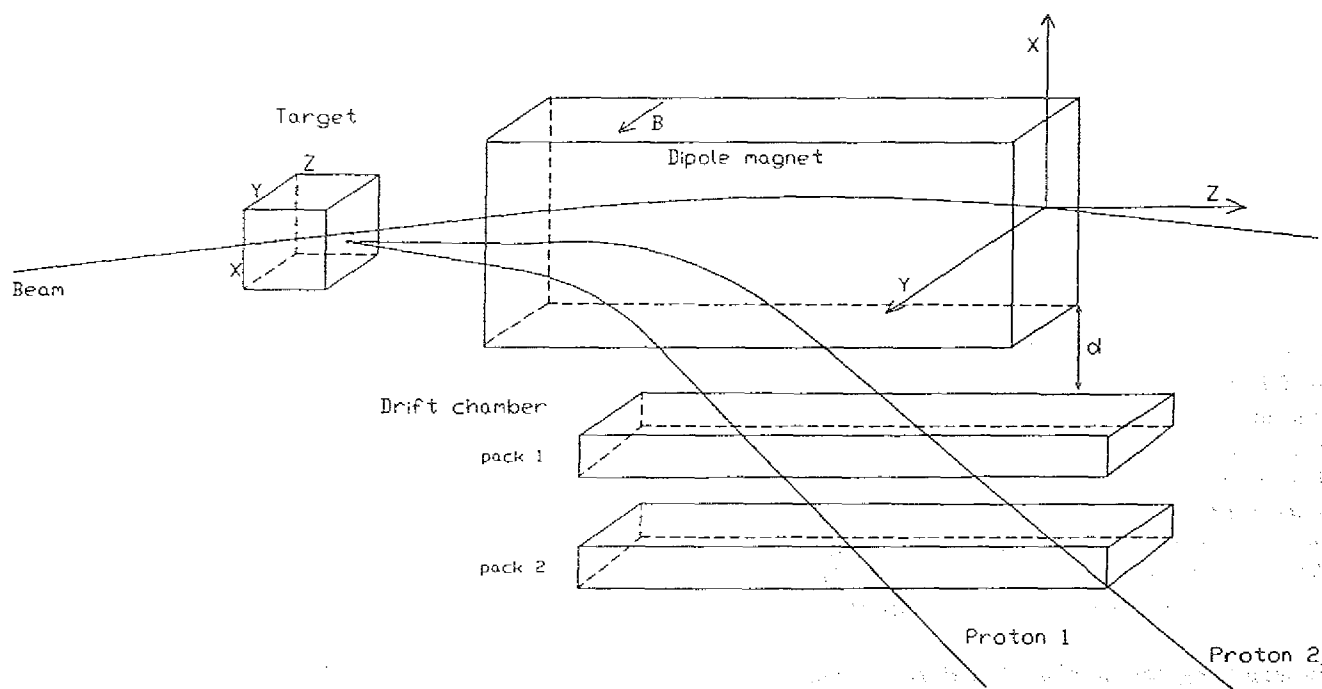


Figure 1: Experimental lay-out

The simulation of the detector system performance² led to the following inaccuracies of the reconstructed track parameters:

standard deviations:

of track position $\sigma_x = \sigma_y = 0.2 \text{ mm}$
of track angle $\sigma_\alpha = \sigma_\beta = 0.2 \text{ mrad}$

The spectrum of invariant mass M of the resonances was assumed in form of three very narrow peaks separated by 1 MeV (987.5, 988.5, 989.5 MeV) with a width of 0.02 MeV each. The outgoing protons

The following dependences of reconstruction accuracy were investigated:

1. on the edge length of a cube target; parameters given in table 1 indicate that up to the edge length of 5 mm the quality of reconstruction is acceptable;
2. on the shape of the parallelepiped target with dimensions indicated in Fig. 1; as can be seen from the results listed in table 2 only the target width influences significantly the quality of reconstruction while the target height and depth in the beam direction are practically without influence;

Target Dimensions [x, y, z] (width,height,depth)	Γ_1 [MeV]	Γ_2 [MeV]	Γ_3 [MeV]	r_1	r_2
[1, 1, 1]	0.22	0.26	0.17	0.00	0.00
[3, 3, 3]	0.31	0.30	0.24	0.06	0.01
[5, 5, 5]	0.52	0.33	0.33	0.33	0.12
[10,10,10]	0.57	0.62	0.60	0.67	0.58

Table I: Reconstruction quality parameters for different sizes of the cubic target.

Target Dimensions [x, y, z] (width,height,depth)	Γ_1 [MeV]	Γ_2 [MeV]	Γ_3 [MeV]	r_1	r_2
[5, 5, 5]	0.53	0.33	0.33	0.33	0.12
[5, 5, 1]	0.40	0.30	0.31	0.32	0.17
[5, 1, 5]	0.45	0.40	0.30	0.38	0.13
[1, 5, 5]	0.23	0.25	0.25	0.00	0.00

Table II: Reconstruction quality parameters for different parallelepiped targets.

3. on the performance parameters of the detection system; the variations of the parameters σ_x , σ_y and σ_α and σ_β revealed a rather small sensitivity of reconstruction accuracy to the angle parameters up to 0.5 mrad and to the position parameters up to 0.5 mm what puts the limits on additional inaccuracies due to the straggling effects.

4. on inaccuracy of beam momentum Δp_b , of the field strength ΔB and on the distance between magnet edge and detector system d ; the performed calculations showed that up to $\Delta p_b/p_b = 10^{-4}$, $\Delta B/B = 10^{-3}$ and $d = 0$ cm, 10 cm or 20 cm the results of reconstruction are practically the same.

Summarizing the results of performed calculations it could be stated that the target width limits the experiment resolution, the straggling up to 0.5 mrad does not influence the accuracy and the beam momentum and field strength inaccuracies inside of limits set by COSY requirements are acceptable for mass spectrum reconstruction.

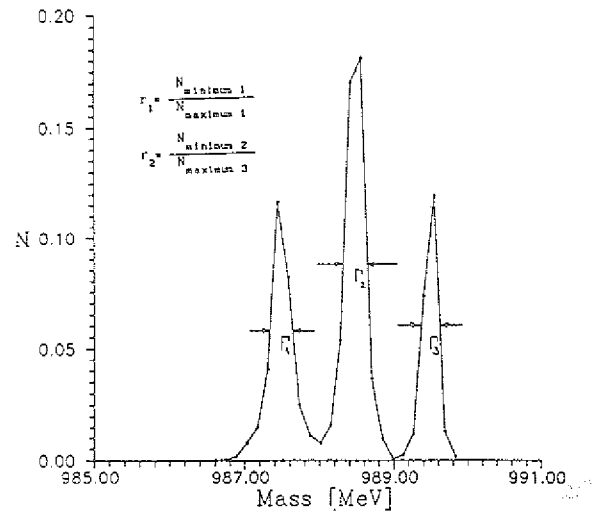


Figure 2: Reconstruction of invariant mass spectrum.

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A Fastbus System for Preparation of the Experiment COSY-11

M. Ziolkowski, W. Oelert

The experimental data which will be produced by the COSY-11 detector will be collected by the Fastbus front end electronics as: multihit TDCs for drift chamber, fast TDCs and ADCs for scintillator hodoscope.

A small Fastbus laboratory system has been set up for use in such applications in order to:

- learn and train software and hardware aspects of Fastbus,
- test Fastbus modules and develop electronics,
- perform data acquisition in the laboratory for detector development.

The system consists of a Fastbus crate controlled by CERN Fastbus Interface CFI which connects to host-dedicated I/O camac register STR 301, see fig. 1.

The microcomputer camac crate controller DSP 6002 provides a link to the IBM-PC computer which executes programs of the user and acts as an inexpensive peripheral server (i.e. hard disc, optical disc, graphic terminal, printer).

The CFI M68k microprocessor controls the execution of standard Fastbus functions implemented on board as microcode routines.

Fortran written programs running on the PC computer call the adopted CERN Standard Fastbus Library [1]. Execution of Fastbus compound operation lists (which are preloaded in the CFI master) may be triggered with an external signal. Readout data may optionally be stored in CFI buffer and then transmitted to the host using the DMA mode.

This simple Fastbus test system has been proven to be a user friendly tool kit in the laboratory with Fastbus instrumentation.

Reference:

1. Users Guide to CERN Fastbus, Routines DD Division, CERN 1988.

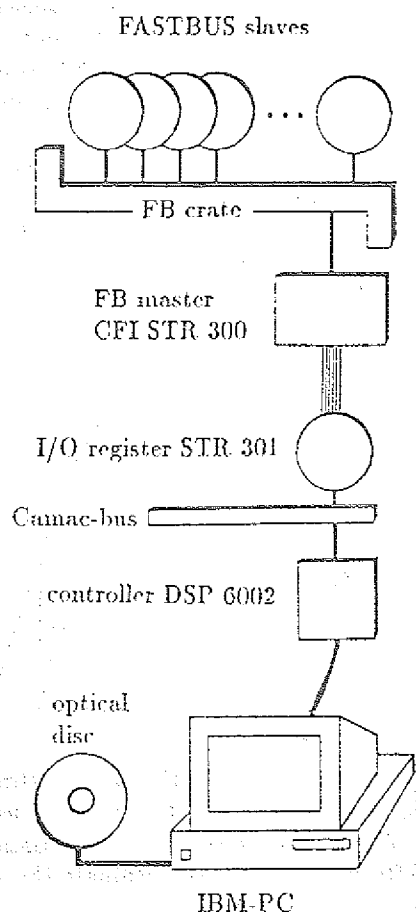
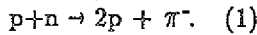


Figure 1: The architecture of the small Fastbus data acquisition system.

The Experiment COSY 10 (Pionic Atoms): A Status Report

P. Cloth, V. Drücke, A. Djalois, D. Filges, D. Grzonka, K. Kilian, H. Machner, H.P. Morsch, D. Protic, G. Riepe, E. Roderburg, P. von Rossen, P. Turek, L. Jarczyk¹, M. Kistryn¹, J. Smyrski¹, A. Strzalkowski¹, W. Drochner^{2,4}, W. Meiling², A. Budzanowski³, D. Dabrowski³, I. Skwirczynska³, W. Erven⁴, J. Holzer⁴, H. Kopp⁴, K. Zwoil⁴, H.S. Plendl⁵, B.J. Lieb⁶, J. Konijn⁷, D. Frekers⁸

The aim of the experiment is the investigation of pionic atoms in low lying states. Because of the strong interaction between captured π^- s and nucleons, which is the quantity of interest, the pions annihilate at a substantial rate before reaching the lowest atomic states. This makes the study of transitions to such states with the help of the classical method, namely to capture degraded π^- s into atomic states, unfeasible. Alternative methods proposed are based on the idea of π^- -production directly in the states of interest. Especially convincing seems the idea of recoilless π^- -production, since this method maximizes the cross section. We will employ the elementary reaction



This process embedded in a nuclear environment is shown in Fig. 1.

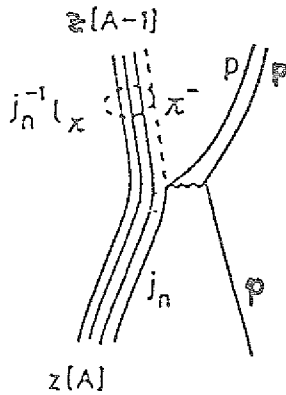


Figure 1: Graph for the reaction leading to a pionic atom.

In principle, the experiment could be made with incident neutrons leading to a bound state for the two particles to be observed in the final state: a deuteron. The advantage of this reaction would be that it is rather easy to measure the four momentum vector of the emerging deuteron. The disadvantage is the low neutron intensity, especially when the four momentum vector of the neutron has to be measured with similar accuracy. A first attempt employing this reaction failed [1]. We, therefore, have chosen protons in the entrance channel. The cross section for this reaction on an atomic nucleus is given by

$$\frac{d\sigma}{d\Omega} = \frac{d\sigma}{d\Omega}(p+n \rightarrow 2p + \pi^-) \times \left| \int \phi_{n,l}^*(\vec{r}) \exp[-i(\vec{k}_{2p} - \vec{k}_n) \cdot \vec{r}] \phi_{\alpha}(\vec{r}) d^3r \right|^2 \quad (2)$$

Here the bound neutron has been assumed to be in a quantum state ϕ_{α} with quantum numbers α while $\phi_{n,l}$ denotes the pionic wave function. It is

selfevident that the spectroscopic factor has a maximum for a momentum transfer $\vec{q} = \vec{k}_{2p} - \vec{k}_n = 0$. This method was successfully applied in the study of hypernuclei.

The experiment will measure the direction and energy of two protons identified as such. Because of the Fermi-momentum of the struck neutron and the well known repulsive $1s$ interaction in the p - p channel a maximum angle of 137 mr between the two probes is expected. The anticipated detector will achieve a sufficient missing mass resolution by a spatial precision of ~ 2 mr and an energy resolution $\Delta E/E \approx 10^{-3}$. The detector is a hybrid: a magnet spectrometer and a solid state detector.

The magnet spectrometer BIG KARL will detect protons emerging into a cone with ± 28 mr opening. In order to allow for 3 dimensional track reconstruction the spectrometer was redesigned into a Q3D2 structure. Three high pole-tip-field quadrupoles (on loan from DESY) provide the needed focussing. In Fig. 2 we show the structure of envelopes for this "angle-sensitive" operation mode. Calculations performed with the program TURTLE indicate an "optical amplification" of two in the y -direction, i.e. any uncertainty in a trajectory measurement in this direction doubles when calculating back the trajectory from the target to the spectrometer entrance.

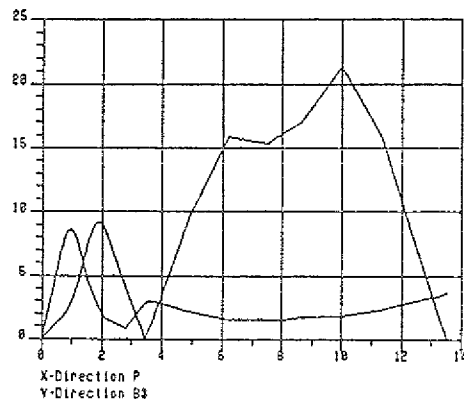


Figure 2: Envelopes for the beam through the magnetic Q3D2-spectrometer in x - and y -direction. The abscissa gives the flight path through BIG KARL in m and on the ordinate the values are given in cm.

The detector in the focal plane will be a stack of two times six multiwire-drift-chambers (MWDC) followed by a trigger hodoscope. The two planes of the MWDC will allow to measure a track with position resolution of 0.1 mm and 0.25 mm in x - and y -direction, respectively. Readout electronics is in common with experiment COSY 11. Time digitizing electronic bought in this year (together

with experiments COSY 1, 2 and 6) will allow for double hit resolution.

The trigger hodoscope will consist of organic scintillator pads of 20 cm times 40 cm each. There will be again two planes approximately 1 m apart with time of flight measurement to reduce background. Each plane will consist of two layers of pads shifted by one half width with respect to each other. By this way a coarse spatial resolution of 10 cm will be achieved.

Around the acceptance of the magnet spectrometer charged particles will be detected by the GERMANIUM WALL. In Fig. 3 the cross section through this device is shown.

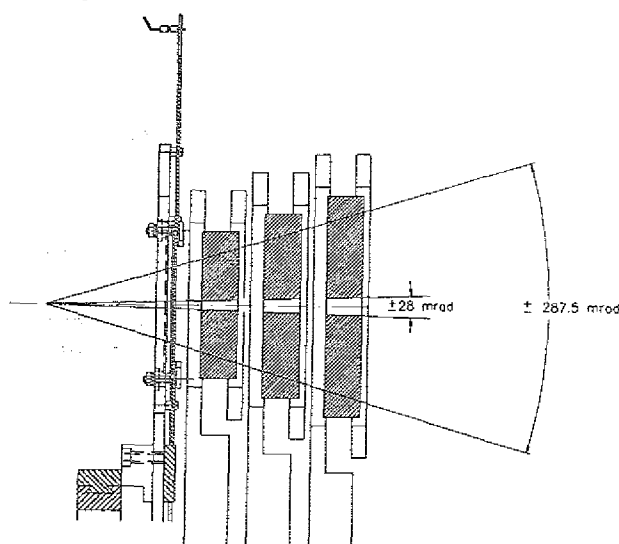


Figure 3: Cross section through the Ge-wall.

It consists of four Ge-diodes, each having a hole in its middle for reaction products going to BIG KARL and the primary beam as well. The total thickness of 52 mm allows the stopping of 1850 MeV/c protons. The three thick calorimeter diodes (each 17 mm thick) will have a structure of 32 pies on their front face. The advantages of this are beyond multi-hit detection smaller capacitances of each subdiode and therefore reduced noise and in addition higher counting rates possible in the following analog amplification circuits compared to only one diode. The first 1 mm thick diode will give a ΔE -signal for particle identification. To allow for the measurement of the emission angles (ϑ, φ) this diode will cover structures in its front and back side as well. After the encouraging results with 2 times 50 stripes (see separate contribution to this report [2]) we decided to use a rotational invariant structure like the underlying physics. Such a structure will also lead to a uniform rate distribution over the stripes. Therefore, the structure will be 200 left bended Archimedes' spirals on the front and 200 right bended on the back side. The grooves between the spirals will be 25 μm wide. This diode is shown in Fig. 4 but with only 50 stripes on each side for clarity. By the crossing points between the spirals, one of which is shown in the figure) 20.000 pixels were defined. The Ge-wall defines a solid angle of 23.6 msr in the laboratory

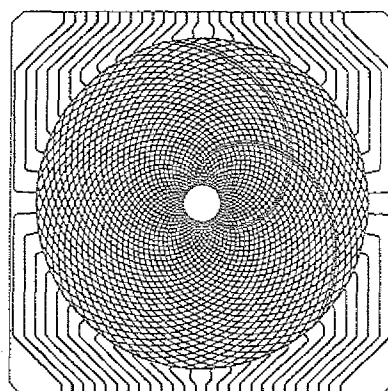


Figure 4: Simplified view of the first Ge-diode with shown 50 (out of 200) spirals on each side.

system with a spatial resolution of $\pm 1.2 \mu\text{sr}$. The polar angle will be measured to $\pm 1.4 \text{ mr}$. These numbers are all for the geometry shown in Fig. 3.

Data digitising will be performed completely by CAMAC electronics making use of the FERA-bus. To achieve sufficient energy resolution Silena ADC's were chosen while time digitising will be made with Le Croy FERA modules. In addition a FERA compatible hit-pattern presently under design will identify the pixels fired. In order to derandomize the data and take different conversion times into account a dedicated CAMAC module was designed much in the spirit of the earlier used MSU [3].

Data readout will be through a VIC bus to a multi-processor VME-system employing a VIC-CAMAC controller developed by us in Jülich [3]. This system will store the data on tape and will allow on-line sorting on a work-station for experiment control. For this task standard CERN software including PAW was successfully tested. Further details are found in a dedicated contribution in this report [5].

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A Study of the $p+d \rightarrow \eta+{}^3\text{He}$ Reaction with a 4π Detector Geometry

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There are four light mesons with quantum numbers 0 for charge, strangeness and isospin. Two of them are pseudoscalar mesons belonging to the $1^P = 0^-$ nonet: η and η' and two are their heavier vector equivalents belonging to the 1^- nonet: ω and ϕ . These mesons differ from the closely related π^0 and ρ^0 , respectively by having an ss component in their wave function. A part of the COSY 10 detector seems to be ideally suited to measure excitation functions for the production of the above mentioned four neutral mesons (see Fig. 1).

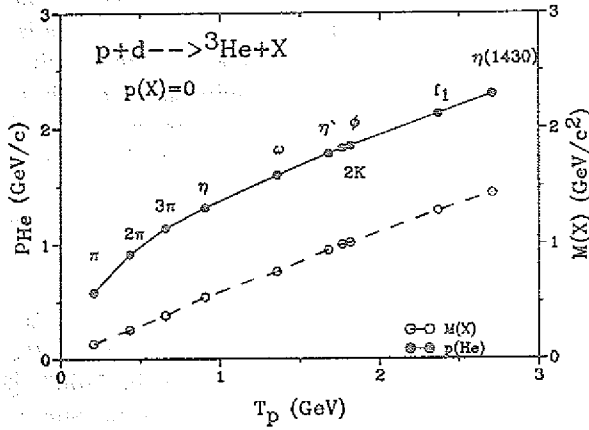


Figure 1: Upper curve and left axis: The maximum of the ${}^3\text{He}$ -momentum for different mesons under the assumption of recoil free production; lower curve and right axis: The masses of the indicated mesons as function of the proton kinetic energy.

All particles can be produced at COSY and the maximum momentum of the ${}^3\text{He}$ -particles can be measured. Of special interest is the lightest of these mesons, the η which has the recently attracted much attention due to the unexpected large cross section in the reaction under consideration close to threshold [1]. Three different reaction mechanisms can in principle contribute:

- The η is produced in a one step process with the nucleons being lumped into the phase space of a ${}^3\text{He}$.
- The η is produced in a two step process with an intermediate π^+ [2].
- A virtual π excites a nucleon resonance which then decays into η and nucleon.

In all cases the three nucleons have to fuse. In the last case only the $N^*(1535)$ contributes to this possibility close to threshold. We, therefore, propose to measure the excitation function for the η -meson from threshold up to 1850 MeV/c proton beam

momentum. In this momentum range the magnetic spectrometer plus the Ge-wall of COSY 10 are a 4π detector in the cm-system (see Fig. 2).

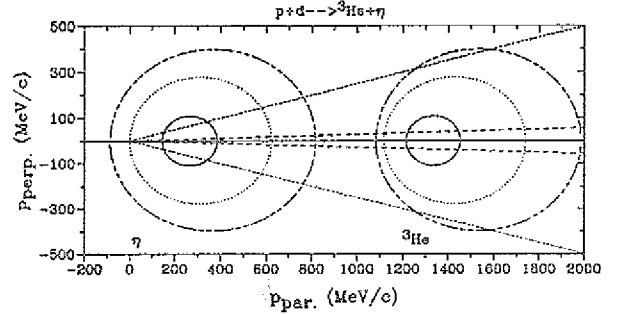


Figure 2: Kinematical loci in a polar diagram of the η 's and ${}^3\text{He}$ -ions for different bombarding momenta: 1600 MeV/c (solid curves), 1740 MeV/c (dotted curves), and 1900 MeV/c (long dashed, 2x short dashed). Also shown are the angular openings of the magnetic spectrometer (long dashed) and the Ge-wall (short dashed).

First Monte Carlo simulations are shown in Figs. 3 and 4 for two different bombarding energies. A "background" due to 2π production is also shown. From these simulations we expect to good signal to background ratio.

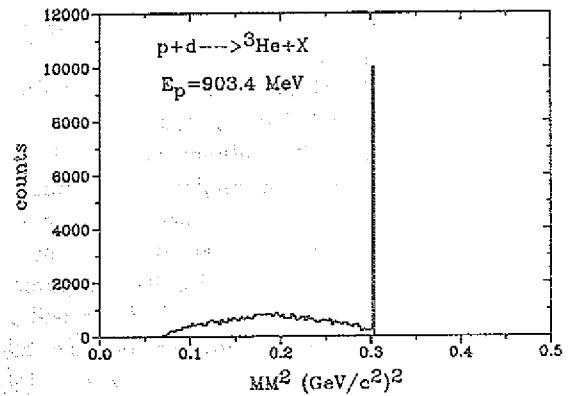


Figure 3: Monte Carlo simulation for η and 2π production for the indicated energy.

Essential for these results are the good beam quality and a rather thin liquid deuterium target. Such a target is presently under construction in Jülich [3].

With 2 mm thick target we expect to produce and measure approximately 0.5 η 's per second for a bombarding energy close to threshold. For higher energies a drop of the cross section by one order of magnitude will still lead to a reasonable count rate.

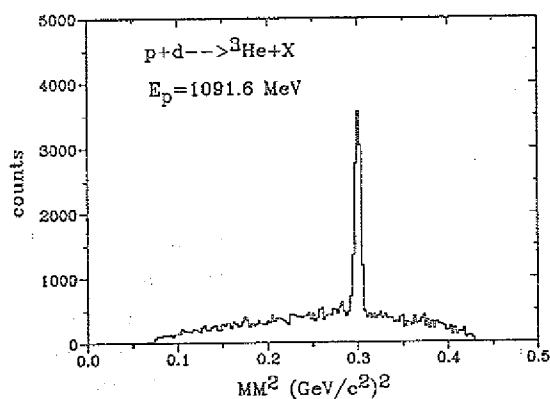


Figure 4: Same as Fig. 3, but for a different bombarding energy.

As soon as a polarized proton beam will be available, we will employ it to measure, in addition, the total cross sections and all angular distributions and the analysing power. With these informations a better understanding of the reaction mechanism should be possible.

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Monte Carlo Simulations of Energy Dose Profiles of Protons in Water for Studies of Possible Medical Applications of the COSY Beam

V. Drücke, P. Cloth, D. Filges, U. Linz, R. Maier

In connection with recent ideas to use the COSY proton beam for medical applications energy dose calculations have been conducted as an introduction to a more comprehensive dosimetry program, which shall provide the physical base for treatment protocol development.

To develop the application potential of protons for medical treatment, it is necessary to use particle transport simulation codes, which can - besides dE/dx -energy loss - consider nuclear interaction with secondary particle production. This is due to proton energy requirements upto 200 MeV. It is also desirable to compare proton dose effects with that of other particles in use today as photons, electrons, neutrons and pions.

The most suitable Monte Carlo system for this purpose seems to be the HERMES system (High Energy Radiation Monte Carlo Elaborate System) /1/, which can treat the transport and interaction mechanism of all these particles. Therefore, the computational work has to be shared between several Monte Carlo codes, each solving a specific part of the total problem. The HERMES system and its particle and detector data flow is illustrated in Fig. 1.

Currently three Monte Carlo codes are implemented, namely HETC, MORSE-CG, and EGS4, besides some auxiliary codes. HETC (High Energy Transport Code) is implemented in a version based on the original coding. The particles treated by HETC are p , n , π^+ , π^- , D , T , ^3He , α , and light heavy ions with $A \leq 20$. The physical model of HETC consists of mainly two parts, the intranuclear cascade model and the evaporation model, performing

- charged particle ionization energy loss (dE/dx),
- primary particle small angle multiple Coulomb scattering
- elastic collisions for proton and neutron nucleus collisions
- production of residual nuclei and residual excitation energy.

MORSE-CG is a multigroup particle transport program to simulate the histories of neutral particles, which may originate from neutron or gamma sources. MORSE-CG simulates the following physical process

- elastic neutron nucleus collisions including hydrogen collisions,
- nonelastic neutron nucleus collisions,
- capture and fission, gamma production and transport, productive nuclear collisions ($n, 2n$ etc.)

The EGS4 computer code provide a Monte Carlo simulation of the coupled transport of electrons and photons with energies from a few keV upto several TeV. The particles treated by EGS4 are e^+ , e^- and γ , which may originate from fixed sources of these particles or from electromagnetic particle sources

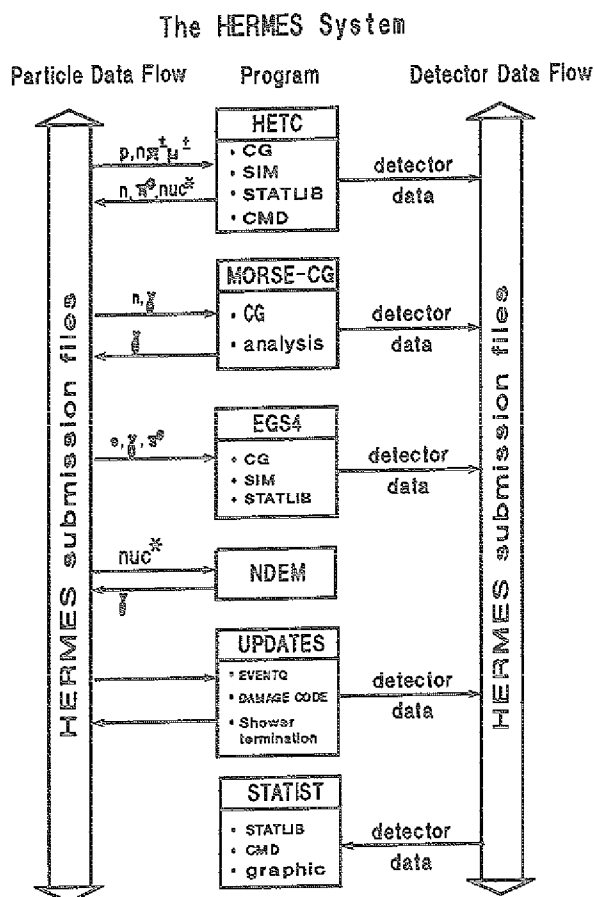


Fig. 1: Organization of the HERMES program system for simulation studies on beam materials interaction

(e^+ , e^- , γ 's in phase space). EGS4 simulates the following physical processes:

- small and wide angle scattering
- bremsstrahlung, pair production, Compton scattering, Moeller effect, Bhabha effect, photoelectric effect
- annihilation
- continuous e^+ , e^- energy loss

As an example the result of an irradiation protocol to gain a proton induced energy dose plateau in water is shown in Fig. 2.

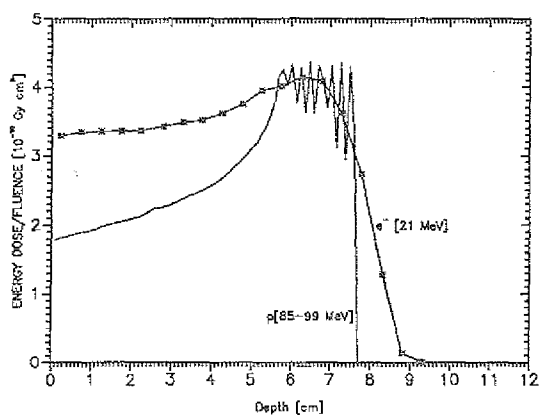


Fig. 2: Energy dose distribution in water of protons and electrons

The proton induced energy dose plateau was simulated by modulation of beam energy between 85 and 99 MeV and intensity. For comparison to the dose effects of protons the dose profile of an electron beam of 21 MeV is shown as calculated by the EGS4 code. The electron beam of 21 MeV is the standard maximum energy of the current electron irradiation facilities in therapy.

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Spin Correlation in $\bar{\Lambda}\Lambda$ -Production

K. Kilian, W. Oelert, K. Röhrich, Th. Sefzick, M. Ziolkowski
for the PS185 collaboration

Due to the rather simple structure of the Λ -particle its spin dynamics can be related to the strange quark content. Quark spin correlations of the $\bar{s}s$ -pair are imaged by the spin correlations of the $\bar{\Lambda}\Lambda$ -pair.

The spin-vectors of the two hyperons produced in the interreaction $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ can be correlated to each other. A matrix with nine spin-correlation-elements [1] describes this property of the produced particle-pair. The parity conservation of the strong interaction implies that the elements C_{xy} , C_{yx} , C_{yz} , C_{zy} must be equal to zero. Additionally the elements C_{xz} , C_{zx} are equal because of charge conjugation symmetry. The spin correlation parameters C_{ij} appear in the complete description of the differential cross section angular correlations in the distribution of the decay baryons [1]:

$$W(\bar{X}, \bar{Y}, \bar{Z}, X, Y, Z) \sim [(1 + \bar{\alpha}P_{\bar{Y}}\bar{Y} + \alpha P_Y Y) + (C_{xx}\bar{X}X + C_{yy}\bar{Y}Y + C_{zz}\bar{Z}Z + C_{xz}\bar{X}Z + C_{zx}\bar{Z}X)] \quad (1)$$

where:

$$(\bar{X}, \bar{Y}, \bar{Z})(X, Y, Z) \quad (2)$$

are the direction cosines of the decay baryons in the $\bar{\Lambda}$ and Λ rest frames, $P_Y, P_{\bar{Y}}$ is the polarisation for $\bar{\Lambda}, \Lambda$, respectively. $d\sigma/d\Omega$, P , and C_{ij} are all functions of the $\bar{\Lambda}\Lambda$ scattering angle. The analysis of Monte-Carlo events indicates significant inefficiencies, see Fig. 1a, due to detector constraints and due to limitations in the evaluation code.

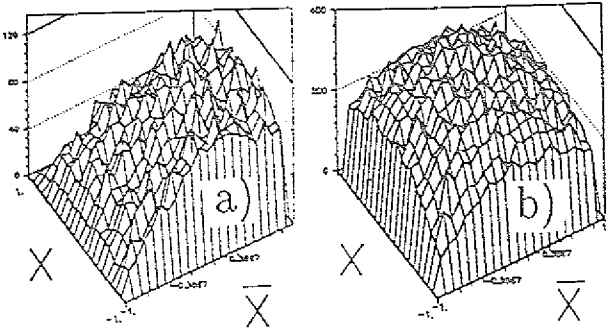


Figure 1: Distribution of Monte-Carlo events in $(\bar{X}, X)$ space for $\bar{\Lambda}$ -production angle around 90° ($\bar{\Lambda}\Lambda$ rest frame) a) without correction and b) after symmetrisation.

Among others a C_{xx} pseudo correlation for Monte-Carlo events, shown on Fig. 2a, is observed. An appropriate correction method is necessary to obtain the true spin correlation parameters. A method has been developed based on the used symmetries in the process $\bar{p}p \rightarrow \bar{\Lambda}\Lambda \rightarrow \bar{p}\pi^+\pi^-\pi^-$: for each event (Eq. 2) there are three others (Eq. 3), which give exactly the same results. This property results from Eq. 1 and is a consequence of P, C and CP symmetry.

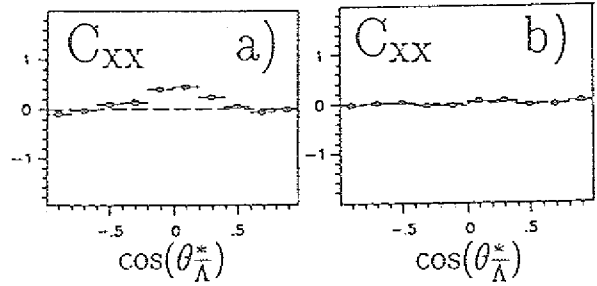


Figure 2: C_{xx} parameter as a function of cosine of the $\bar{\Lambda}$ -production angle for Monte-Carlo data a) before correction, b) after correction.

$$\begin{aligned} &(-\bar{X}, \bar{Y}, -\bar{Z}) \quad , \quad (-X, Y, -Z) \\ &(X, -Y, Z) \quad , \quad (\bar{X}, -\bar{Y}, \bar{Z}) \\ &(-X, -Y, -Z) \quad , \quad (-\bar{X}, -\bar{Y}, -\bar{Z}) \end{aligned} \quad (3)$$

Adding for each event (Eq. 2) these three other symmetric combinations (Eq. 3) an effective decrease and smoothing of the reported inefficiency (Fig. 1a) is observed (Fig. 1b). Corrected with this new efficiency data result in much smoother C_{ij} -parameter distributions, see Fig. 2b.

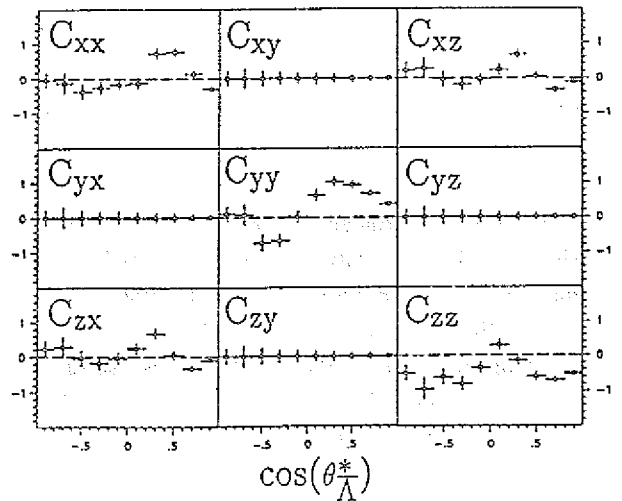


Figure 3: C_{ij} matrix for corrected experimental data (1919 MeV/c).

Figure 3 presents the appropriate corrected matrix for 34000 experimental events (1.919 MeV/c). Further details will be described in [2]. It is shown that the described method works well to extract significant spin correlations in the $\bar{\Lambda}\Lambda$ -production experiment.

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Feasibility of CP-Invariance in $\bar{\Lambda}\Lambda$ Hyperon Decay

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The PS185 collaboration [1] at LEAR/CERN uses a threshold decay spectrometer to study the anti-hyperon-hyperon production. Besides other observables as total and differential cross section and spin correlations [1,2] the polarization has been analysed in the $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ reaction for the kinematically correlated $\bar{\Lambda}$'s and Λ 's independently, see Fig. 1. An apparent difference would signal a CP-violating effect in the weak decay parameter α , since CP-conservation requires zero for the relative quantity:

$$A = (\alpha_{\Lambda}^P \Lambda + \alpha_{\bar{\Lambda}}^P \bar{\Lambda}) / (\alpha_{\Lambda}^P \Lambda - \alpha_{\bar{\Lambda}}^P \bar{\Lambda}).$$

Combining data samples at momenta 1546 MeV/c, 1642 MeV/c and 1695 MeV/c [1,3], one obtains the current best value for A

$$\langle A \rangle = -0.013 \pm 0.029.$$

The error quoted is only the statistical one and the value obtained for the entire center of mass angular range.

The quantity A to be measured from the $\bar{\Lambda}\Lambda$ system has been theoretically determined to be in the order of 10^{-4} to 10^{-5} [4].

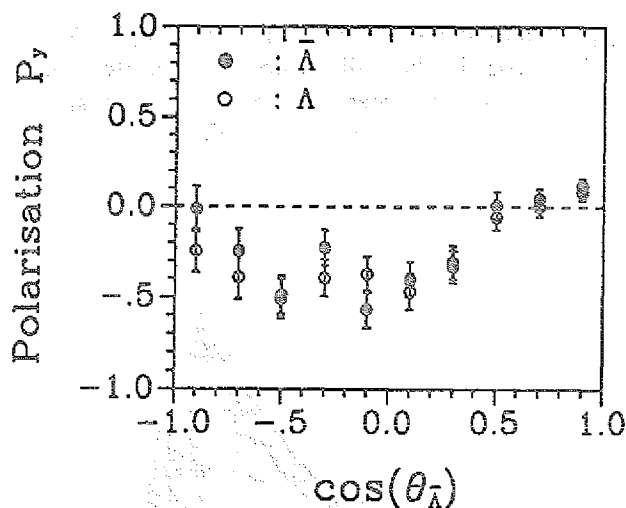


Figure 1: The polarisation of the $\bar{\Lambda}$ and Λ as a function of the cosine of its center of mass production angle.

Based on the result and experience obtained from the PS185 collaboration an international study group at CERN investigated the feasibility of a CP-invariance measurement. This working group [5] concluded: Such an experiment could be performed at an upgraded LEAR machine. The

chosen momentum would be right below the $\bar{\Lambda}\Sigma^0$ + c.c. threshold (1653 MeV/c) to avoid background due to the $\Sigma^0 \rightarrow \Lambda\gamma$ decay. With an internal hydrogen cluster target and an antiproton stack of about $2 \cdot 10^{11}$ revolving with 3.32 MHz at 1650 MeV/c beam momentum the peak luminosity is $6.6 \cdot 10^{31} \text{ cm}^{-2} \text{ sec}^{-1}$. Beam and target have to be unpolarized.

The aimed accuracy of $\leq 10^{-4}$ corresponds to about $5 \cdot 10^9$ analysed events. The principle of this experiment is based on some assumptions which have to be controlled for the systematic uncertainty: C, P and CP are 100 % conserved in the hadronic production process, beam and target have to be unpolarized and there must not be any selective depolarisation of $\bar{\Lambda}$ or Λ before their decay, possible detector asymmetries, different interactions of particles and antiparticles with matter and also the possible misidentification of pion and proton.

It is estimated that an experiment at an upgraded LEAR machine with an internal jet target can take several 10^9 $\bar{\Lambda}\Lambda$ events within one year of data taking, thus obtaining the statistics required for a determination of direct CP violation in a process outside the neutral K system with an accuracy better than 10^{-4} .

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In experiments where only charged particle tracks are measured one has to be aware that a non-negligible amount of reactions with additional uncharged particles can occur. In the case of the $\bar{p}p \rightarrow \bar{\Lambda}\Lambda \rightarrow \bar{p}\pi^+p\pi^-$ -reaction one gets additional $\bar{p}p \rightarrow \bar{\Lambda}\Lambda x^0$ -events with x^0 representing an uncharged particle like γ or π^0 . Measuring the 4-vectors $\mathbb{P}_{\bar{\Lambda}}$ and $\mathbb{P}_{\Lambda}$ by analyzing the charged delayed Λ -decays and taking into account the known 4-vectors of the incoming $\bar{p}$ and p there is sufficient information to determine the 4-vector of an unobserved neutral particle :

$$\mathbb{P}_{x^0} = \mathbb{P}_{\bar{p}} + \mathbb{P}_p - \mathbb{P}_{\bar{\Lambda}} - \mathbb{P}_{\Lambda} .$$

Missing mass calculations supply additional information about the event. The missing mass can be derived from the equation above easily :

$$M_{x^0}^2 = \mathbb{P}_{x^0}^2 = (\mathbb{P}_{\bar{p}} + \mathbb{P}_p - \mathbb{P}_{\bar{\Lambda}} - \mathbb{P}_{\Lambda})^2 .$$

Here $\mathbb{P}_{\Lambda}$ and $\mathbb{P}_{\bar{\Lambda}}$ are used. With only one of these 4-vectors the missing mass of the "opposite" particle system can be extracted

$$M_{(\Lambda x^0)}^2 = \mathbb{P}_{(\Lambda x^0)}^2 = (\mathbb{P}_{\bar{p}} + \mathbb{P}_p - \mathbb{P}_{\bar{\Lambda}})^2$$

supplying information about the Λx^0 -system.

For the two-body $\bar{\Lambda}\Lambda$ -reaction the "opposite" missing mass squared has to be $m_{\Lambda}^2 = 1.2446 \text{ (GeV/c}^2\text{)}^2$ and the missing mass squared has to be zero. Those events will be located in one point in Fig.1 and 2. For $\bar{p}p \rightarrow \bar{\Lambda}\Sigma^0 \rightarrow \bar{\Lambda}\Lambda\gamma \rightarrow \bar{p}\pi^+p\pi^-\gamma$ the directly produced Λ (or $\bar{\Lambda}$) will supply the correct missing mass for the opposite particle (Σ^0) but the other $\bar{\Lambda}$ (or Λ) will lead to some value greater then the Λ mass squared and restricted by the allowed missing mass region (Dalitz-plot) for $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\gamma$. So the $\bar{\Lambda}\Sigma^0$ -reaction will produce a cross-like (in the crossing point region there is no way to decide which Λ has been produced directly) structure growing out of the uniformly scattered background. The only way to distinguish between $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$, $\bar{p}p \rightarrow \bar{\Lambda}\Sigma^0 \rightarrow \bar{\Lambda}\Lambda\gamma$, and $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\gamma$ is to look at the "opposite" missing mass values because in all cases the missing mass is zero.

For $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\pi^0$ there is a good possibility to suppress $\bar{\Lambda}\Lambda\gamma$ -events by requiring the missing mass squared to be $m_{\pi^0}^2 = 0.0182 \text{ (GeV/c}^2\text{)}^2$. Again, because of three body kinematics, there is an allowed missing mass region for $\bar{\Lambda}\Lambda\pi^0$ which will be continuously populated.

It is obvious to see that missing mass calculations provide a good possibility to distinguish between several reactions showing the same signature by cutting on missing mass and "opposite" missing mass values. This strategy is the same for other reactions (like e.g. $pp \rightarrow pp\gamma$ and $pp \rightarrow pp\pi^0$) if the four momentum vectors of the charged particles are reconstructed but nothing is known about an additional neutral particle.

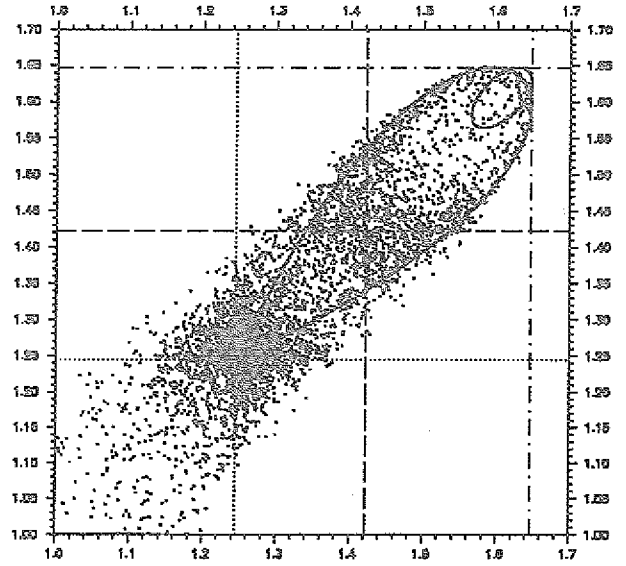


Fig.1 Dalitz-plots for $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\gamma$ and $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\pi^0$ (little ellipse) in the 2-dimensional "opposite" missing mass squared representation. The scales are in $(\text{GeV/c}^2)^2$. Data are cut around $m_{x^0}^2 = 0$. m_{Λ}^2 (dotted), $m_{x^0}^2$ (dashed), and the kinematical limit (dashed-dotted) are marked.

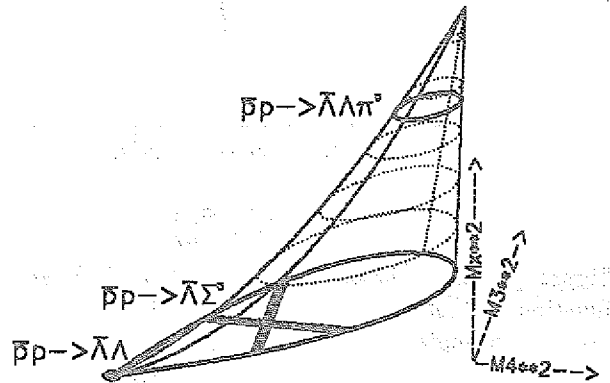


Fig.2 Physically allowed space for $\bar{p}p \rightarrow \bar{\Lambda}\Lambda x^0$ -events. At ground level ($m_{x^0}^2 = 0$) the Dalitz-plot from Fig.1 for $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\gamma$ is shown, at higher $m_{x^0}^2$ -values the Dalitz-plot for $\bar{p}p \rightarrow \bar{\Lambda}\Lambda\pi^0$ is shown.

A Trackfinding Procedure on the Basis of Event Rotation

K. Kilian, A. Kurtenbach, W. Oelert, K. Röhrich, E. Roderburg

The LEAR experiment PS185 at CERN uses sets of multiwire proportional and drift chambers for particle tracking. The chambers consist of 4 sets of wireplanes rotated by 45° relative to each other. We are developing a new, geometrical method to reconstruct these tracks from their hits in the detector.

The idea of this "Rotation Method" is the following: We regard the detector elements, which have been hit, of one of the 4 sets of planes, and project these on a reference plane. Normally we start by choosing the plane parallel to the chosen set orientation. Thus at any point on this reference plane the number of the projections onto this point is the number of hits aligned perpendicular to the reference plane.

By rotation of the chamber around an axis, which has the same orientation as the wires of the chosen set, the number of projections can be maximized at certain points on the reference plane. With no magnetic detector every track is straight, and therefore such a maximum indicates a particle track. Then the angle of rotation and the position of the maximum define a plane including the track. Repetition of this procedure with the other 3 sets of wireplanes allows the reconstruction of three dimensional particle tracks.

Thus the main effort to be made is a search of maximum projections in the reference planes. The maxima are two dimensional peaks because of the physical size of the projected detector elements. A track is found, if the maximum contains more than a defined number of hits. Not all of the detector elements can contribute to certain angles, because a track may leave the detector under large angles or just cut into a corner. Therefore, the distributions of projections have to be normalized to the number of possible projections.

Table 1 compares the results of the described method to the standard data analysis. Two tracks at negative track angles are easily found. The track at 0° is not found by our method.

Table 1: Comparison between rotation and standard analysis. For explanations see text.

rotation method		standard analysis	
direction ($^\circ$)	origin (mm)	direction ($^\circ$)	origin (mm)
15.84	.104		
15.44	.104		
14.63	.207	14.93	.254
		0.0	1.57
-32.50	-.593	-30.43	-.672
-34.16	.121	-33.57	.143

It turns out, that it is a fake and is rejected in the later analysis. The track at positive angles has three solutions in our method. This is due to the peak finding procedure. Certainly the finite resolution of the detector influences the results obtained. Thus we get several solutions with similar quality. A procedure to optimize this ambiguity is under development as well as the attempt to obtain more definite criteria of the maxima.

The geometrical rotation and projection of hits in an event is an alternative way of track finding, with the advantage of being independent of combinatorics and the detector type. Furthermore, since all detector planes are handled in the same way, the outlined procedure is the most unbiased way of analyzing the data. The method was evaluated on a basic level to prove that it working. The search for maxima has still to be optimized, since broad projections (hit clusters) lead to several solutions for the tracks. An investigation of the computing speed has not yet been done, since the algorithms are not fully developed until now.

Reference:

1. K. Kilian et al, Proposal P49, CERN/PSCC/81-69

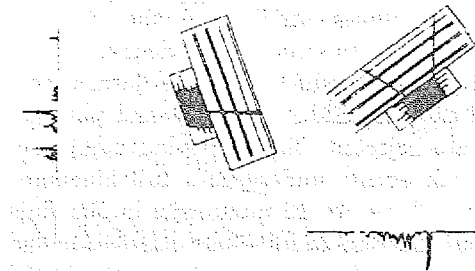


Figure 1: An event from PS185 in two orientations. The projections onto the reference plane are displayed on the left and bottom.

O. Steinkamp, K. Kilian, W. Oelert, K. Röhrich, B. Stugu
for the JETSET Collaboration

The experiment JETSET at CERN's Low Energy Antiproton Ring (LEAR) [1] searches for hadronic resonances in antiproton-proton annihilation, especially in gluon-rich channels like $\bar{p}p \rightarrow \phi\phi$. This channel is of special interest, because the simplest reaction mechanism has no connected quark-lines and thus is suppressed by the OZI rule. This suppression might however be overcome by an intermediate resonant gluonic state.

Previous experiments on $\phi\phi$ production in hadronic channels and in radiative J/Ψ -decay have shown structure in the whole mass region from threshold to $2.5 \text{ GeV}/c^2$. In particular an experiment at BNL, investigating the reaction $\pi^-p \rightarrow \phi\phi n$, claims to

have found three $J^{PC} = 2^{++}$ resonances which the authors interpret as tensor glueball states [2]. The JETSET experiment detects ϕ s in their decay channel into two charged kaons, which has a branching ratio of 49 % giving rise to the signature of four charged tracks for the process $\bar{p}p \rightarrow \phi\phi \rightarrow 4K$. An estimate for the cross-sections we get from an early bubble-chamber experiment, that was able to isolate six $\bar{p}p \rightarrow 4K$ events in the momentum range from 1.6 to 2.2 GeV/c , one of which was identified as $\phi\phi$. The authors give an average cross-section of $3.8 \pm 1.7 \text{ } \mu\text{b}$ for the reaction $\bar{p}p \rightarrow 4K$ and an estimate of 600 nb for the intermediate state $\bar{p}p \rightarrow \phi\phi$ [3].

In order to obtain the high luminosities required due to the low expected cross-section, the JETSET experiment employs an internal hydrogen cluster jet target.

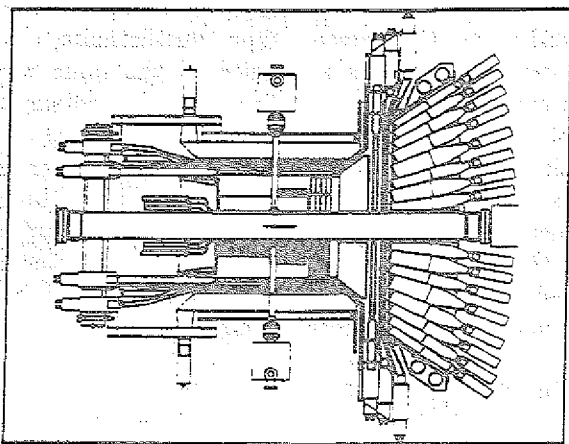


Figure 1: View of the JETSET-Detector

Fig. 1 shows a plan view of the compact detector surrounding the target region. It consists of sub-components that perform the tasks of triggering, tracking and particle identification. The online trigger relies on the multiplicity of charged particles seen in two sets of scintillators (one set of 60 strips surrounding the beampipe, and an "outer" three layer hodoscope [4] and on a system of segmented threshold Cerenkov counters that veto fast pions. An electromagnetic calorimeter and a gamma veto enable the rejection of photon events at the trigger level and the reconstruction of neutral particles (π^0, η) in the offline analysis. A tracker allows the reconstruction of particle directions, and silicon dE/dx -counters assist in the determination of the particle velocities.

The detector was mounted during the second half of 1990 and the LEAR shutdown in the winter of 1991. After a short "pilot run" in December 1990 we were able to identify the first 4-kaon-events. In the course of 1991 the experiment got three beam-times of 1-2 weeks each. In April and July a systematic scan over the momentum range from 1.2 to 2.0 GeV/c (corresponding to an energy range from 2.15 to 2.43 GeV) in steps of 0.1 GeV/c was performed. During a third run in October the region from 1.39 GeV/c to 1.480 GeV/c was scanned in small steps of .015 GeV/c . In parallel to the channel $\bar{p}p \rightarrow \phi\phi$ data were collected on the channel $\bar{p}p \rightarrow K_S^0 K_S^0$.

Three independent methods for calculating the luminosity are being developed. One uses the decay rate of the antiproton intensity stored in LEAR. The two others are based on dedicated triggers measuring the elastic cross-section that run in parallel to our main physics triggers. All three methods yield a value of 2 to $3 \times 10^{29} \text{ cm}^{-2}\text{s}^{-1}$ for the luminosity. This number is expected to rise for future runs, due to a higher LEAR intensity and to an improved target configuration.

A full Geant-based Monte Carlo has been developed and allows studies of the detector acceptance and the reconstruction efficiency.

In a three-step offline reduction of the data, events with kinematically forbidden topology and pathological events are rejected. The surviving sample of clean four-track events undergoes a full kinematical evaluation. There are 16 unknowns in the final state (the four-vectors of the four particles), and twelve measurements (the four-vector of the initial $\bar{p}p$ -state and the directions of the outgoing particles). Assuming the particle masses to be those of kaons, a kinematical reconstruction is possible,

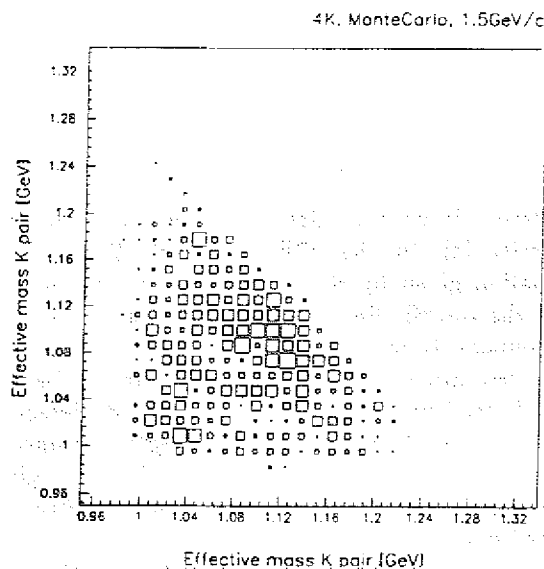


Figure 2: Two-dimensional Goldhaber plot of reconstructed invariant masses of k from Monte Carlo generated phase-space distributed four kaon events at 1.5 GeV/c incident $\bar{p}$ -momentum.

however giving several solutions for each event. A certain fraction of these solutions can be rejected because the calculated momenta for individual kaons are not compatible with the energy loss in the silicon pads or with the response of the Cerenkov counters. For every valid solution the invariant masses of each combination of kaon pairs are calculated, giving rise to three entries in a two-dimensional so-called Goldhaber plot of $M(K_i K_j)$ versus $M(K_k K_l)$. Note that, since the sign of the charge of the kaons is not measured, an unambiguous assignment of the kaon pairs is not possible. For a given center-of-mass energy the limits of the Goldhaber plot are given by the kinematics and form a triangle. A phase-space-distributed four kaon sample leads to a flat, featureless distribution within these borders (see Fig. 2). For a pure $\phi\phi$ -sample the Goldhaber plot would show a peak at the crossing of the ϕ -mass on each axis, plus a band near the hypotenuse of the triangle, that is due to the wrong combinations (recall that for each "correct" there are also two "wrong" entries; see Fig. 3). Fig. 4 shows a Goldhaber plot of data taken during the July run at 1.4 GeV/c. A conclusive statement on actual cross-sections certainly needs more thorough studies of the detector acceptance and of the reconstruction efficiency, but a surprisingly clear $\phi\phi$ -signal is seen. The analysis of the data at all measured momenta is in good progress and we are looking forward to interesting insights into this important aspect of physics.

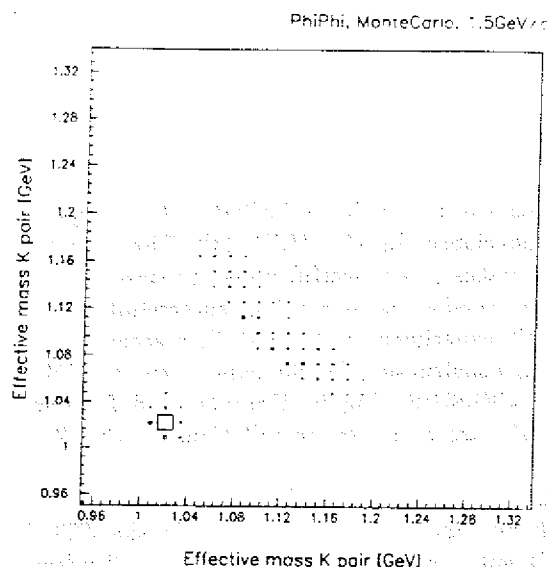


Figure 3: Goldhaber plot from Monte Carlo generated $\phi\phi$ events at 1.5 GeV/c incident beam-momentum.

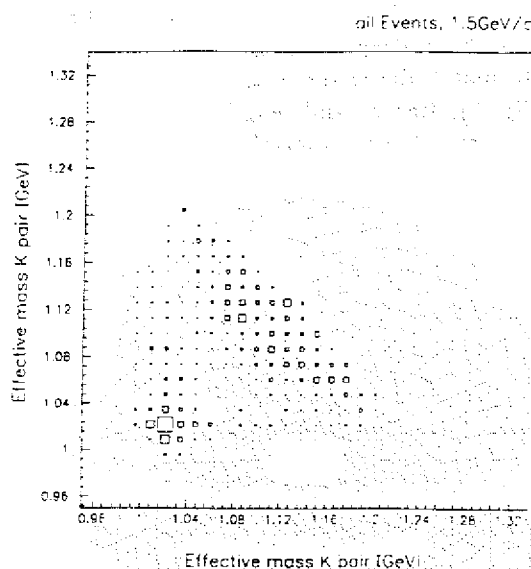


Figure 4: Goldhaber plot of invariant masses of kaon pairs, from data taken at an incident anti-proton momentum of 1.5 GeV/c. Note that the results are not yet acceptance-corrected.

References:

1. proposal P97, CERN/PSCC/86-23
2. A. Etkin et al., Phys. Lett. 201B (1988) 568
3. J. Davidson et al., Phys. Rev. D9 (1074) 77
4. O. Steinkamp et al., this report

A three-layer circular scintillator hodoscope has been developed in the IKP1 [1]. The symmetric design makes it a powerful multi-purpose detector. A first version is operating successfully in the JETSET experiment at CERN [2], a second version is under construction for the experiment PROMICE at the CELSIUS-ring in Uppsala, and for one of the major experimental installations at COSY, the Time-of-Flight spectrometer (TOF) very similar but larger hodoscopes will be made.

The hodoscope configuration used in the JETSET experiment, consists of a forward part ("endcap"), covering polar angles up to 45° and a barrel part covering angles from 45° to 135° . Both parts have three layers, one segmented into straight pieces, the other two segmented into pieces bent in opposite sense. The coincidence of one scintillator in each layer defines a "pixel", providing a spatial information on the particle trajectory. The outer radius of the endcap is 58 cm, the inner radius is made to match the shape of the JETSET beampipe (see. Fig. 1). Radius and length of the barrel part are 35 cm and 70 cm, respectively.

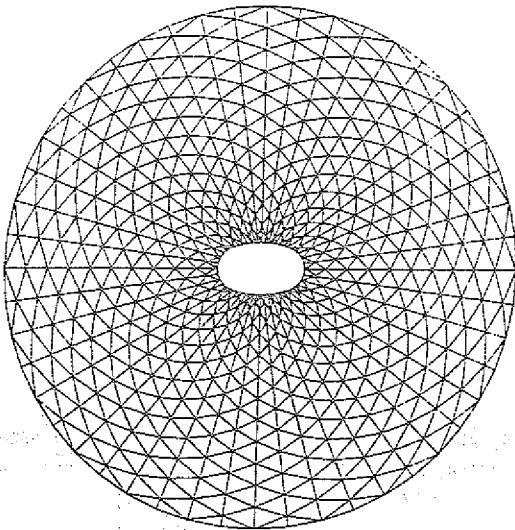


Figure 1: Schematic drawing of the endcap of the hodoscope.

The charged particle multiplicity information extracted from the three layers is used in the main trigger scheme [3] of JETSET. The timing of the straight scintillators defines the event-time (T_0) of triggers for delayed neutral decays (in particular one of the main physics channels, $\bar{p}p \rightarrow K_S^0 K_S^0$). A special trigger using pixels formed by the bent

scintillators is used to determine and monitor the luminosity [4]. In the offline analysis the spatial information given by the pixels plays an important role in the trackfinding algorithm. A time-of-flight measurement is in principle possible, the intrinsic timing resolution has been measured to be better than 0.8 ns. But due to the very compact detector design leading to short time-of-flight distances, and to a not too well defined T_0 it does not give significant information in the JETSET environment.

It turns out, however, that a determination of particle velocities seems to be possible by measuring the energy loss (dE/dx) of charged particles traversing the scintillators. The correlation between the particle velocity β and mean energy loss per unit length (Bethe and Bloch) is roughly $dE/dx \propto 1/\beta^2$. This correlation is smeared out by a series of statistical processes in the energy deposit itself (e.g. Landau distribution) and detector performance. This limits the resolution of the dE/dx measurement. Fig. 2 shows an ADC spectrum measured in one straight scintillator with the so-called pixel trigger [4].

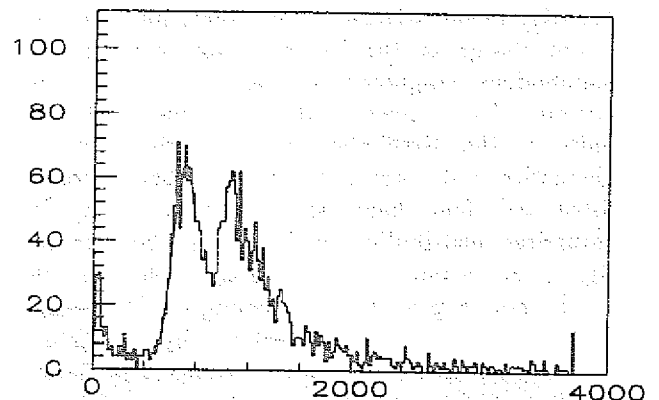


Figure 2: Spectrum of ADC values measured in a straight scintillator for one pixel. The peak at lower values is due to minimum ionizing particles ($\beta \approx 1$), the peak at higher values stems from elastically scattered protons with a β of 0.73.

This trigger measures elastic proton antiproton scattering at 90° in the center of mass system, thereby selecting protons of a given polar angle and a given β in the laboratory system. The spectrum shows a clear double-peak structure. The peak at higher values is due to the desired elastic events, the peak at lower values is due to a contamination with minimum ionizing particles that passed through the hardware trigger. By applying stricter cuts to select clean elastic events in the offline analysis this peak can be completely eliminated. From the width and the separation of the two peaks one can estimate a β -resolution of the order of 10 %. It should be possible to improve this value by taking the mean (or the truncated mean) of the three dE/dx measurements in the three layers of the hodoscope.

The light-collection efficiency is strongly dependent on the position of irradiation. For the bent scintillators variations of up to 30 % have been shown [5]. For this reason a careful calibration of the scintillators is necessary. This can be accomplished using the signals from minimum ionizing particles. Fig. 3 shows calibrated means of ADC distributions for pp elastic scattering at a series of beam-momenta (in 0.1 GeV/c steps from 1.2 GeV/c to 2.0 GeV/c) and for different pixels along the scintillator (corresponding to different intervals of the polar angle) plotted against the β of the particle calculated from kinematics. The plot shows a clear correlation between the calculated β and the measured energy loss.

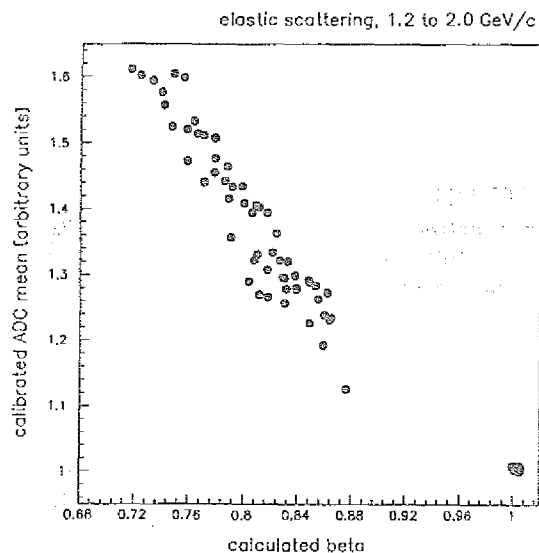


Figure 3: Calibrated means of ADC distributions (y-axis, in arbitrary units) plotted against the calculated β (x-axis), for elastic scattering at beam-momenta from 1.2 GeV/c to 2.0 GeV/c and polar angles from 15° to 30° .

References:

1. Th. Sefzick, Diplomarbeit, Bonn 1988,
A. Empl et al., Jül-Spez-562, April 1990, p. 213,
G. Decker et al., Jül-Spez-2462, 1991, p. 212.
2. proposal P97, CERN/PSCC/86-23
3. D. Maeckelburg et al., this report
4. K. Röhrich et al., this report
5. T. Empl, Diplomarbeit, Jül-Spez. 510, 1989

Determination of Luminosity with the Pixel Trigger of the Circular Scintillator Hodoscope

K. Röhrich, K. Kilian, W. Oelert, O. Steinkamp, B. Stugu
for the PS202 collaboration

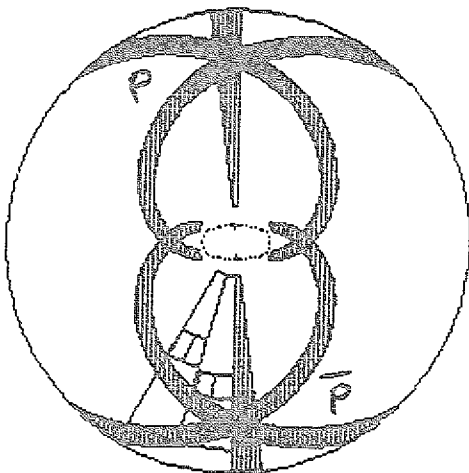
The JETSET experiment at CERN [1] measures the production of strange mesons in antiproton proton annihilations on an internal cluster hydrogen target in LEAR at CERN. There is no direct measurement of beam intensity or target thickness possible. The luminosity

$$L = N_p \cdot \nu_{\text{LEAR}} \cdot \rho_{\text{jet}} \cdot \int f(\bar{p})f(p)$$

might be determined from the jet density ρ_{jet} , the number of antiprotons in the machine N_p , the revolution frequency ν_{LEAR} and the geometrical overlap $\int f(\bar{p})f(p)$ of beam and jet. Because these numbers have large errors, (especially the overlap integral between beam and jet), an independent measurement of the luminosity is necessary.

We have set up a simple coincidence logic between the outer bend scintillation counters in the forward direction defining two pairs of pixels to measure the elastic antiproton proton scattering cross section at symmetric kinematics (90° in the center of mass system). The pixel positions are opposite in the horizontal and vertical axes of the experiment, thus coplanar of the accepted particles and providing two parallel measurements of the cross section.

11/07/1991 23:31:12.42 EVENT N:0 16065



TRIGGER BIT 6 Y -> X

Figure 1: An event of PS202 in the forward detectors, taken with the pixel trigger. Small boxes are calorimeter hits, empty wedges show hit Cerenkov counters and shaded are the forward hodoscope counters.

Figure 1 shows an event taken with the pixel trigger. The pixels are defined by the overlap of two bend scintillators. The recoiling proton penetrates the upper pixel, the antiproton is clearly identified by its annihilation in the calorimeter and the clustering of hits around its track.

Even so the trigger selects mainly elastic events — about 40 % of the raw triggers are finally accepted as elastic events — other reaction channels are accepted by the logic. In the event analysis those are rejected to a large fraction by simple multiplicity cuts. After track finding and fitting, the final reduction is done by checking elastic kinematics for each event.

The number of elastic events has to be corrected for deadtime and trigger and analysis efficiencies to get the true fraction of elastic events in the trigger. The deadtime easily can be calculated from the recorded rates. The trigger acceptance and analysis efficiency have been derived analyzing Monte Carlo generated events, taking the full detector (multiple scattering, energy loss, annihilations, ..) into account. The reconstruction efficiency is 66.0 ± 3.0 %. Events have been generated for the four pixels in ranges of $\pm 5^\circ$ and $\pm 10^\circ$ of polar and azimuthal angles, respectively. Thus defined, the geometrical acceptance is 11.0 ± 0.34 %. The luminosity for the experiments is then determined by comparison to formerly measured cross sections [2–4]. It can be determined with an accuracy of about 12 %, which originates from the uncertainty of the differential cross sections (2 %) and the statistical error from the pixel trigger (around 3 %). Acceptance and efficiency are determined with an statistical accuracy of about 5.5 % from Monte Carlo events. The systematic error mainly arises from differences between the Monte Carlo program and the real apparatus, especially from the uncertainty in the real position of the counters and the target. It can be estimated to be less than 2.5 %.

For the JETSET runs in April and July 1991 the luminosity at a beam momentum of 1500 MeV/c was around $(2.4 \pm 0.27) \cdot 10^{29} \text{ cm}^{-2}\text{sec}^{-1}$ and $(2.0 \pm 0.23) \cdot 10^{29} \text{ cm}^{-2}\text{sec}^{-1}$, respectively.

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1. proposal P97, CERN/PSCC/86–23
2. E. Eisenhandler et al., Nucl. Phys. B113 (1976) 1.
3. T.C. Baron et al., Nucl. Phys. B32 (1971) 66.
4. D.L. Parker et al., Nucl. Phys. B32 (1971) 29.

The JETSET Trigger System for the Circular Scintillator Hodoscope

D. Maeckelburg^a, W. Oelert, K. Röhrich, T. Sefzick, O. Steinkamp

The JETSET experiment at CERN [1] consists of a compact detector system surrounding the interaction region of the LEAR antiproton beam and a hydrogen cluster jet target. The JETSET collaboration is searching for exotic states of nuclear matter such as hybrids or glueballs, by looking for structures in the excitation function of gluon rich pp annihilation channels. The main reaction channel JETSET is searching for has four K mesons emerging directly from the pp interaction. Another channel is the $K_S K_S$ final state, which is detected via the delayed decay of the K_S into charged pions. In addition, reactions with uncharged particles in the exit channels are investigated.

To select the wanted events out of the huge number of interactions, sophisticated triggers provide the gates for the ADCs and the start signals for the TDCs in the experiment as well as a strobe for the readout. They are segmented to determine the multiplicity of particles in the event and to allow kinematical criteria to be applied online. The JETSET triggers use signals from both scintillation and Cerenkov counters and a calorimeter.

A major contribution from the IKP to the experiment is a circular forward and a barrel (45° to 135°) hodoscope, both consisting of three layers of straight and bend scintillators [2]. The hodoscopes surround all other JETSET detectors except the calorimeter. They are used in a fast trigger as a space point (pixel) providing device. We have setup a logic to exploit the features of the hodoscopes as part of the general JETSET trigger.

The analog signals from the hodoscopes are split and fed into ADCs and discriminators. The discriminated signals are recorded with TDCs. The discriminators used provide besides the logic signals an analog multiplicity signal. Both are used in the trigger, which follows two paths.

To determine the multiplicity of charged particles in the hodoscopes, the analog multiplicity information is used. The signals from all the discriminators of a single hodoscope layer are added. A multichannel discriminator with threshold settings according to multiplicities from 1 to 10 is used to transform the analog signal into a bit pattern onto a differential ECL bus. This bit pattern has ones in a row up to the actual multiplicity, from whereon it

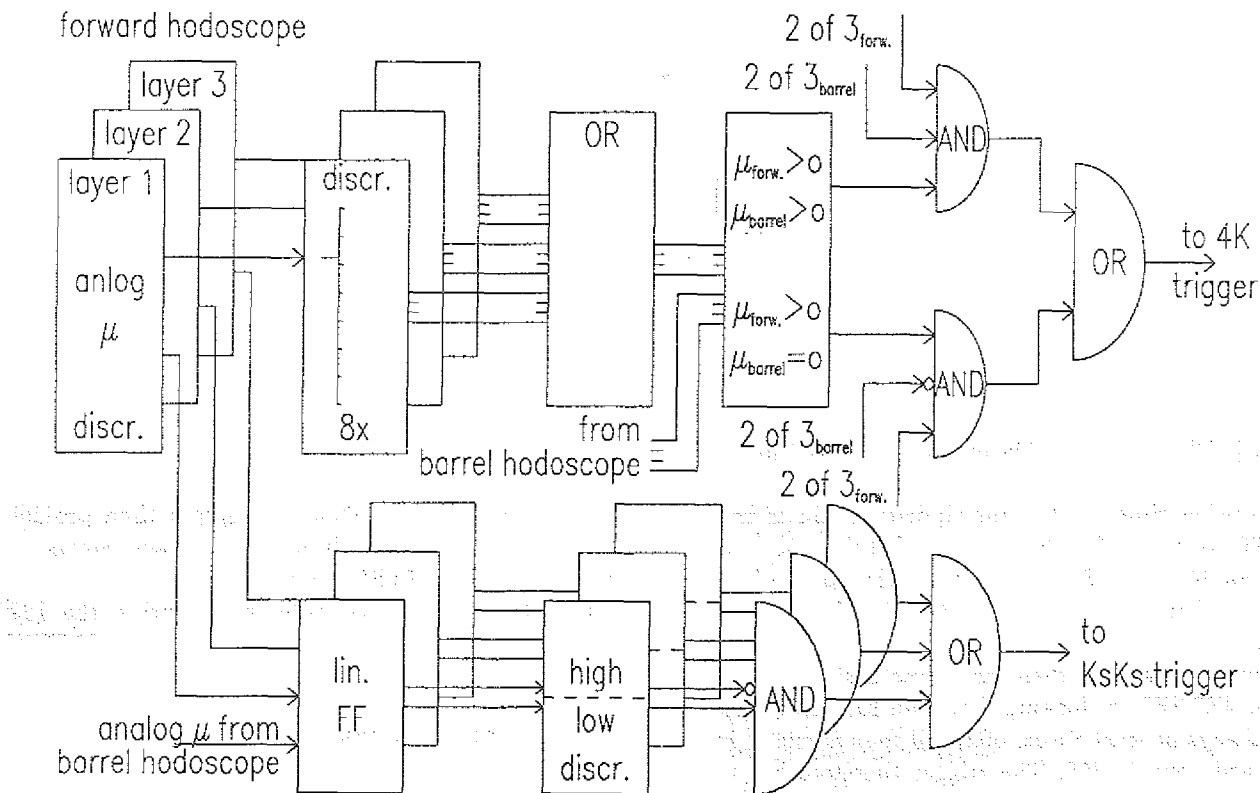


Figure 1: The outer scintillator multiplicity logic.

has 0's. The patterns of the three layers are "OR"ed, selecting the highest of the three multiplicities from the layers of a hodoscope. In a logic unit [3] eight different conditions can be applied to the multiplicity bit pattern and allow for eight different, but parallel decisions. Here the forward and barrel hodoscope multiplicities are combined and the final multiplicity trigger signal is generated. Due to the geometrical shape of the hodoscope counters, coincidences between counters of different layers define pixels. To form all possible combinations of counters the electronics needed was not available. Therefore a simple logic checks only, that at least two out of the three layers in each hodoscope have had a hit. This is a minimum requirement to assure that there is at least one good track candidate. The logic signals from all the counters of a single layer are "OR"ed. The three "OR" signals of a hodoscope are combined in all combinations of two in coincidences. The "OR" of the output of the three coincidences then gives the condition, that two out of three layers had a signal.

In figs. 1 and 2 a sketch of the two trigger logics is displayed. A further application of the hodoscopes as a luminosity monitor is described elsewhere [4].

From $K_S K_S$ decays we expect four charged pions. Because of their light mass, there is no such forward peaking kinematic. To accept also those events, we require a total multiplicity of 4 for the sum of multiplicities in the forward and barrel detector. The analog multiplicities from corresponding layers of the forward and barrel hodoscope are added. Discriminators select minimum multiplicities (fig. 1 bottom). The lower multiplicity is defined by the signal from the discriminator with the lower threshold. The upper limit comes from the discriminator with the higher threshold, whose output is used to veto events with too high multiplicity. After this multiplicity determination in the single layers the signals are "OR"ed to provide the final multiplicity trigger.

Events with only neutral particles have signals only from the calorimeter. The hodoscopes are therefore used as veto counters. As a veto signal the "two out of three" trigger is used, because it allows for a noise hit, but has the minimum requirement for a charged particle.

The different hodoscope informations are combined with the signals from other detectors to form the final event trigger. The triggers for the different

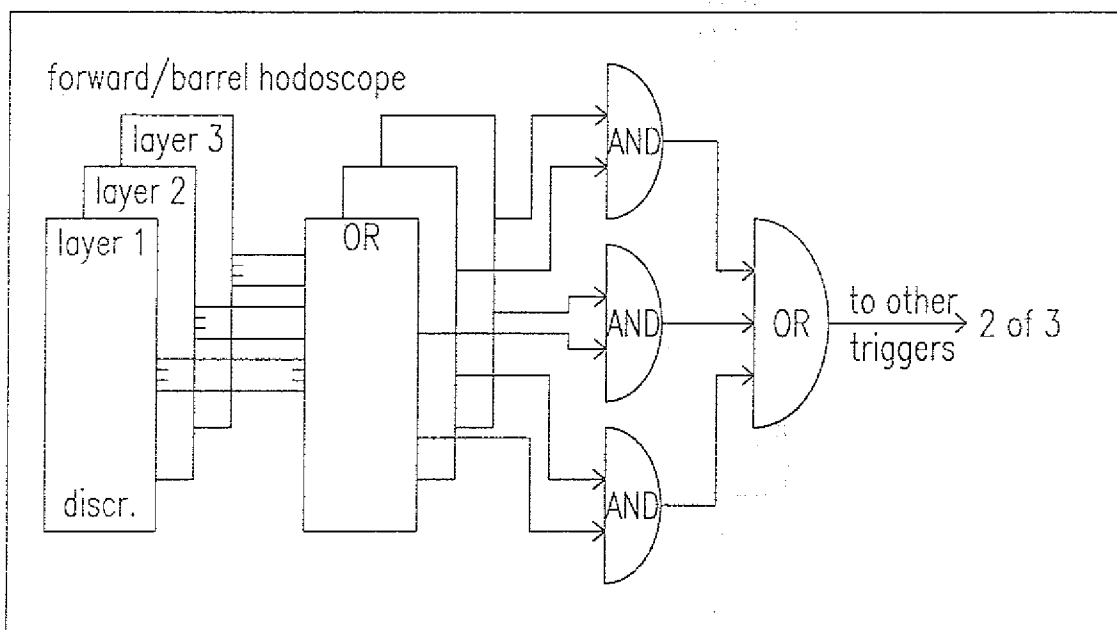


Figure 2: The outer scintillator "two of three" logic.

The decision time for the multiplicity chain is less than 90 nsec and for the "two of three" logic it is less than 55 nsec. For the combination of both it needs another 20 nsec, in total the decision time is less than 120 nsec.

As mentioned above, there are three main types of events, JETSET is looking for. The four K events have always at least three, often all four particles in a forward cone of 45° . The trigger therefore has to allow a multiplicity 3 or 4 in the forward detector. In the barrel the only allowed multiplicity is 1. It is used as a veto counter if all 4 particles are in the forward hodoscope.

events are "OR"ed. This final signal then provides the gates for the ADCs and TDCs and starts the readout of the JETSET apparatus.

The trigger was successfully used during the 1990 and 1991 data taking periods of the JETSET collaboration.

a ZEL, KFA Jülich

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3. LeCroy, New York, ECLINE model 2365 8LM.
4. K. Röhrich et al., JETSET note 91-24.

Correction of the JETSET Experimental Data by Monte-Carlo Simulation

B. Stugu, K. Kilian, W. Oelert, K. Röhrich, O. Steinkamp

The aim of the experiment PS202 (JETSET) at CERN's LEAR is to measure the cross sections for the two reactions $\bar{p}p \rightarrow 4K$ and $\bar{p}p \rightarrow \phi\phi \rightarrow 4K$ as function of energy. In order to obtain an estimate of the cross sections, one has to carefully evaluate all entries in the expression: $\sigma = (N - N_b) / L \alpha \epsilon$, where N is the number of examples of the reaction found, N_b is the estimated number of background events contaminating the sample, L is the integrated luminosity, and $\alpha \epsilon$ is the combined efficiency stemming from the geometrical acceptance and efficiency of the detector and of the data analysis chain.

To evaluate $\alpha \epsilon$ a detailed Monte-Carlo simulation program is developed. This program generates the kinematics of the events and tracks them through the detector. The output of the program is coded raw data, which in turn is analyzed as true data by the JETSET analysis chain. Unfortunately the CPU requirements for use of the program are large whenever reasonable statistics is required to get through the whole analysis chain. It turns out that computer resources at the KFA in Jülich are available for production of reasonable amounts of MC data. A total of 30.000 events of each of the two reactions were generated at each of the nine lab. momenta where experimental data were collected. The CPU time required for this operation was about 75 hours on the central IBM computer of the KFA. These data have been copied to tape, and are made available to all members of the JETSET collaboration.

Efficiencies can be estimated when running the data through the JETSET analysis chain. In fig. 1 we see the efficiency obtained as function of lab. momentum for the two reactions, assuming that the events are distributed isotropically in phase space.

The essential proof for $\phi\phi$ events is to see a peak in the bidimensional plot where the invariant mass of one pair of kaons is plotted against the mass of the other pair. To get an estimate for the background under the peak, the detailed acceptance in this plot has to be known. In fig. 2 we depict such an acceptance plot. The experimental data have to be corrected bin by bin in this plot when estimating the number of events which were genuine $\phi\phi$. Note that the acceptance is enhanced in the $\phi\phi$ region. This correction should thus be done very carefully.

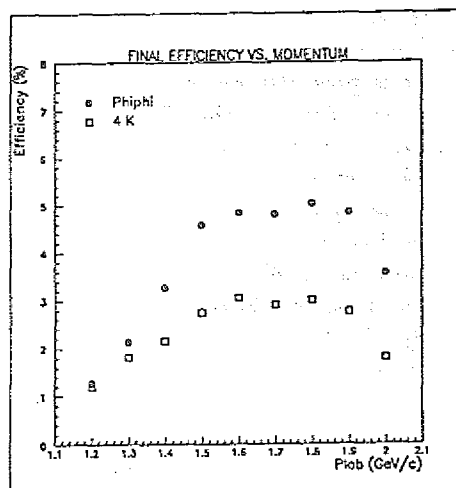


Figure 1: Final reconstruction efficiency vs lab. momentum.

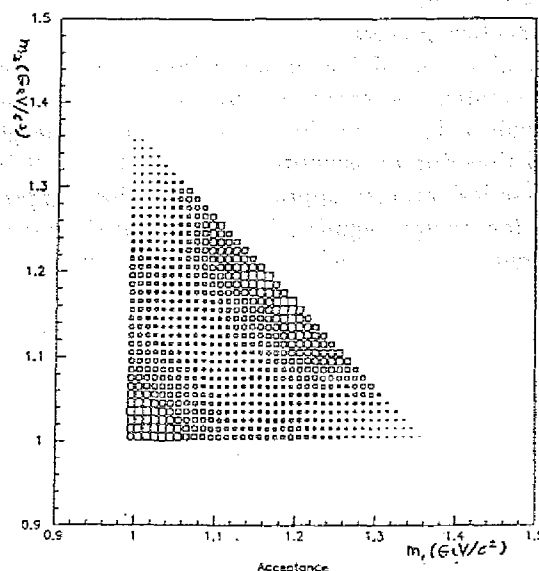


Figure 2: Acceptance for $\bar{p}p \rightarrow 4K$ in the bidimensional mass plot for a lab. momentum of 1.8 GeV/c. The size of the squares is proportional to the acceptance in each bin.

Heavy Hyperon—Antihyperon Production

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Super LEAR, a storage ring for antiprotons with superconductive magnets and stochastic beam cooling in the momentum range from 2 GeV/c up to 12 GeV/c is presently being discussed at CERN. A natural extension of the $S=1$ hyperon—antihyperon physics program, which is studied successfully at LEAR, would be the investigation of $S=2$, $S=3$ or $C=1$ heavy flavour production.

Hyperon—antihyperon ($Y\bar{Y}$) production from the proton—antiproton ($p\bar{p}$) interaction requires the creation of at least one heavy quark pair ($s\bar{s}$, $c\bar{c}$ or $b\bar{b}$) preceded by the annihilation of an equivalent number of $q\bar{q}$ pairs in the initial state.

Direct observation of quark dynamics appears to be impossible due to colour confinements in baryon— or meson—like bags. Due to the parity non—conserving weak decay of hyperons, large asymmetry parameters with respect to the polarizations and spin correlations may allow one to make use of the self—analysing features to determine spin observables of the final state particles (Y or $\bar{Y}$) which possibly can be related to the quarks and $q\bar{q}$ production process.

The gluon—mediated $q\bar{q}$ annihilation followed by a $q\bar{q}$ creation is characterized in the high energy domain (Fig. 1 top) by the exchange of one gluon and thus due to asymptotic freedom the underlying dynamical process appears to be rather simple. In the low energy regime (Fig. 1 bottom) the situation becomes more complicated due to infrared slavery and the nature of the possible multi gluonic intermediate matter is very complex.

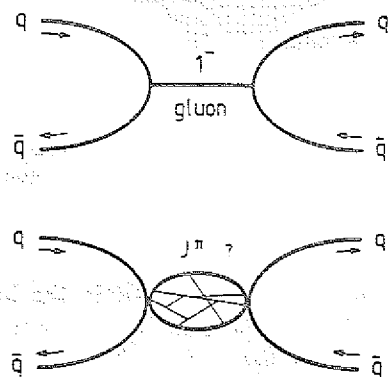


Figure 1: Basic mechanisms for $q\bar{q}$ annihilation and creation in the gluon—mediated picture. High momentum and low momentum characteristics are shown on top and bottom, respectively.

It seems that calculations employing the constituent quark model for describing the $q\bar{q}$ creation preceded by the annihilation process — since they do not involve the quark degrees of freedom explicitly — are on a somewhat less conclusive ground than those hidden flavour production considerations referring to the t —channel meson exchange model. Figure 2 presents hyperon—antihyperon production diagrams of simplest order in the quark—gluon and the meson—nucleon framework, respectively. The recently available calculations [1] on the threshold $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ reaction show that the following ingredients seem to play appreciable roles:

- initial and final state interactions, often developed from one—boson—exchange versions of NN potentials, resulting in strong short—range absorption;
- strong tensor interaction due to combined exchange of the $K(494)$, $K^*(892)$, and even $K^{**}(1430)$ mesons;
- both s —wave and p —wave transition amplitudes at low momenta; for excess energies at or above about 50 MeV even higher partial waves come into play;
- coupled—channels effects;
- a possible existence and influence of resonant behaviour at/or below the threshold [2].

The theoretical investigations mentioned so far have been essentially limited to the threshold region of the $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ reaction with excess energies ≤ 100 MeV. Kroll et al. [3] give predictions i) at higher energies well above threshold and ii) for other hyperon—antihyperon channels with heavy flavour following the $p\bar{p}$ annihilation process. These authors introduce within a constituent scattering model spatially extended diquarks as elementary constituents of baryons in addition to point—like quarks. The diquark description seems to be essential for reactions requiring the annihilation of two $q\bar{q}$ pairs in the $\bar{p}p \rightarrow \bar{\Lambda}\Lambda$ reaction and in the consideration of spin effects.

The special feature of the hyperons is their flavour—changing weak decay. Due to the large mass of the charmed hyperons (>2.2 GeV/c²) there is a large phase space available for several decay modes. The strange hyperons, on the other hand, show rather simple decay features with high branching ratios and group in two categories, see Fig. 3 for examples:

- the charged hyperons having $c\tau$ values of a few (2—4) cm; decaying to a charged meson and a Λ^0 which then decays to the πp system,

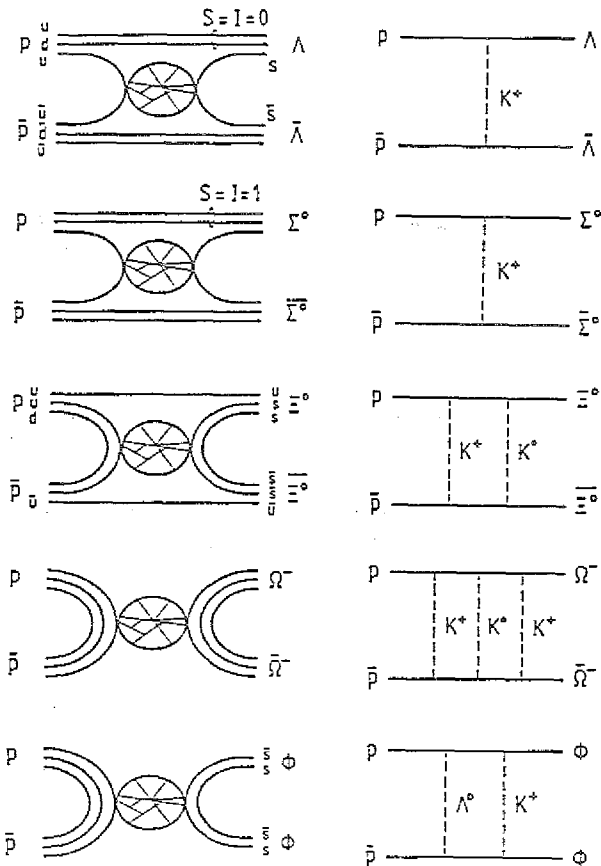


Figure 2: Quarkline and meson exchange diagrams for hyperon-antihyperon and $\phi\phi$ -meson production.

allows a multiplicity step of two (charged Y and $\bar{Y}$) to six to be applied as a clean trigger,

- b) the neutral hyperons having $c\tau$ values of some (7–9) centimeters and thereafter decaying via the Λ^0 to a two-prong πp system, allows a neutral-charged trigger to be applied experimentally.

These clear decay signatures result naturally in a similar experimental setup as used successfully by the PS185 collaboration at LEAR [4]. Embedding this detector into a time-of-flight spectrometer as discussed for COSY in Jülich [5] would certainly increase the reliability and precision drastically. A detailed detector device has to be worked out. But due to the special decay feature of the strange hyperons, and the consequent trigger logic as described above, a setup at an external beam line seems to be essential.

A systematic study of antihyperon-hyperon production with increasing strangeness content (see Fig. 2) is interesting due to the change of quark dynamics.

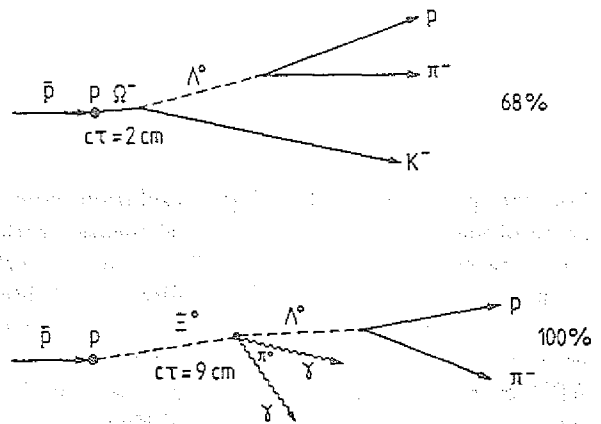


Figure 3: Examples of decay modes for two hyperons.

There are two special features of interest for the $\bar{\Omega}\Omega$ production: i) two spin 3/2 objects are created out of the pp interaction and ii) since the Ω is a $SU(3)$ colour singlet state consisting of strange quarks only, the $\bar{\Omega}\Omega$ -pair will be produced after total annihilation of the valence quarks in the initial state. Here investigating the excitation function of the $\bar{\Omega}\Omega$ production is suggested.

As outlined in [7] it is certainly an interesting question whether there is or is not a correlation between the incident $\bar{p}$ and the produced $\bar{\Omega}$.

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Performance Study of Doubly Focussing Crystals for a Low Energy Bragg Spectrometer

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For the preparation of a high resolution study of antiprotonic X-rays in light elements with a focussing crystal spectrometer [1,2], a major direction of investigation concerns the realization of ideally curved crystals. In the energy range of interest between 1.7 keV and 11 keV quartz, Si, LiF and ADP crystals are candidates to be used as X-ray reflectors in a Bragg reflection diffractometer. Of special interest is the option to shape the crystals spherically because of the doubly focussing behaviour. The realization depends critically on the perfectness of the lattice structure and the elastical properties of the crystals. Therefore we started a systematical study to develop a technique to bend crystals elastically to spherical shape [3].

In the conventional method the crystal was fixed with epoxy onto the concave surface of an optical lens by pressing it with the help of a convex lens during the hardening time of the glue.

In an alternative procedure the crystal is put into a pressure vessel and glued onto the concave lens while a pressure of about 8 bars is acting.

The curvature of the final mounting is checked by means of an optical set up using the imaging properties of the concave surface. As a crucial test the crystal mounting is placed in the spectrometer and irradiated with characteristic fluorescence X-rays excited by means of an X-ray tube with various targets.

The X-ray image is registered in a two dimensional MWPC with a position resolution of about $150 \mu\text{m}$. As a typical example the $K\alpha$ doublets of Cl(2.62 keV) and Cu(8.04 keV) reflected by a Si(111) crystal in first (Cl) and third order (Cu) are shown in Fig. 1.

From the fit of the spectra with Voigt-functions information about the crystal rocking curve and the natural line width has been extracted. In the case of Cl it is necessary, for a satisfactory description of the spectra, to use a rocking curve that is composed of a Gaussian and a Lorentzian contribution. The extracted crystal reflection width is about 30 arc sec for the first order reflection.

In the case of Cu the width of the rocking curve is negligible as compared to the natural line width. Therefore, a precise determination of the natural line width of the $K\alpha$ lines became possible. A value for the $K\alpha_1$ of $2.7 \pm 0.1 \text{ eV}$ and for the $K\alpha_2$ of $3.2 \pm 0.1 \text{ eV}$ was extracted. The difference in the widths is attributed to a different strength of the corresponding Coster-Kronig transitions.

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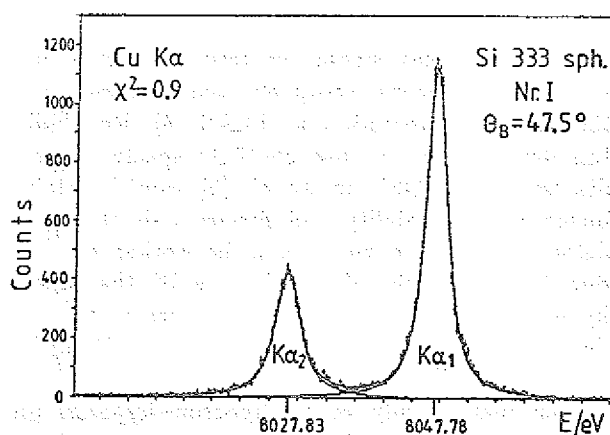
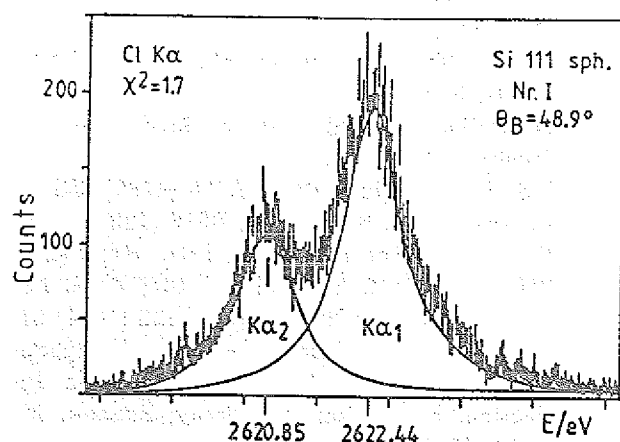


Fig. 1: $K\alpha$ doublets of chlorine and copper measured with a spherically bent Si(111) crystal. The fit is the folding of the natural line width and the response function of the apparatus. In Cl the natural line width is 0.93 eV [4], thus dominating the total width. In the case of Cu the crystal reflection width is negligible as compared to the natural width.

γ -Decay of the Giant Dipole Resonance in Hot Rotating Rare Earth Nuclei

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D.J. Horen², I.Y. Lee², D.G. Sarantites³, D.W. Stracener³, and R.L. Varner²

We have investigated the temperature and spin dependence of the Giant Dipole Resonance (GDR) in highly excited rare earth nuclei by measuring continuous γ -ray spectra in the reaction $^{16}\text{O} + ^{159}\text{Tb}$. The experiments have been performed using beams of 160, 216, 300 and 400 MeV ^{16}O ions from the Holifield Heavy Ion Research Facility at Oak Ridge. The γ -rays were detected in a 4π crystal ball consisting of 72 NaI detectors and in 4 arrays of each 19 BaF₂ detectors in coincidence with complete and incomplete fusion products. Neutrons and γ -rays were separated by time-of-flight and in case of the BaF₂ additionally by pulse shapes. Results from the data analysis are discussed in refs. 1-3.

A critical point in the analysis is the separation of neutrons and γ -rays at high energies. Since the neutron efficiency of NaI crystals is very high a separation of neutron and γ events as clean as possible is required. High energy neutrons can hardly be distinguished from γ events by the time-of-flight technique. However, the main contribution of high energy neutrons, mainly due to preequilibrium emission, is at forward angles. Therefore, we created new γ energy spectra summing up only those NaI detectors having an angle of $\theta_{\text{lab}} > 130^\circ$. In addition, we added these detectors which were placed in forward direction at a distance of 1 m to the target. Due to the much larger target distance these detectors allow a neutron discrimination which is very clean.

Further, a new response function was applied to the NaI spectra which takes into account the scattering of high energy γ -rays into neighbouring detectors. In the previous analysis this effect was only accounted for in the scanning software.

Finally, the treatment of the hard photon yield in the γ -spectra, which is mainly due to the bremsstrahlung effect in the reaction process, has been improved. A reasonable description of the slope of the bremsstrahlung yield was achieved in applying a systematics which is discussed in ref. 4. The amplitude was obtained from a fit to the data. By this, we obtain a description of the bremsstrahlung contribution which is in agreement with other available data. In fig. 1 four γ spectra are shown for 300 and 400 MeV at different momentum transfers with the bremsstrahlung yield indicated by the solid line. The slope parameter E_0 obtained from the systematics in ref. 4 provides a reasonable description of the data above 22 MeV.

The above discussed improvements in the treatment of the NaI spectra were necessary to allow a significantly better analysis of the data. The complete analysis of the new NaI spectra is in progress. First results already show slightly different GDR strength functions.

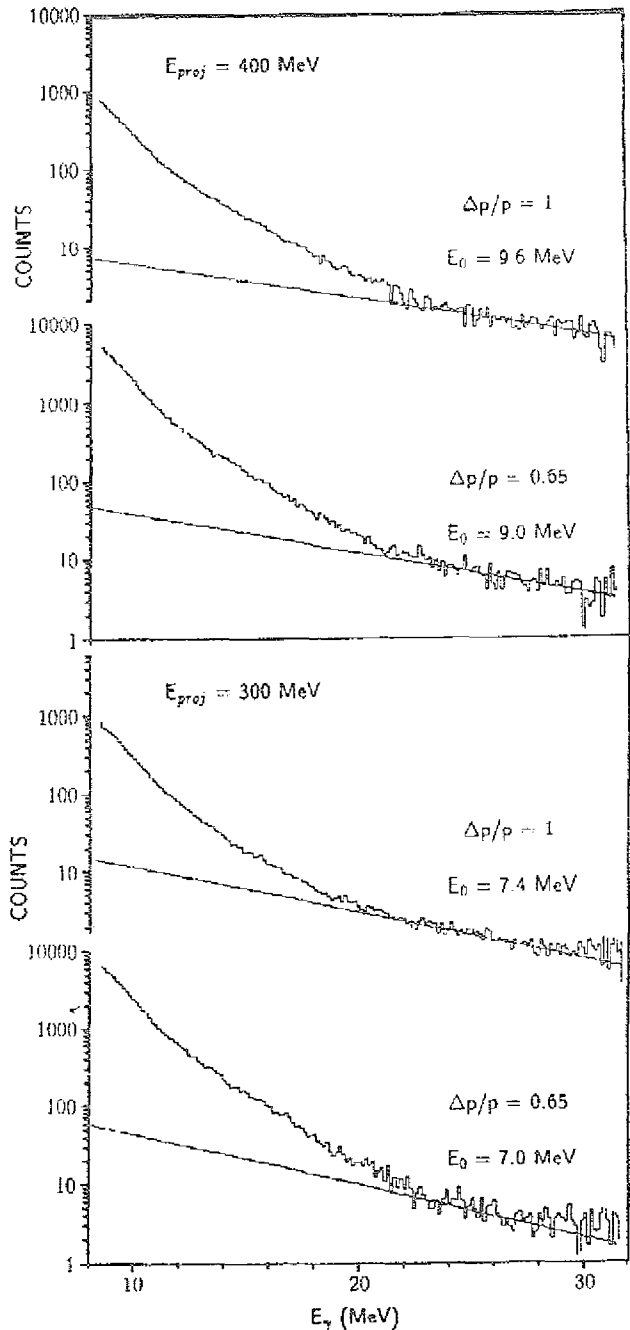


Figure 1: γ -energy spectra for 300 and 400 MeV at different momentum transfers. The bremsstrahlung yield, described by an exponential function with slope parameter E_0 , is shown by the solid lines.

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Radial Excitation of the Nucleon to the P_{11} (1440 MeV) Resonance in Alpha-Proton Scattering

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The detailed experimental study of the baryon properties presents an exciting challenge in hadronic physics for the next decade since it is related to the understanding of the structure of QCD in the non-perturbative regime. One of the most basic degrees of freedom is the size of the baryon. Its static properties are given e.g. by the charge densities studied in electron scattering. Dynamical properties of the size degree of freedom can be studied in radial modes (isoscalar monopole excitations) which give information on the compressibility of the system. A candidate for a radial mode of the nucleon is the Roper resonance P_{11} (1440 MeV).

The experiments on the Roper resonance excitation [1] in α -scattering at SATURNE at a beam momentum of 7 GeV/c were completed by taking additional data in the region of the π -threshold and measuring the target thickness to determine absolute cross sections. Missing energy spectra ($\Omega = E_i - E_f$) were measured at four scattering angles of 0.8, 2.0, 3.2 and 4.1 degree, those obtained for the two smallest angles are shown in fig. 1. The irregularities in some parts of the spectra are caused by errors in the trajectory reconstruction due to small field distortions in the wire chambers. The spectra indicate a strong rise of the yield above the π -threshold and a pronounced structure above 400 MeV.

In the spectra we expect two contributions; in addition to direct excitation of the proton we expect also excitation of the projectile, mainly to the Δ (1232 MeV) resonance. Whereas for α -p scattering the target Δ excitation should be forbidden, there are no selection rules which inhibit Δ excitation of the projectile. By emission of a pion this excitation decays favorably back into the α -particle ground state observed in the detector. Evidence for this process is found in older spectra of α -proton scattering measured at SATURNE [2], where a strong enhancement was observed above the π -threshold. Calculations of projectile excitation to the Δ resonance have been performed within the impulse approximation using the meson exchange model [3], in which these spectra could be well described. The appearance of this contribution in the missing energy spectra close to the pion threshold is due to the projectile velocity giving rise to a Lorentz boost. Here, it should be mentioned, that the spectral shape obtained for the projectile excitation is quite independent of the Δ resonance parameters. Even assuming, e.g. a width of the Δ resonance in the projectile twice as large as in the

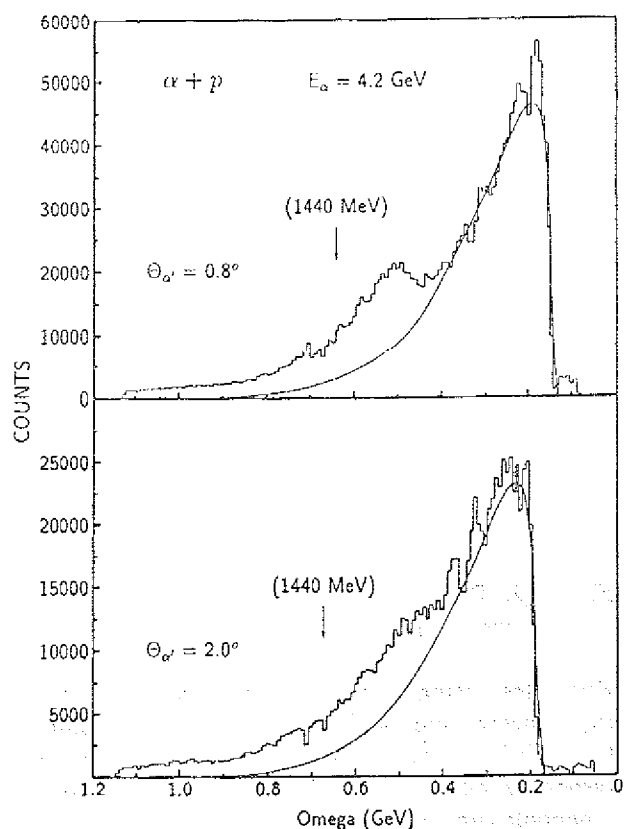


Figure 1: Missing energy spectra of inelastically scattered α -particles on a hydrogen target. The solid lines show the spectral shapes calculated for projectile excitation which were adjusted to the data above the π -threshold.

free case would hardly change the shape in the spectral distribution. A calculation of this contribution, adjusted in the spectral height to the data above the π -threshold, gives rise to the solid lines in fig. 1. The shape of the spectrum is quite well reproduced up to a value of Ω of 400 MeV. At larger values of Ω we observe a structure indicating a strong excitation of the P_{11} (1440 MeV) resonance region. The difference spectrum — in which the projectile excitation is subtracted from the measured spectrum — is shown in fig. 2 which shows a pronounced peak. Above 0.9 GeV the yield is rather flat. This is the region of the D_{13} (1520 MeV) and S_{11} (1535 MeV) resonance which can be excited in inelastic α -scattering by $L=1$ transfer. Differential cross sections are given in fig. 3. They show a very steep angular dependence characteristic of monopole transitions (discussed below).

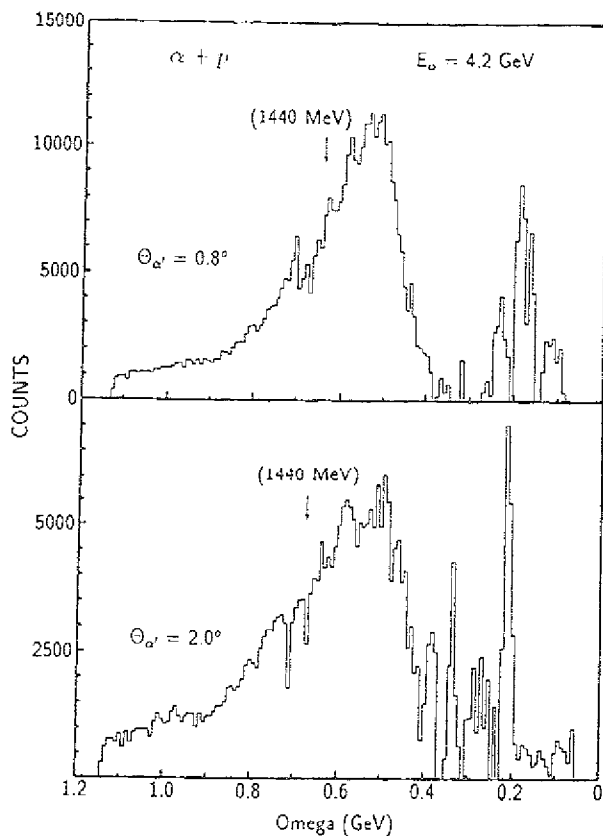


Figure 2: Difference spectra of the data and the calculations in fig. 1.

Estimates using simple scaling assumptions of experimental monopole cross sections [4] indicate that the large cross sections obtained for the Roper resonance excitation imply a large fraction of the monopole sum rule strength.

In the experimental spectra in fig. 2 the monopole strength peaks at a value of Omega which corresponds to an excitation energy of 410–420 MeV with a width of about 120 MeV. Due to the momentum transfer dependence of the α -particle form factor, which has to be taken into account, the strength is shifted to lower values of Omega. To obtain the monopole strength function, the form factor dependence has to be unfolded. This is only known for small momentum transfers [5]; the extrapolation gives rise to uncertainties in the extracted resonance parameters. Our estimates give a shift of only 30–50 MeV in the peak energy and a 30–60 MeV increase in the resonance width indicating a P_{11} -resonance at a mass of about 1400 MeV with a width of about 160–170 MeV. The energy is rather low in comparison with the average energy from π -N phase shift analyses [6]; the position corresponding to 1440 MeV (by including the effect of the α -form factor) is indicated by the arrows in figs. 1 and 2. An analysis is in progress in which the differential cross sections are described in a folding model approach including distorted wave effects. The obtained shape of the angular distribution (normalized in height to the data) is in quite good agreement with the data (solid line in fig. 3).

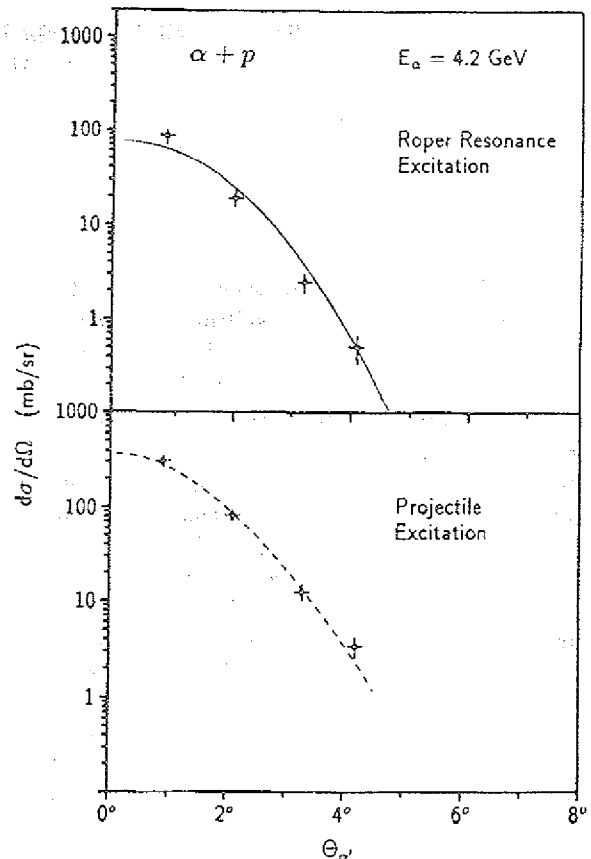


Figure 3: Differential cross sections for excitation of the P_{11} resonance excitation (upper part) and of the projectile (lower part). The errors given are mainly due to uncertainties in the separation of target and projectile excitation. The solid line shows the calculated shape for monopole excitation. The dashed lines indicate the consistency of both angular distributions discussed in the text.

Further support for a monopole character of the excitation results from a comparison of the data in the upper and lower part of fig. 3. Projectile excitation is dominated by the Δ resonance which presents also a monopole ($L=0$) excitation (however, of spin-isospin structure). Therefore, the angular distribution should be similar for both cases. However, different from target excitation, a pion is emitted in the projectile excitation process. Due to the pion emission the α -direction is smeared out and lies in a cone with an opening angle depending on the energy of the decaying resonance. In this case the monopole angular distribution should be folded over all pion angles. The data in the lower part of fig. 3 are well described by the dashed line assuming a Gaussian folding with a mean width corresponding to $E_x = 300$ MeV. This indicates a consistent description of both processes.

The nucleon compressibility K_N is related to the excitation energy E_x of the monopole mode by operator sum rules [7]. This leads to

$$K_N = (m_0/\hbar^2) E_x^2 < r_N^2 >, \quad (1)$$

where m_0 is the effective mass in the kinetic energy term and $\langle r_N^2 \rangle$ relates to the rms radius of the isoscalar density.

In the constituent quark model m_0 is the constituent quark mass m_q . However, in this model the ground state rms radius is directly related to the oscillator constant by the virial theorem giving rise to a nucleon radius which is far too small, ($\langle r_N^2 \rangle \sim 0.37 \text{ fm}^2$). Therefore, the radial properties can only be understood by the inclusion of a meson cloud. Assuming a radial vibration of the meson cloud gives $m_0 = 2 m_\pi$. The value of $\langle r_N^2 \rangle$ is obtained from the rms radius difference of proton and neutron density, which give $\langle r_N^2 \rangle = 0.776 \text{ fm}$. This leads to K_N in the range of 1150–1370 MeV, which can be considered as a lower limit since the effective mass m_0 may be higher, also we expect monopole strength at higher E_x due to quark excitation. From simple estimates we observe only 50–70 % of the sum rule strength in our peak. This gives values of the compressibility K_N in the order of 1500–2200 MeV. Here, it should also be mentioned that eq. (1) has been derived from a one density approximation which might give results somewhat different from a correct treatment of quarks and mesons. For a better determination of K_N more experimental information is needed.

This can be obtained from the inelastic cross sections which should allow to extract the Roper resonance $\langle r^2 \rangle$ -transition matrix element. Values of the compressibility can be derived from different models. For the constituent quark model a value of K_N of 3 GeV is obtained using a harmonic oscillator with $\hbar\omega = 500 \text{ MeV}$. More complicated versions of the model as well as MIT bag estimates are discussed in ref. 8 and give K_N in the order of 900–1200 MeV.

Presently, we are trying to describe the absolute cross sections in a folding approach in which the different densities involved are folded with an effective isoscalar interaction.

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ZEUS Barrel Calorimeter Assembly and Testing

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The ZEUS detector is one of the experiments used at the high energy beam collider HERA-DESY (820 GeV protons and 30 GeV electrons). Within the frame of the ZEUS-DESY-collaboration ANL-DESY-KFA have worked together in the assembling and final preparing of the barrel calorimeter modules for the ZEUS detector [1] and see Fig. 2. Within the ZEUS-DESY collaboration and the AM-ZEUS group the fabrication of barrel calorimeter modules for the ZEUS detector (BCAL) began at Jülich in February 1990 with the stacking of hadronic sections (HAC) and in September 1990 was expanded to include complete assembly of modules. As described in a previous publication in this series [2], the ZEUS calorimeter uses the sampling techniques with depleted uranium in the form of 3 mm thick plates as the passive element alternated with plastic scintillator as the active element. Scintillator light is collected by wave length shifter panels over both sides of each module. By connecting these panels to photomultiplier tubes, 81 for each side, the energy deposited in the calorimeter by a particle is read out from each side of the module. Ten of the 32 BCAL modules were assembled at Jülich from the HAC sections stacked there and the electromagnetic (EMC) sections stacked at Argonne.

The remaining components came from the Ohio State and Louisiana State Universities, University of Wisconsin and the Virginia Polytechnic Institute.

Modules fully assembled at Argonne underwent extensive quality assurance checks. Modules were checked by measuring their response to hardened cosmic radiation and the response to a 2 mCi Cobalt point source, which moved through conduits in the module. Although it was not practical to duplicate this effort at Jülich, a testing station was constructed, which allowed the photoelectron yield of each section of the calorimeter to be measured. This parameter is a critical characteristic of the module. It directly affects the energy resolution attainable and was used as a measure of the quality of modules assembled at Jülich.

The tasks carried out by this testing station were to supply the control voltage to the Cockcroft-Walton PMT bases, where the high voltage was generated and to read out the charge generated by each PMT as a result of the activity of the uranium. This "DC"

current in each PMT is used to monitor the stability of modules in the ZEUS detector. The PMT voltage is adjusted to attain a definite current of each section of the calorimeter module. The photoelectron yield may then be computed from the required voltage and the known gain of the PMT.

The data acquisition was done using an IBM PC interfaced to a CAMAC controlled system, that provided the PMT CW-base control voltage (multiplexed CAMAC DAC modules) and readout of the PMT charge (multiplexed CAMAC ADC modules).

The mean photoelectron yield of the EMC and HAC sections of the calorimeter modules assembled at Jülich were in good agreement with similar measurements for modules at Argonne. This comparison is shown in Fig. 1 for one of the dual modules readouts; similar agreement was observed for the other module readout.

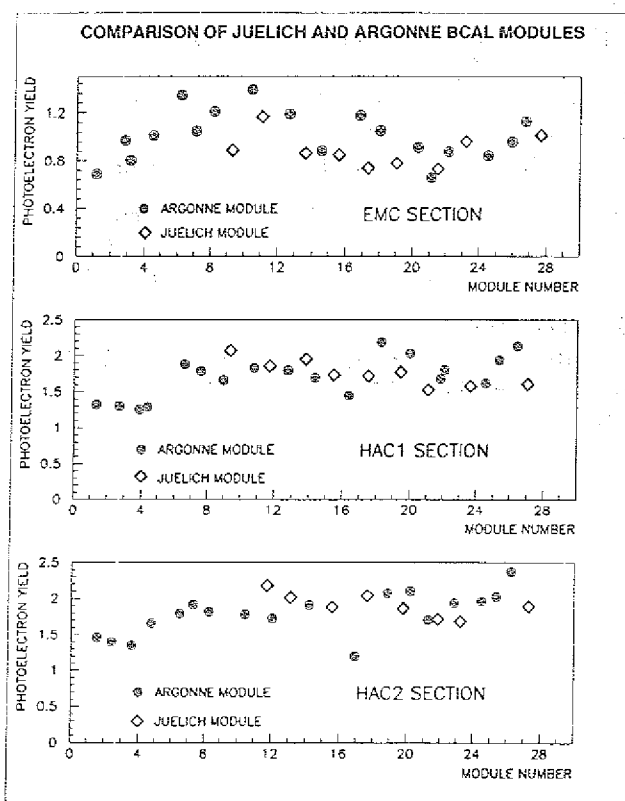


Fig. 1: Comparison of Jülich and Argonne BCAL modules

All of the modules assembled at Jülich were inserted into the ZEUS detector at the HERA hall in May 1991. The barrel calorimeter was successfully tested in September 1991.

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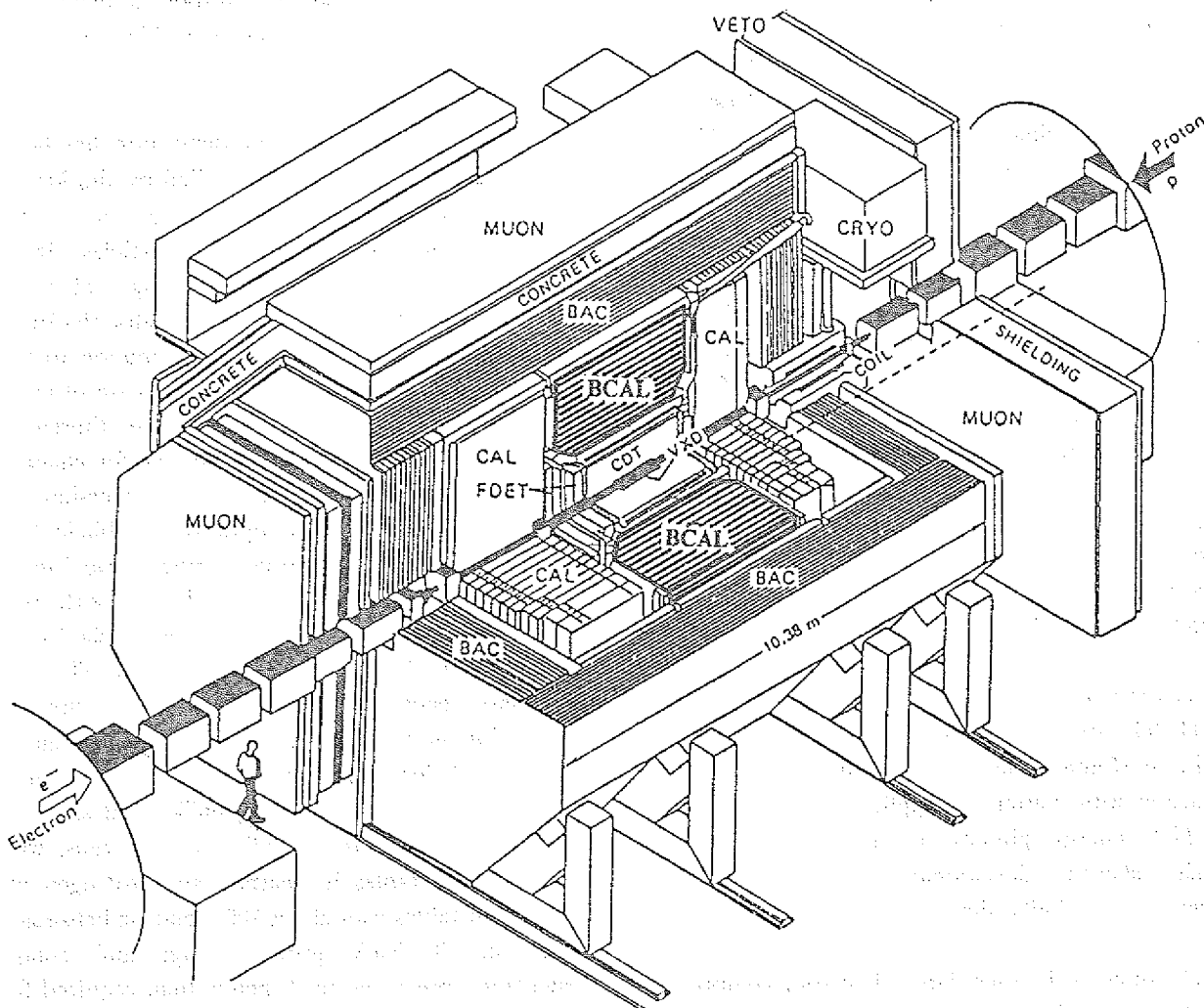


Fig. 2: View of the HERA-ZEUS detector with the barrel calorimeter (BCAL) in the center

Development of a Vectorized Monte Carlo Code for the Simulation of Particle Transport in the HERA-ZEUS Calorimeter

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The large and heavily structured particle detectors used in high energy physics require detailed Monte Carlo simulation for interpretation of their responses /1/. The need for computer power along with such experiments is extremely high and tends to become prohibitive with the new generation of accelerators as e.g. the ZEUS detector at the HERA ring, which has come into operation recently. One of the main reasons for the growth in computer time consumption is the simple relationship between the energy of a particle incident into the detector and the number of secondary particles in the induced shower. This number is roughly proportional to the incident energy. In the so called analogue Monte Carlo, this means computer time need increases with accelerator energy and will cause problems, unless computer efficiency is increased at the same time. Also, the structural complexity of high energy detectors increases with energy and makes event reconstruction more difficult, increasing the need for support by Monte Carlo simulation again /1/.

Several years ago we started a research program at the HLRZ super computing center to develop vectorized Monte Carlo simulation codes for particle interaction with matter for applications in medium and high energy physics, namely simulation of particle shower development in the ZEUS calorimeter and similar detectors.

Particle showers in such large detectors comprise particle energies from the energy of the incident high energy particle down to thermal energies. Due to the high multiplicity of cascade particles with low energy, the Monte Carlo programs for low energy particle transport (mainly neutrons and photons) are the most time consuming ones. Therefore, we developed a vectorized particle transport code comprising a geometry package to allow the description of complex geometrical detector structures, a nuclear event generator using tabulated particle cross sections for all reactions of interest in the simulation of single nuclear events in the CM system, conserving energy and momentum in each collision, and a tracking routine for neutral particles /2,3/.

In a high energy detector normally the energy deposited by the cascade particles in the detector material is the physical quantity of interest, and hence, the fraction that is really detected, e.g. the so

called visible energy in the case of a scintillation detector. In the low energy part of the cascade (≤ 20 MeV) the energy is transported mainly by neutrons, whereas the energy deposition proceeds via stopping of recoils.

Up to now, therefore, energy deposition has been computed as neutron flux multiplied by the kerma factor, which represents the effect of energy deposition by recoiling charged particles. In a detector like ZEUS, however, the ranges of recoil protons are comparable to the scintillator thickness (2.5 mm). This fact leads to a contribution to the signal noise, which cannot be seized in case of using the kerma factor. In the last year we, therefore, developed a vectorized tracking routine for charged particles with an extension for conversion to scintillation light. This latter options use table lookup in an energy range table to determine energy loss of a particle. The usual procedure to determine the rest energy by inverse table lookup was replaced by lookup in an additional range energy table. This was necessary, because we perform the lookups by calculating the address of the appropriate element in a table with logarithmic argument. This procedure can be successfully vectorized, unlike real stepping through the table. In a number of test runs, with energy range tables for carbon and hydrogen we found that tables with about 100 elements between 1 keV and 20 MeV provide range and energy calculation precision much better than required for detector thicknesses in the order of millimeters. The energy range table as material data is read in by the program, which computes from this range the energy table. Also it calculates a visible energy versus deposited energy table to be used with the scintillation light detector. This table is calculated from the energy range table using Birk's law to account for the quenching effect.

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Neutron Spectra Measurements and Monte Carlo Simulations of (p,xn) Reactions at 597 MeV and 800 MeV

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Double differential neutron production cross sections were measured at the LAMPF-WNR facility at Los Alamos between 100 MeV and 800 MeV incident proton energy using time-of-flight techniques. Virtually systematic results were gained by a large number of different angles, energy points and target materials from low to high atomic masses. The measured spectra were compared to detailed Monte Carlo simulations using the KFA version of the intra-nuclear cascade code HETC. The results are also compared with earlier measurements and those of other authors. Here we present the latest series of these measurements, which were performed at 597 MeV and 800 MeV incident proton energy.

The experiments were performed by the LANL-KFA collaboration in the period from 1981 to 1990. Single measurements in this series and in one case comparison with Monte Carlo simulations have been published earlier /1/. In Table 1 an overview about the recent measurements is given.

Table 1. Thin Target (p,xn) Measurements

Energy (MeV)	Target Materials	Angle (degree)
597	Be,B,C,N,O,Al, Fe,Pb,U	30,60,120,150
800	Be,B,C,N,O,Al, Cd,Fe,W,Pb,U	30,60,120,150
800	C,Al,Ni,Ta,W, Pb,U	7.5,30

The intranuclear cascade evaporation model (INCE), as implemented in the high energy radiation transport code HETC of the HERMES system /2,3/, is used in the calculation of double differential cross sections of proton induced neutron production /1/.

To achieve optimum comparison, identical energy intervals as used in the experimental analysis were

provided in the HETC calculations. Depending on target material, emission angle and incident proton energy, the number of spallation events varied from 2.5×10^5 to 6.0×10^5 to achieve reasonable statistics in the calculations.

Some examples out of this recent series of measurements together with the respective Monte Carlo simulations are presented in Figs. 1-2. The kind of presentation (per units of lethargy) was chosen here to pronounce discrepancies between experimental and simulated curves. With these and the earlier measurements proton induced neutron production cross sections are available now for a wide range of target masses, incident energies and neutron emission angles. Discrepancies observed between earlier measurements and INCE simulations, become considerably smaller with the latest measurements. In Fig. 3 a comparison is made with older data of Ref. /4/. The disagreement between measurements and HETC-INCE simulations is clearly larger in the older measurements, than in the most recent ones. This is particularly true for the backward angles. Details are given in Ref. /5,6/.

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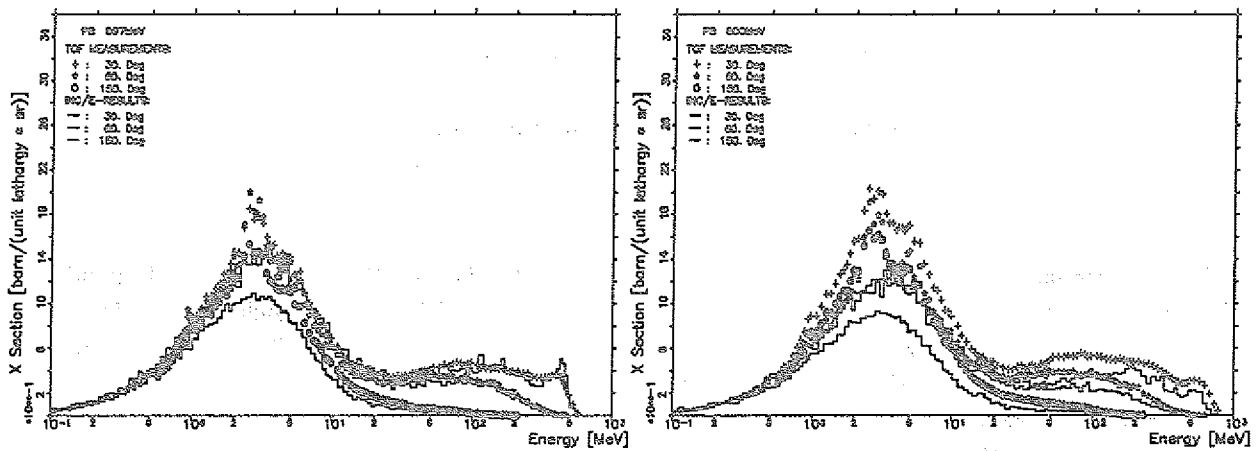


Fig. 1 Neutrons from 597 MeV protons (left) and 800 MeV protons (right) on lead

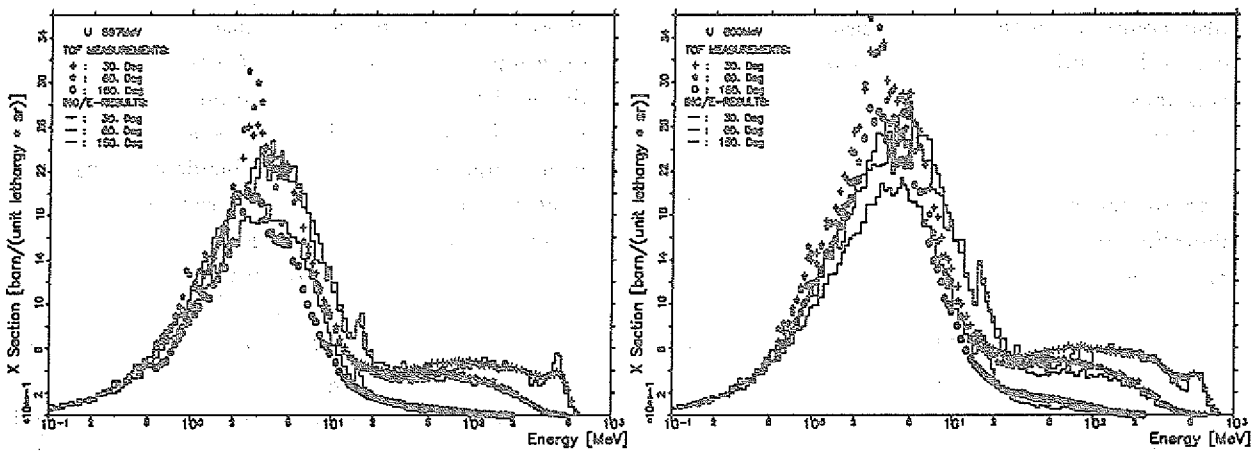


Fig. 2 Neutrons from 597 MeV protons (left) and 800 MeV protons (right) on uranium

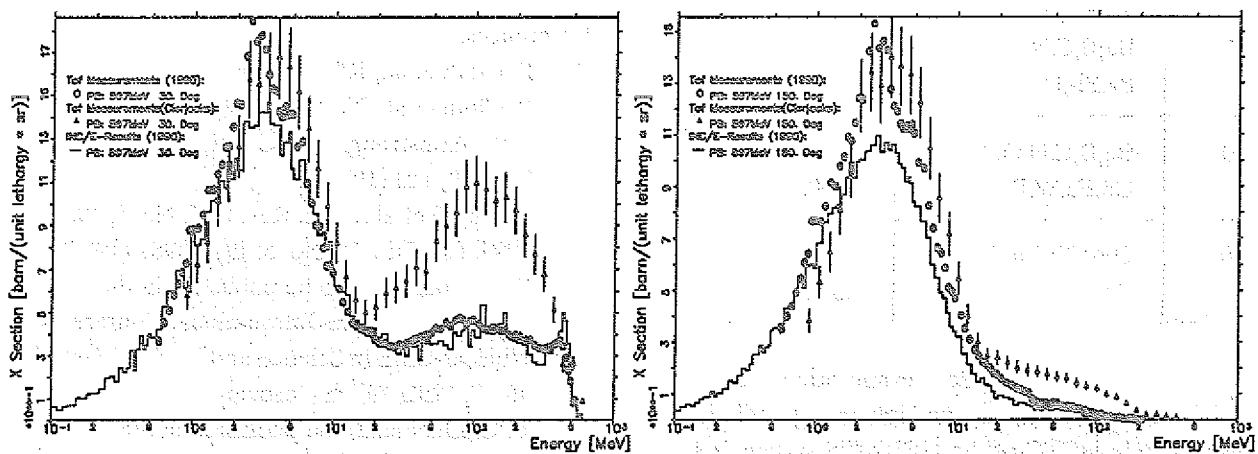


Fig. 3 Example for a comparison with older data of Ref. /4/ - Neutrons from 597 MeV protons on lead - (30-degree on left, 150-degree on right)

Simulation Experiments for Planetary Gamma-Ray Spectroscopy by Means of Thick Target High-Energy Proton Irradiations

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The scientific objective of this experiment was to simulate the emission of planetary gamma rays that are leaking from a planet's surface as a result of the interaction of the galactic cosmic rays with the planet. A detailed investigation was needed for the upcoming NASA Mars Observer Mission (launch September 1992). The Gamma-Ray Spectrometer (GRS) on board the spacecraft was designed to measure the intensities of gamma-ray lines, characteristic of each element detected, that emerge from the planetary surface [1]. The inversion of the measured gamma-ray spectra will provide data on the elemental composition of the surface of Mars. These data depend strongly on the quality of the applied inversion model, which was to be tested by an accelerator experiment.

Very recently, we simulated the process of planetary gamma-ray production by irradiating a so-called "Thick Target" with energetic protons at the French accelerator Saturne. The gamma-ray flux emitted from the Thick Target was measured by a high-resolution gamma-ray spectrometer consisting of a high-purity germanium detector (HPGe), a charged particle anti-coincidence shield, and complex electronics.

The large-volume n-type HPGe detector, a special design by PGT, was equipped with a transistor reset preamplifier to avoid saturation resulting from interactions of energetic protons and gamma rays with the germanium crystal (for example, a 1.5 GeV proton deposits about 40 MeV in a 5 cm long Ge crystal). In a typical run, about 950 protons and 12200 gamma rays per second were detected, while the reset rate was about 65 per sec. The efficiency of the HPGe detector was carefully determined in an energy range from 0.1 to 10 MeV over a large field of view by using a rotational device to position calibrated gamma-ray sources at any angle between ± 80 degree azimuth and declination.

Since a HPGe detector is sensitive to gamma rays, charged particles and to less extent to neutrons, energetic protons and pions had to be screened out to avoid interferences with the gamma-ray measurements. Therefore, the detector was surrounded by a plastic scintillator anti-coincidence

shield that vetoed events resulting from charged particles. A sophisticated electronic circuit determined the total system dead time to correct for count rate losses.

The electronic signals resulting from the gamma rays were routed in two electronic units with respect to alternate beam on and off. In addition, the general gamma-ray background produced by surrounding material in the experimental hall was determined by placing remotely a lead shield in front of the Thick Target having a hole to allow the proton beam to hit the target. Four types of gamma-ray spectra were accumulated sequentially: target spectra taken during proton beam on and off, then background spectra taken during beam on and off for the same amount of time. These measurements were performed periodically.

To simulate the Martian surface the target used was large enough to stop the primary protons and to confine the secondary particles produced as a result of nuclear cascade processes with neutrons as the most abundant species. The Thick Target had the outer dimensions of 1.5x1.5x1.8 meters and weighted about 30 metric tons when filled with iron. The Thick Target was designed to satisfy versatility of material content and safety requirements for crane operation. It was divided into three major parts: the surface target, the inner core, and the outer sleeve. The surface target contained the materials to be studied; it faced the incoming proton beam and was the major source of prompt gamma-rays. The purpose of the inner core was to expose the primary proton beam to the same material along its total stopping range. The outer sleeve contained enough iron to confine most of the secondary particles produced inside the Thick Target. Inner core and surface target could be rapidly exchanged after having removed the upper part of the iron sleeve.

Two different types of material were used for bombardment: iron and basalt. The iron target was the pure water-free case and simple in composition, while the basaltic rock contained major and many minor rock forming elements including water that effects neutron moderation, significantly.

During the initial process of beam tuning the spectrometer was removed from the experimental area to prevent radiation damage /2/. Two large plastic scintillators with different design were positioned in the main beam sustaining high proton count rates (typically 1.5 million protons per spill), while five small scintillation detectors were placed at various distances from the beam line. These detectors were used for the fine tuning of the pencil shaped main beam and the surrounding halo that had to be minimized. The first proton beam had an energy of 2.5 GeV, a beam repetition rate of 3.5 sec, and a spill time of 0.5 sec. After beam tuning the gamma-ray spectrometer was positioned at a distance of 50 cm in front of the Thick Target with an off-set to the beam axis of 25 cm.

The first target material, a basaltic rock (gabbro), was bombarded by protons for about 21 hours. At a designated time proton energy was changed to

1.5 GeV with a new beam repetition rate of 1.6 sec and an identical spill time of 0.5 sec. The beam was to be retuned at the beam line, which required storage of the HPGe detector in a radiation safe position again. Following this, bombardment of the same target configuration continued for about 24 hours.

In a fast craning process, except for the outer sleeve, the surface and the inner core were replaced with iron and measurements were carried out for 40 hours in a similar manner as the rock target. A gamma-ray spectrum (Fig. 1) taken for about 4 hours during beam on shows the most prominent lines of iron ($\text{Fe} = \text{capture reaction}$, $\text{Fe}^* = \text{scattering reaction}$), silicon, hydrogen, germanium (triangular shaped), and annihilation /3/. The lines are sitting on a relatively high gamma-ray continuum as expected.

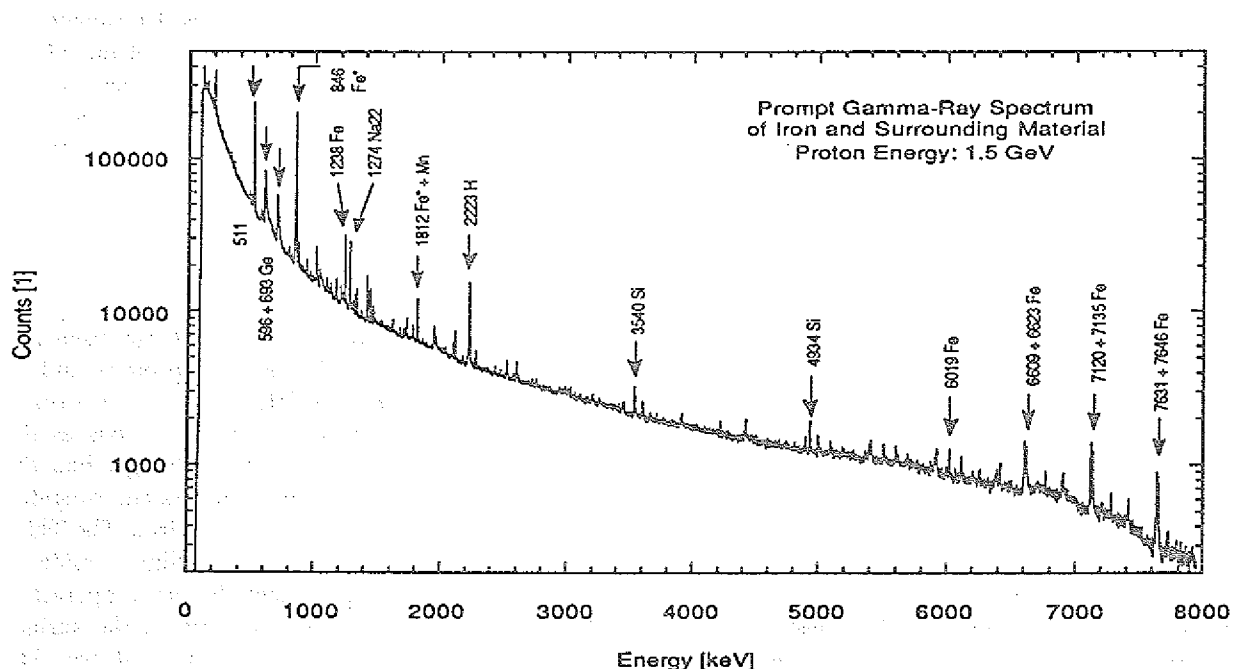


Fig. 1: Prompt gamma ray spectrum from 1.5 GeV protons on iron target

During most of the time neutrons were also measured with a small spectrometer containing two helium and BF_3 filled proportional counters with one counter wrapped in a thin cadmium foil. In this way the neutron flux was determined in the thermal and epithermal energy range.

Over the whole period of measurements the accelerator provided an impressively stable proton beam that permitted the acquisition of excellent data.

Simulations calculations of the complete "Thick Target - HPGe detector" setup are presently underway by using the HERMES code system /4/.

The calculated production rate of discrete-energy gamma-ray lines will be compared with the measured data in order to verify the accuracy of the inversion model.

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A Meson Production Experiment in Uppsala

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The large cross section for η production in the reaction $pd \rightarrow {}^3\text{He} \eta$ close to threshold that has been found in measurements at SATURNE [1], has recently increased the interest in the η channel of light ion collisions. Of central interest apart from the interaction of the η with nuclei resp. nucleons is the production mechanism itself. Several attempts have been made to explain the fact that the η cross section, measured at backward angles, is comparable to the π^0 cross section at the same momentum transfer despite of the difference in their mass. The elementary production process near threshold is presently thought to be dominated by the S_{11} nucleon resonance $N^*(1535)$,

$$\pi + N \rightarrow N^*(1535) \rightarrow N + \eta,$$

with exchange of a virtual meson (π, ρ, η) between two nucleons in analogy to pion production. On the contrary to pion production, the non-resonant contribution to the η production is negligible due

to the weak $NN\eta$ coupling [2]. But the theoretical cross sections calculated on the basis of this production process are still too small. This indicates that possibly collective effects and 3-nucleon mechanisms must also be taken into consideration, as e.g. double scattering of the intermediate mesons [3].

To study the production of η 's and their interaction with nuclear matter in more detail, the mesonic reaction channels in (p,p) and (p,d) interactions, suggested by the PROMICE group [4], will be measured at the cooler-storage ring CELSIUS in Uppsala by the PROMICE/WASA collaboration. The high momentum resolution ($\Delta p/p \sim 2 \cdot 10^{-4}$) of the cooled circulating beam in this storage ring is ideally suited for particle production near threshold and studies of rare processes. An internal cluster jet will be used as ultra thin target ($\sim 10^{-9}$ g/cm²). The experimental setup is shown in figure 1. The electromagnetic calorimeter consists of two 7×8 matrices of 30 cm long CsI(Na) crystals (16 X_0)

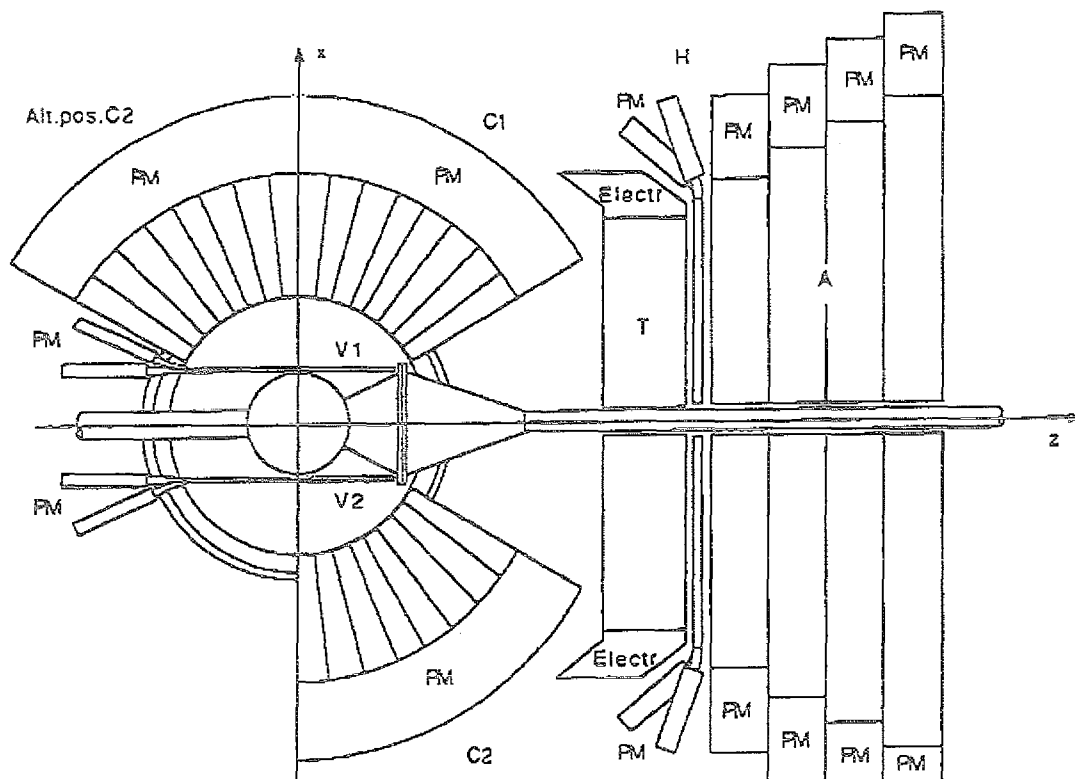


Figure 1: Top view of the PROMICE experimental setup. Indicated are the forward spectrometer with the tracker (T), the hodoscope (H) and the range telescope (A). The wide angle electromagnetic calorimeters (C1,C2) are shown with an alternative location for the C2 calorimeter. V1 and V2 are plastic scintillators covering the front of the calorimeters.

* PROduction of Mesons In CELsius

** Wide Angle Shower Apparatus

forming two spherical sectors covering about $60^\circ \times 60^\circ$.

The forward spectrometer has three parts: A tracker consisting of four blocks of thin walled ($10^{-4} X_0$) straw chambers, a range telescope with four layers of thick plastic scintillators, and between them a three layer circular plastic scintillator hodoscope for triggering purposes, track reconstruction, and ΔE measurements. This hodoscope is of the same type as used in the JETSET experiment at CERN and as under construction for the time of flight spectrometer at COSY Jülich. The three layers consist of 5 mm thick plastic scintillators with a radius of 58 cm: 48 straight triangle shaped in the first plane, 24 left and 24 right twisted scintillators in the second and third plane. A coincidence between signals from two or all three planes defines a pixel being the geometrical overlap area of the corresponding scintillator modules.

For the total cross section about 2 η 's per second and μb will be detected at a luminosity of $3 \cdot 10^{30} \text{ cm}^{-2} \text{ s}^{-1}$. Using kinematical constraints, the final event sample can be made practically background free. An angular resolution of 1° (fwhm) is provided by the scintillator hodoscope in scattering angle and 4° (fwhm) in azimuthal angle, which is adequate to distinguish between several charged tracks, so that the hodoscope can be used to resolve multi-hit ambiguities and to trigger on a particular hit multiplicity. In the reaction $pd \rightarrow {}^3\text{He} \eta$ a trigger of multiplicity 1 is demanded, corresponding to the recoiling ${}^3\text{He}$ nucleus.

Neutral pions which will be the dominant background will be measured via their decay photons which are detected in the electromagnetic calorimeter unless the ${}^3\text{He}$ recoils can be observed at small angles. The total detection efficiency for neutral pions is about 1 %.

During 1991, data have been accumulated using the H_2 cluster target and proton beams of 500 and 1150 MeV with intensities of $1\text{--}3 \cdot 10^{10}$ protons.

Furthermore, short test run was performed in the end of October with 275 MeV protons and electron cooling applied during part of the time. Finally we took data during the last week of November at 500 MeV (p+p) for calibration, with "slow ramping" from 270 to 310 MeV over the $pp \rightarrow pp\pi^0$ threshold, and a second flat top ramping run over the $pd \rightarrow {}^3\text{He} \pi^0$ and $pd \rightarrow pd\pi^0$ thresholds (180 — 240 MeV).

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Is there Shape Coexistence in ^{99}Zr ?

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The sudden onset of ground-state deformation at $N=60$ in the region of neutron-rich $A \sim 100$ nuclei is connected with the occurrence of shape coexistence in this mass region. Prominent examples¹⁻⁵⁾ are the $N=59$ nuclei ^{97}Sr and ^{98}Y where spherical ground states coexist with deformed structures at excitation energies around 0.5 MeV. Evidence for a collective E2 transition has also been found⁶⁾ in their isotone ^{99}Zr . At the time of its discovery this transition could not be interpreted since very little was known about rotational structures in odd-mass nuclei of this region of deformation. With more detailed knowledge about bands in these nuclei and their interpretation in terms of rotational models which is now available, a new attempt to understand the structure of ^{99}Zr has been made. These activities include new measurements on the γ -ray intensities and level lifetimes at JOSEF as well as particle-rotor coupling calculations.

Fig. 1 shows properties of ^{97}Sr and ^{99}Zr which are relevant in this context. These isotones are characterized^{7,8)} through a set of three low lying states and a bunching of levels around 500 keV. A well developed rotational band is built on the 585 keV level in ^{97}Sr . The energy differences between the three band members are close to those of ground-state bands in this region, and large quadrupole moments have been determined through the measurement of the level lifetimes^{4,9)}. A second band with a collective intraband E2 transition but with a very small level spacing seems⁵⁾ to be built on the 645 keV level. The three lowest levels of ^{97}Sr are interpreted as spherical structures and their lifetimes are well reproduced through IBFM⁵⁾ calculations. The ground state and the 308 keV state have mainly $s_{1/2}$ and $g_{7/2}$ single particle character, respectively, while the 167 keV level is rather mixed with strong $d_{3/2}$ single particle and $g_{7/2} \otimes 2^+$ components.

No spin and parity measurements have been possible for ^{99}Zr but the similarity of the three lowest-lying levels to those of ^{97}Sr is striking. If the reduced transition probabilities from the IBFM calculations for ^{97}Sr are scaled to ^{99}Zr then lifetimes of 0.7 and 273 ns result for the first and second excited states of ^{99}Zr , respectively, in accordance with the measured values, see Fig. 1. Therefore it is suggestive to assume that the three lowest levels of ^{99}Zr have the same nature as those in ^{97}Sr . This implies spherical nature of ^{99}Zr in its ground state.

Evidence for the existence of deformed states in ^{99}Zr

is provided through the half-life⁶⁾ of the 667 keV level of 8.7(5) ns. This level is depopulated mainly via the low-energy transition of 53.3 keV. The conversion coefficient of $\alpha = 2.0(6)$ identifies the 53.3 keV γ ray as an M1/E2 transition with a mixing ratio of $|\delta| = 0.41(15)$. The deduced transition probabilities amount to $3.4(18) \cdot 10^{-3}$ and 183(103) single particle units (spu) for the M1 and E2 fractions, respectively. The large uncertainties are mainly due to the uncertainties of the mixing ratio of the 53.3 keV transition. The E2 transition probability corresponds to a quadrupole moment of $|Q_0| = 3.9(11)b$ and to a value of the deformation parameter of $\beta_q = 0.37(10)$ if $K=3/2$ or $5/2$ and prolate shape are assumed.

The energy of 53.3 keV is much lower than that of unperturbed intraband transitions in this mass region. In order to see whether this can be explained through band mixing, particle-rotor coupling calculations have been performed. The same parameters were used as in the case of $^{103}\text{Mo}_{51}$ which has successfully been described¹⁰⁾ in terms of this model. The calculations have been performed in the range of $\beta_q = 0.37 \pm 0.10$. It was found that the $[541]3/2^-$ assignment to the 614 keV level (with probable spin of $3/2$ or $5/2$) which would correspond to the interpretation of the 645 keV level in ^{97}Sr , poses problems since this band head lies high in energy, see Fig. 2. Moreover, the experimental values for $t_{1/2}(667 \text{ keV})$ and for $|\delta(53.3 \text{ keV})|$ are not reproduced in the calculations. A better fit is obtained if the main configuration $[422]3/2^+$ is assigned to the 614 keV level. A strong mixture with the $[411]3/2^+$ configuration affects the properties of this band. The results are given in Fig. 2. For the partial half-life of the 53.3 keV transition a value of 7.1 ns is found and the mixing ratio amounts to 0.52, compared to the measured values of 10.4(43) ns and 0.41, respectively. The energy of the 53.3 keV transition is not so well reproduced, a finding which might be improved through fine tuning of the parameters.

In summary, there is evidence for the occurrence of shape isomerism in ^{99}Zr as in its immediate neighbours ^{97}Sr and ^{98}Y . The available data are, however, not sufficient yet to draw firm conclusions. Further measurements are needed, especially concerning the relative γ -ray intensities and the multipolarities of the 53.3 and 92.0 keV transitions in order to allow a determination of the multipolarity of the latter and, thus, a more precise value of the enhancement of the E2 fraction of the former. With fission product intensities which are higher than those available at

JOSEF it might be possible to identify further low-energy transitions among the group of levels above 500 keV, the main candidates for deformed structures. The new half-lives of the 575 and 852 keV levels which have been determined at JOSEF (see Fig. 1), will be very useful for an interpretation of this nucleus. Thus, if either the 194.4 or the 276.4 keV line turns out to be a cross-over transition in a rotational band with E2 nature, then it has an enhancement of more than 100 spu.

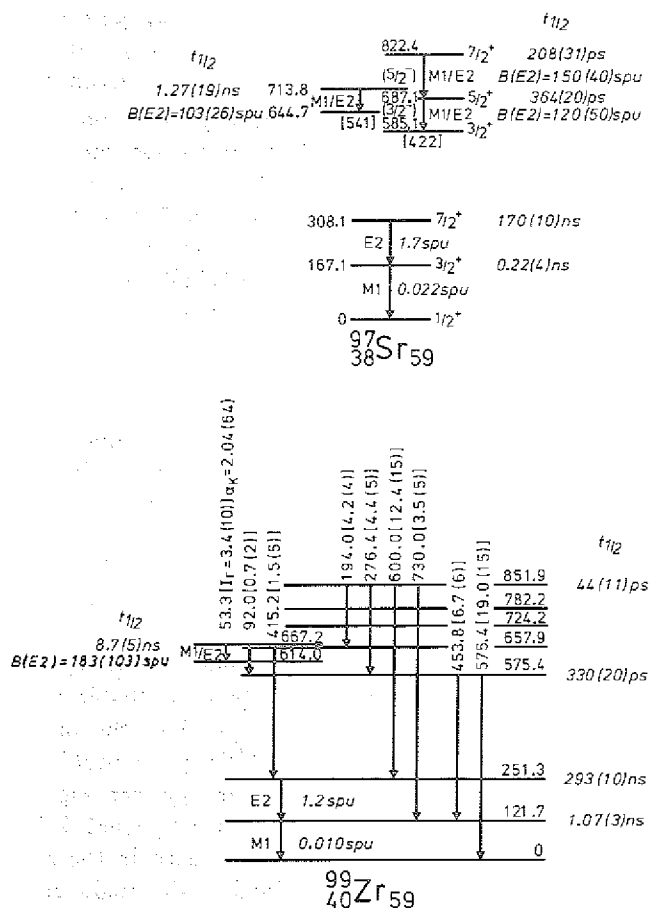


Fig. 1. Partial level schemes⁵⁻⁸⁾ of ^{97}Sr and ^{99}Zr . The transition of 92.0 keV in ^{99}Zr has been identified¹¹⁾ at the fission product separator OSTIS. The half-lives of the levels at 122, 575 and 852 keV result from recent β - γ -time studies¹²⁾ at JOSEF. The relative γ -ray intensities for ^{99}Zr are weighted averages of published data^{7,13)} and of the results from additional γ - γ coincidence measurements at JOSEF.

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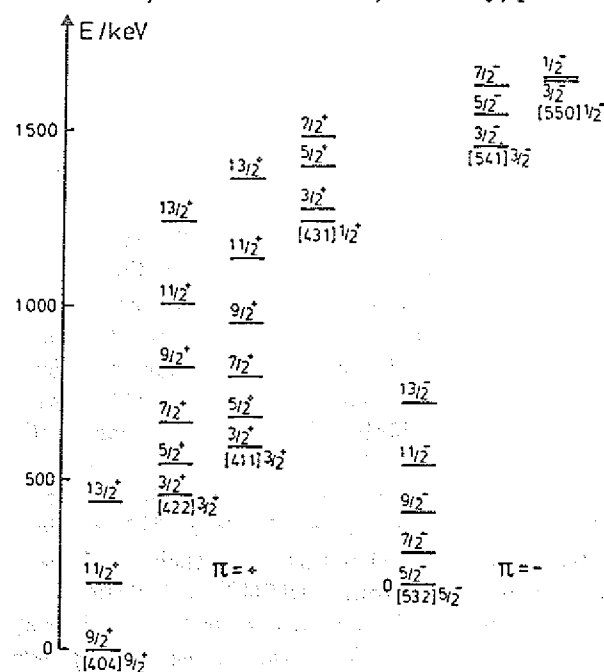


Fig. 2. Low-spin and low-energy results of particle-rotor calculations for ^{99}Zr with the Nilsson and pairing parameters from the studies¹⁰⁾ of ^{103}Mo . Further parameters are: Coriolis coupling attenuation factor = 0.9; core energy = 144.6 keV from $E_2+(^{98}\text{Sr})$.

Nilsson Parameters at $N=61$ and Particle-Rotor Calculations for ^{103}Mo

M. Liang, H. Ohm, I. Ragnarsson* and K. Sistemich

Lifetimes of levels of neutron-rich deformed nuclei around $A=100$ have been measured at the fission-product separator JOSEF with the new β - γ - γ triple coincidence technique for studies in the ps region [1,2]. The results are crucial input information for tests of the applicability of the classical models for rotational nuclei in this new region of deformations.

In ^{103}Mo the lifetimes of five levels have been determined, see [3] and Fig. 1. The deformation has been deduced as $\beta_q = 0.34(1)$ from the half-life of the first excited state. Thus, this nucleus is well deformed, and the published experimental information on the energies and branching ratios of the ground-state band [4-6] together with the new lifetime information provides a good basis for theoretical studies. Therefore, particle-rotor coupling calculations have been performed for this isotope.

A first step of these investigations was the determination of the model parameters, in particular of the Nilsson parameters κ and μ . It was found that the proposed sets (A) $\kappa = 0.080$, $\mu = 0.22$ [7] and (B) $\kappa = 0.066$, $\mu = 0.36$ [5,8] did not lead to a satisfactory description of the data: the former could explain the half-life of the first excited state but did not reproduce the large staggering of the ground state band built on the $[411]3/2$ configuration [5,6], see Fig. 2. The other set did well for the staggering but badly for the half-life.

Therefore optimized parameters have been deduced for ^{103}Mo . This optimization was achieved through a fit of the results of the calculations to the energies in the ground-state band, the half-life of the first excited state, the expectation value $\langle s_z \rangle$ of the spin projection along the nuclear symmetry axis which can be deduced from the ratio $(g_K - g_R)/Q_0 = -0.171(8)$ with $Q_0 = 3.84(18)\text{b}$ [3]. Additionally, it has been assumed that the 641 keV level has the configuration $[422]3/2^+$, a proposition based on the low $\log ft$ value of 4.9 for its β -population from the $[422]5/2^+$ ground state [9] of ^{103}Nb . The deduced parameter set is

$$\begin{aligned}\kappa &= 0.084 \\ \mu &= 0.28 \\ f(g_s) &= 0.85 \\ \text{ccrf} &= 0.85\end{aligned}$$

where $f(g_s) = g_s^{\text{eff}}/g_s^{\text{free}}$ and $\text{ccrf} = \text{Coriolis-coupling attenuation factor}$. Standard values were used for the pairing strength parameter G and for $g_R (= Z/A)$.

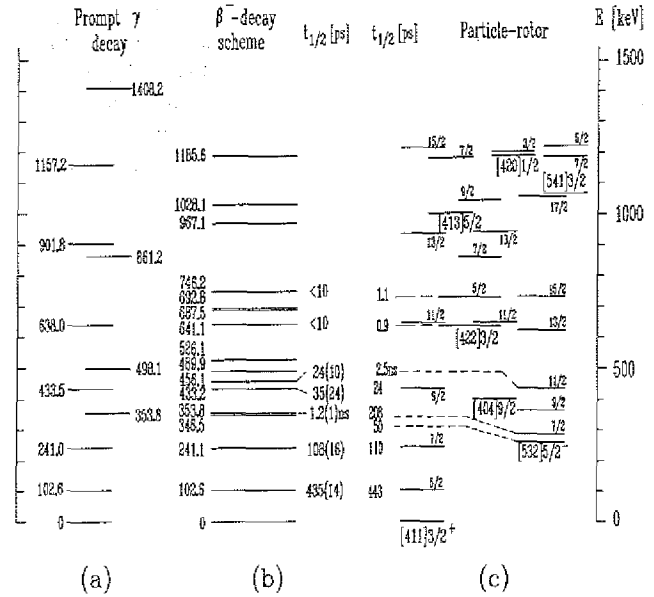


Fig. 1: ^{103}Mo level scheme as populated in prompt γ -decay [8] of fission products (a) and in the β -decay of ^{103}Nb [4] (b). Part (c): band structures obtained in particle-rotor model calculation with adjusted parameters, see text.

g.s.b. in $\text{Mo}103$

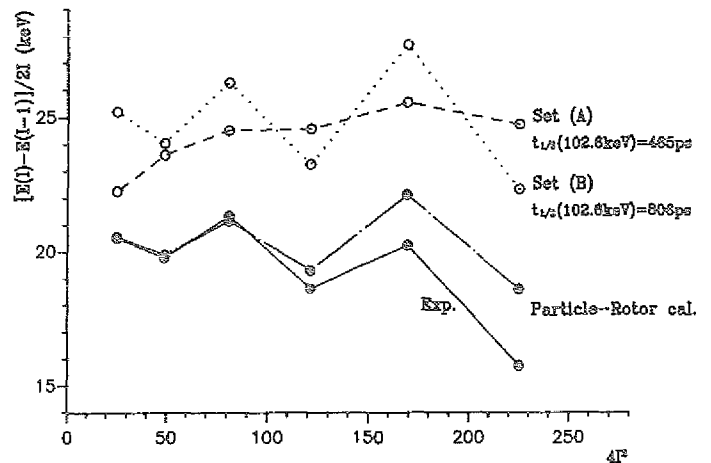


Fig. 2: Energy staggering of the g.s. band in ^{103}Mo . The solid line represents the experimental data. The short-dashed, dotted and dot-dashed lines show the results of particle-rotor model calculations with the parameter sets (A), (B) and with the adjusted parameters, respectively. The curves for set (A) and (B) have been shifted up by 5 keV.

The results which are obtained with these parameters and with the experimental value of the deformation parameter are compared in Fig. 1 to all available data. Good agreement is achieved. An interesting aspect is the fact that the calculations predict the position of the $[532]5/2^-$ band head at about 270 keV and that they result in a strong perturbation of this band and a half-life of about 200 ps of the $7/2^-$ state. It is suggestive to correlate this band [6] with the experimental levels at 347 and 354 keV, where the latter has a half-life of 1.2(1)ns. Since this level is depopulated through γ rays into the first and second excited states of the ground-state band with the probable multipolarity of E1, the ratio of 6 between the measured and calculated value for the half-life is not to be considered as a serious discrepancy.

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Mass-, Charge-, and Angular Momentum Transfer
in $^{106}\text{Cd} + ^{54}\text{Fe}$ Collisions studied by $\gamma\gamma$ -Coincidences

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Recent results obtained with anti-Compton shielded multidetector arrays demonstrate the excellent resolving power of the $\gamma\gamma$ coincidence techniques. The high-quality coincidence data obtained with such arrays allow to study excitations of individual nuclei, that are products of as complex processes as nuclear fission¹⁾. Moreover, it has been shown²⁾ recently, that one can also observe cross-coincidences between γ -rays emitted from two nuclei — the binary reaction products — hereby identifying their primary or secondary character.

In order to explore further the potential application of this technique in heavy-ion reaction studies we performed $\gamma\gamma$ coincidence experiments for the $^{106}\text{Cd} + ^{54}\text{Fe}$ system. The 255 MeV ^{54}Fe beam energy, approximately 30 MeV above the Coulomb barrier, was chosen to enable observation of a wide

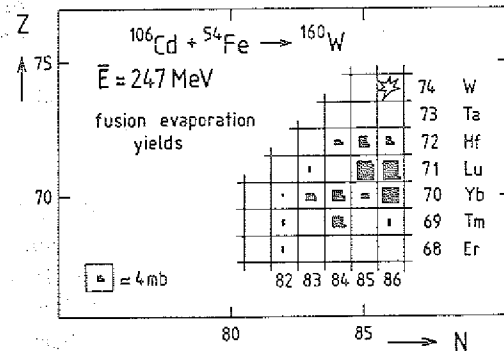


Fig. 1: Production cross sections of evaporation residues for 247 MeV ^{54}Fe on ^{106}Cd .

variety of nuclear processes. A thick target of 1.2 mg/cm^2 90.7% enriched in ^{106}Cd was placed on a 14 mg/cm^2 ^{208}Pb backing to ensure the stopping of virtually all reaction products. The ^{54}Fe beam from

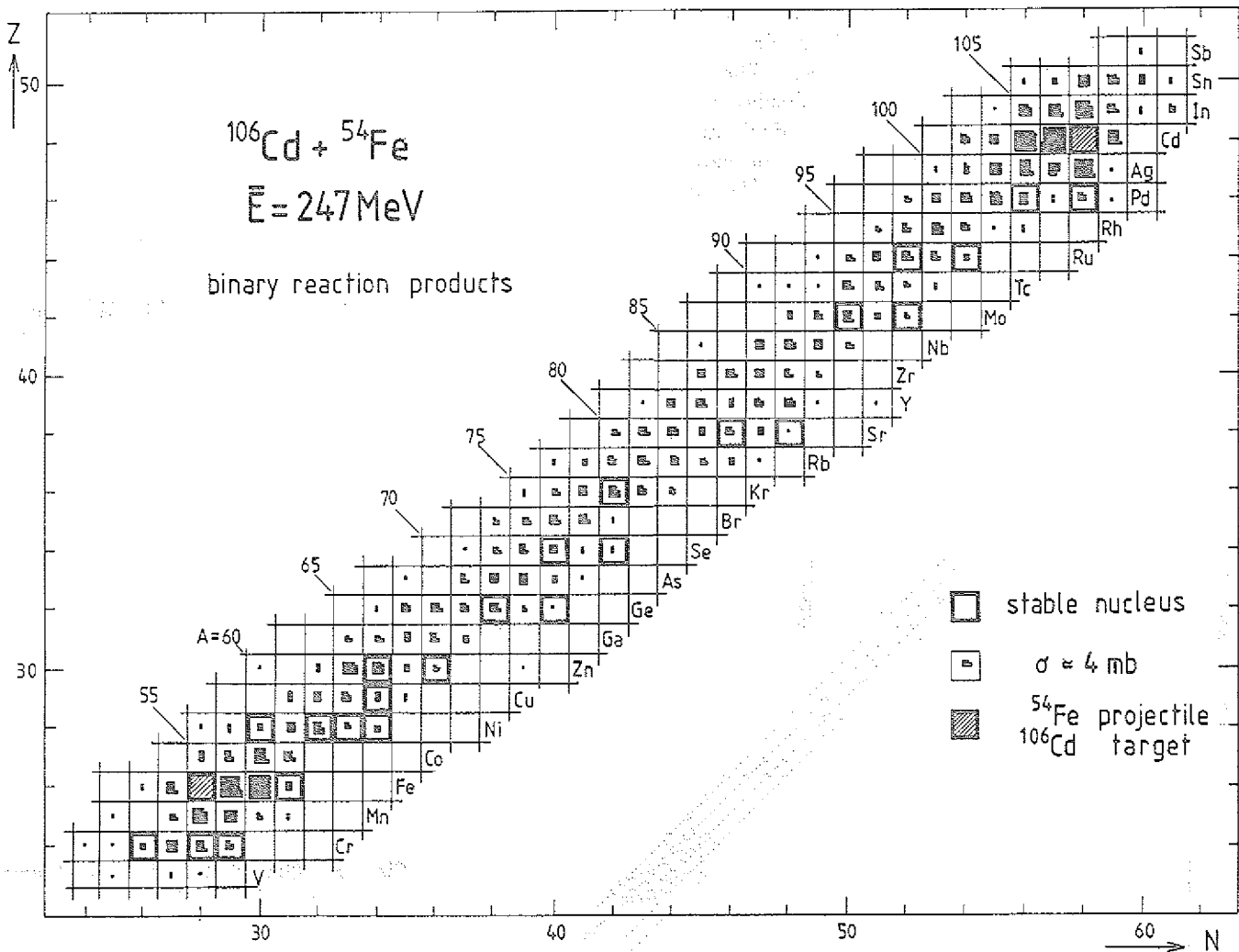


Fig. 2: Production cross sections of binary reaction products for 247 MeV ^{54}Fe on ^{106}Cd .

the VICKSI accelerator at the HMI Berlin was pulsed with 300 ns beam pulse separation, which allowed to separate the prompt in-beam and off-beam (isomeric and radioactive decay) events. The OSIRIS multi-detector ring was oriented perpendicular to the beam direction, and a large 4-sector neutron detector was placed at forward angles. Gamma-X and γ -n coincidences were essential for identification of several higher hitherto unknown $N=84,85$ nuclei produced in fusion evaporation, which are a subject of detailed spectroscopic analysis.

After the experimental run the radioactivity collected in the target was measured and its decay was followed for a period of four months. This data provided accurate production yields of many long-lived products. Quantitative analysis of the off beam coincidence events (between the beam bursts) gave relative yields for short lived radioactive products including $T_{1/2} > 20$ ns isomeric states. More laborious analysis of the prompt in-beam coincidence data confirmed the direct population of nearly 200 nuclei.

For most of them sequences of several coincident, predominantly yrast γ -transitions known from previous studies, could be identified. Some of these in-beam coincidence results were also used to estimate the production yields of stable nuclei and of those with unfavorable radioactive decays. The Figs. 1 and 2 give an approximate graphical display of the production cross sections of the individual nuclei produced through fusion-evaporation or in various binary reactions taking place in $^{106}\text{Cd} + ^{54}\text{Fe}$ collisions; in an extended version of this report numerical

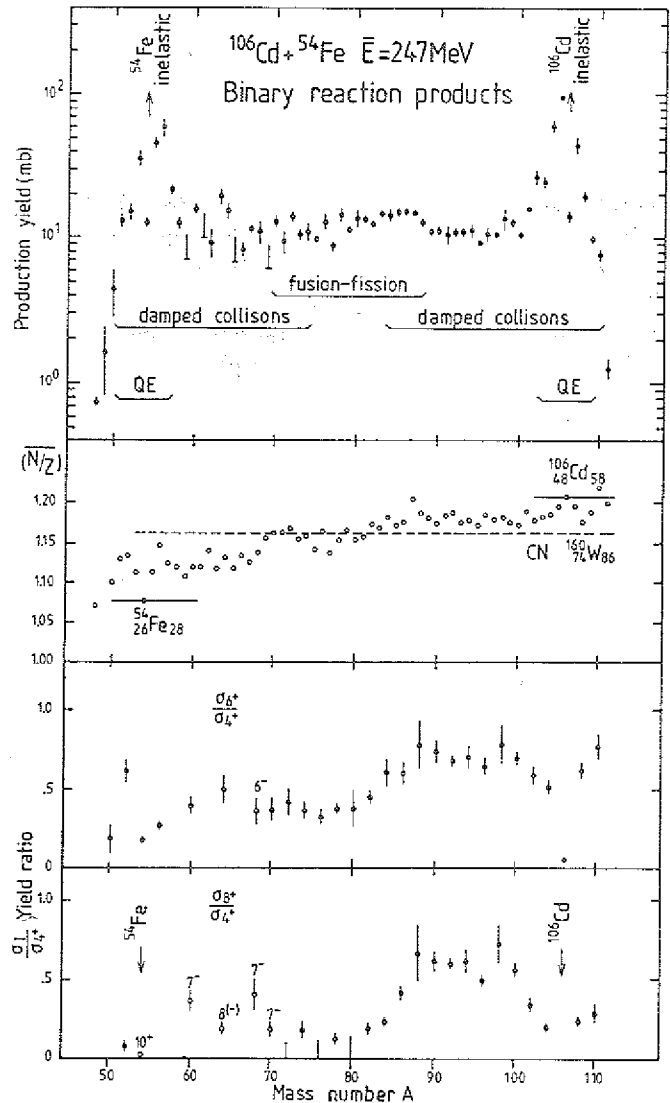


Fig. 4: Production yields, N/Z ratios and direct population of 6^+ and 8^+ levels relative to 4^+ as function of mass

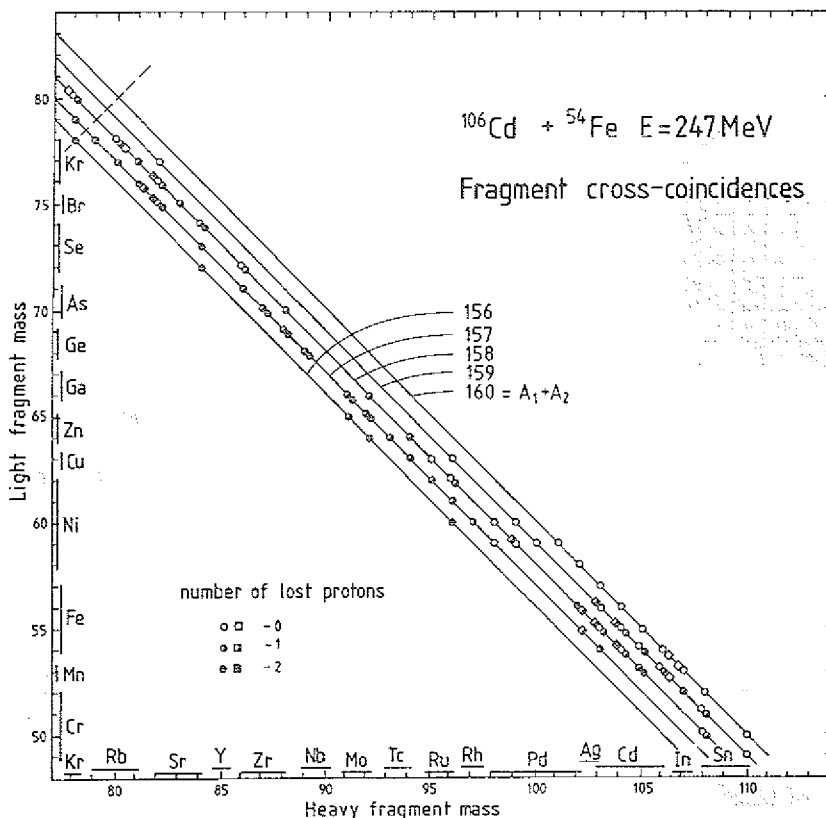


Fig. 3: Observed pairs of binary reaction products.

values will be given. Absolute cross-section normalization was obtained using the results of earlier measurements³⁾ of the α -decaying fusion evaporation products performed for the same reaction with the GSI SHIP facility. The cross-sections for the fusion evaporation reaction channels agree well with HIVAP calculations³⁾. The summation of yields for the respective individual nuclei indicates that the fusion-evaporation process represents only 20% of the total reaction cross-section.

From in-beam coincidence spectra with gates set on specific transitions the population of known states in a given nucleus could be determined. In those spectra one could also recognize the γ -lines from the reaction partner nuclei produced in various binary reactions. Usually more than one partner was observed, as expected for data, which are integrated over the kinetic energy loss. The correlation of light and heavy reaction products established from the observed γ - γ cross-coincidences is summarized in Fig. 3. The points in the figure, independent of the detected coincidence intensity, mark all cases, where the data allowed to conclude firmly the simultaneous appearance of two final products. The location of points on the A_1+A_2 lines indicates how many neutrons and protons were lost (most likely evaporated from the excited primary products) in a process represented by each point. One notices that also in the mass regions around the ^{106}Cd target and ^{54}Fe beam nuclei many processes happen with subsequent emission of nucleons. For processes with larger mass transfer, when the heavy fragment mass is below $A=98$, the points scatter around the $A_1+A_2=157,158$ lines. Taking all of them with equal weight one concludes that on average 2.6 nucleons were emitted, of which 1.2 were protons.

Fig. 4 displays the integral results related to mass, charge and angular momentum transfer. The upper window shows the mass distribution, i.e. the projection of the cross-sections of Fig. 2 on the mass axis. The regions, where the different types of reaction processes are expected to occur are indicated. The average N/Z shown below was calculated for each mass using the measured yields of all observed isobars. It is noteworthy that full charge equilibration, i.e. the compound nucleus N/Z value, was only reached in the region around symmetric fragmentation, apparently corresponding to fission following compound nucleus formation. Outside this region the data seem to indicate that only partial, but a nearly constant degree of charge equilibration was reached. The two bottom figures display the population of 6^+ and 8^+ states relative to the 4^+ state population obtained from coincidences observed for even-even products throughout the full mass spectrum. It is clear that the highest spin population is obtained for heavy fragments produced in damped reactions while fission product nuclei are populated at much lower spin values. Further data analysis still in progress promises more detailed results particularly on the angular momentum transfer.

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Study of the Decay out of a Superdeformed Band in ^{143}Eu

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The connection between the super-deformed (SD) and the normal-deformed (ND) states has not been observed yet owing its difficulty to a statistical, highly fragmented decay path out of the SD bands. This leaves both the spin and the excitation energy of the SD bands undetermined. In the present contribution we are presenting results of a new experimental approach where we take advantage of triple and higher fold coincidences to study the decay out of a SD band.

In order to test this method, we carried out an experiment at the Niels Bohr Institute Tandem Accelerator Laboratory where we used the $(160 \text{ MeV}) \text{ } ^{37}\text{Cl} + ^{110}\text{Pd} \rightarrow ^{143}\text{Eu} + 4n$ reaction. The NORDBALL detector array with 19 Compton-suppressed Ge detectors and 1 LEP detector with the BaF_2 ball for multiplicity and sum energy selection were employed. About 10^9 triple and higher fold coincidence events were collected during the experiment, 10 % of them being 4-fold Ge coincidences.

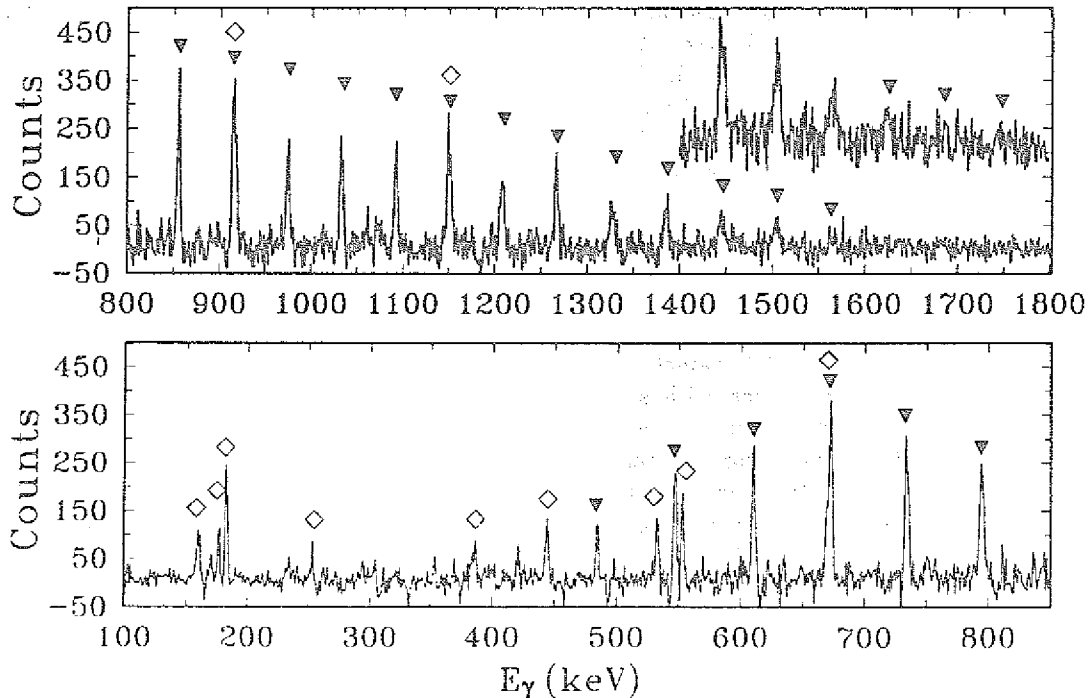


Figure 1: Gamma-ray spectrum coincident with two SD band members. It is a sum of a few clean gates. The band members with energies 484, 547, 609, 672, 733, 794, 854, 914, 973, 1032, 1091, 1150, 1208, 1266, 1325, 1385, 1444, 1503, 1564, 1623, 1686, 1744 keV are marked by (∇) and known low-lying states in ^{143}Eu are denoted by ($\diamond$). The spectrum is shown with a background subtraction and it is corrected for the detector efficiencies.

It is well known from previous works [1] that the decay out of a SD band starts from one or two of the lowest lying known states of the band. After a few consecutive transitions all the decay routes end in a few ND yrast states. We assume that cascades of only two transitions deexcite the SD band through many different intermediate states and they sum up to well defined energies. Using triple coincidence events we sum the energies of two γ rays which are in coincidence with a transition in the SD band to produce a sum spectrum where we would like to identify discrete transitions related to the decay out of the SD band.

Figure 1 shows a spectrum which is double gated on the transitions of the SD band. The transitions between the known [2] low lying states in ^{143}Eu which are populated by the SD band are indicated by an open diamond in the figure. The intensity of the band is about 1 % relative to the total population of ^{143}Eu . This SD band has previously been reported [3] and assigned to ^{142}Eu . In order to clear up the assignment problem, we put one gate on the SD band and another gate on a lower lying transition in ^{142}Eu and in ^{143}Eu , respectively. The SD band can be seen in the latter case, but not in

the former one, leaving little doubt that the SD band belongs to ^{143}Eu . The insert of fig. 1 marks two transitions which are assigned to be the next higher transitions in the band.

The sum spectrum that is the sum of 2 γ -ray energies which are in coincidence with one of the SD transitions is shown in fig. 2.

The peaks marked by arrows in the figure are the candidates which may be assigned to be the sum of linking transitions. The preliminary analysis indicate that the linking is fragmented through many routes, each being very weakly populated. Further data analysis is in progress.

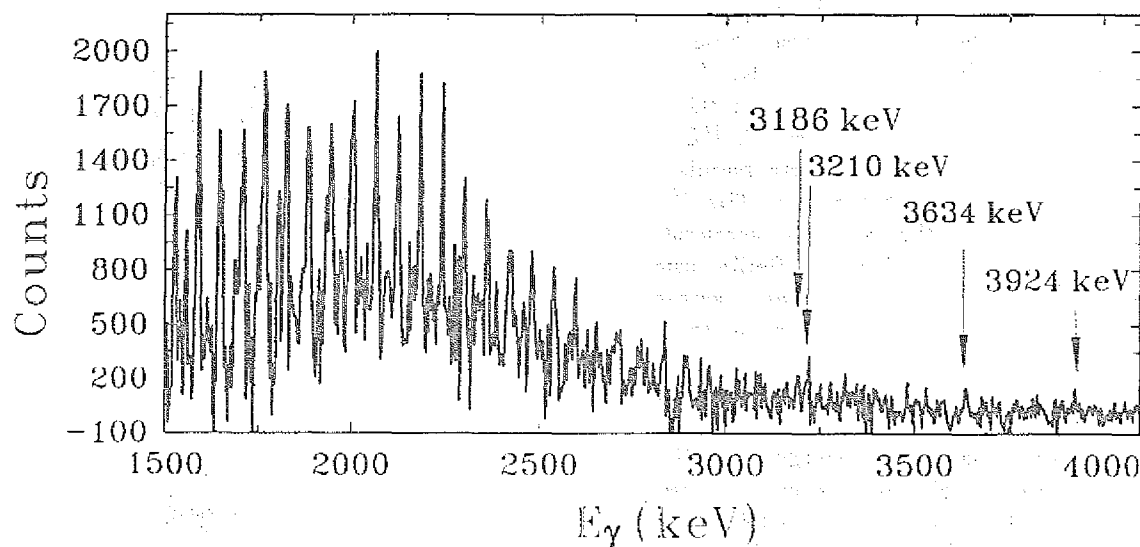


Figure 2: Sum spectrum, displaying the sum of 2 γ -ray energies which are in coincidence with one of the SD band members.

In the energy region between 1.5 MeV to about 3 MeV, all the prominent peaks are identified as being the sum of 2 transitions in the SD band which has a very constant $\theta^{(2)}$. The transitions which connect the SD band with the ND states are expected to be seen above 3 MeV.

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Four experiments were carried out and analysed to search for high-spin states in the Gd isotopes.

In the first two, the beam energy was chosen in order to populate ^{145}Gd most strongly. These experiments were analysed in order to establish superdeformed (SD) states in the nuclei $^{145,146}\text{Gd}$.

The high-spin states in $^{145,146}\text{Gd}$ were populated using a $^{110}\text{Pd}(^{40}\text{Ar}, \text{xn})$, $x=4,5$ reaction. The ^{40}Ar beam was supplied by the VICKSI accelerator system of the Hahn–Meitner–Institut, Berlin, using the CN injector. The γ -radiation was measured with the OSIRIS spectrometer consisting of twelve Compton-suppressed Ge detectors and an inner ball of 48 bismuth germanate (BGO) scintillators, acting as a γ -ray multiplicity filter and calorimeter. (i) To study fast high-energy transitions with optimal energy resolution, we used a target consisting of two self-supporting ^{110}Pd foils with a thickness of 0.5 mg/cm^2 each and a distance of $\approx 0.6 \text{ mm}$ between them. The beam energy was 189 MeV. (ii) We carried out an experiment to measure the effective lifetimes with the Doppler-shift-attenuation (DSA) method. For this experiment a target consisting of a 1.2 mg/cm^2 ^{110}Pd foil backed by a layer of 18 mg/cm^2 ^{197}Au was used. For this experiment the beam energy was 200 MeV. For each event it was required that at least two Ge detectors and the inner ball had fired. We recorded the γ -ray energies and times with respect to the beam burst for the Ge detectors as well as the fold and sum energy for the inner ball. More details on the experimental setup are given in ref. [1].

Analysing the thin target experiment we could observe a ridge structure in the $\gamma\gamma$ -correlation matrix, which is shown in Fig. 1a). According to this structure, SD states have been populated with this reaction. Several indications of discrete SD bands could be extracted from the data of this experiment. One of them led to the discovery of an excited superdeformed band in ^{146}Gd . Information about this band are given in ref. [2]. Our further work concentrated mainly on two other candidates of SD bands. They consist of 8 transitions each, extending over an energy range from about 900 keV to 1300 keV. The energy spacing between the transitions is regular, and has an average value of 58 keV and 56 keV, respectively. Summed coincidence spectra for both bands are depicted in fig. 2. An enhancement of the quality of our current summed spectra should be possible by using improved sorting conditions in the sum-energy-fold plane. The establishment of these conditions

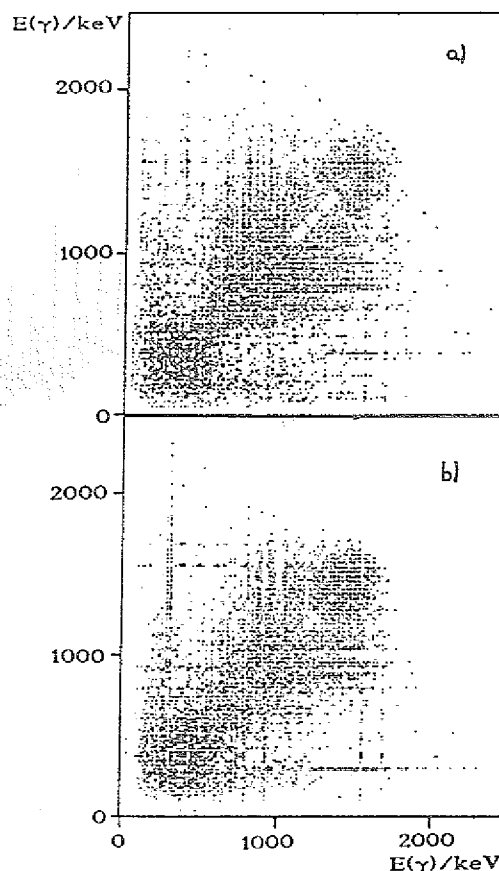


Figure 1: $\gamma\gamma$ -correlation matrices from both thin target experiments, using a $^{110}\text{Pd}(^{40}\text{Ar}, \text{xn})$ reaction at 189 MeV (a) and a $^{108}\text{Pd}(^{40}\text{Ar}, \text{xn})$ reaction at 200 MeV beam energy (b).

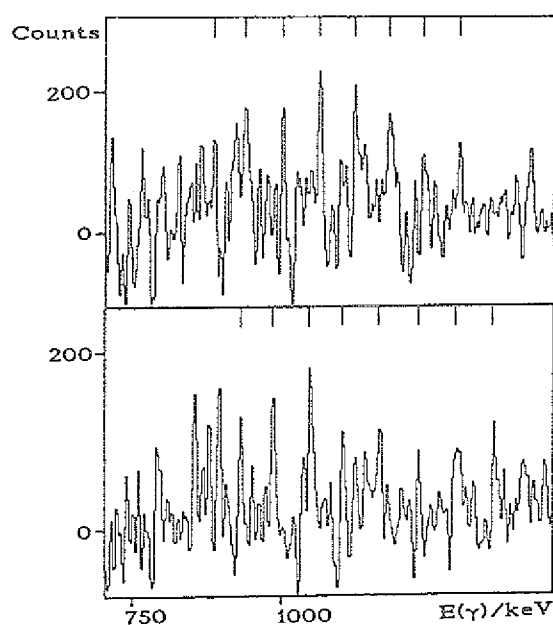


Figure 2: Summed spectra of two sequences of transitions indicating superdeformation in ^{145}Gd .

will be essential for the analysis of the quadrupole moment using the backed target experiment. The extraction of the quadrupole moment is of vital interest to make a final statement about the shape of the nucleus in this configurations.

Two new experiments were made in order to establish superdeformed states mainly in the nucleus ^{144}Gd . The experimental setup and requirements were similar to the ones described above. Using a ^{108}Pd target, ^{144}Gd was populated in a $^{108}\text{Pd}(^{40}\text{Ar}, xn)$ $x=3,4$ reaction. Changing the target from ^{110}Pd to ^{108}Pd allowed a reasonable cross section for the preferred reaction channel, without increasing the fission cross section to an unacceptable value. As it was done successfully before, we used two experiments, with a thin and a backed target respectively. The beam energy was 182 MeV. A look at the $\gamma\gamma$ -correlation matrix in

fig. 1b) shows a ridge structure in the data of the thin target experiment. From this, one can expect to find SD bands also in ^{144}Gd . That is the subject of our current analysis.

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Excited Superdeformed Band in ^{146}Gd

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Excited superdeformed (SD) bands have recently been found in several nuclei in the mass $A=150$ region [1–4]. Here we report on the observation of such a band in ^{146}Gd [5,6]. The yrast SD band in this nucleus was presented in our earlier work [7]. The high-spin states in ^{146}Gd were populated using the $^{110}\text{Pd}(^{40}\text{Ar},4n)$ reaction. The ^{40}Ar beam of 175 MeV was supplied by the VICKSI accelerator of the HMI Berlin. The γ -rays were detected by the OSIRIS spectrometer. Experiments with two self-supporting Pd foils of 0.5 mg/cm² and a backed target (1.2 mg/cm² Pd on 18 mg/cm² ^{197}Au), were carried out. More experimental details are given in ref. [7]. Analyzing the thin-target data we established an excited SD band in ^{146}Gd in addition to the yrast SD band reported before [7]. The band, consisting of 13 approximately equally spaced transitions between 856 keV and 1532 keV, is shown in the sum spectrum of fig. 1.

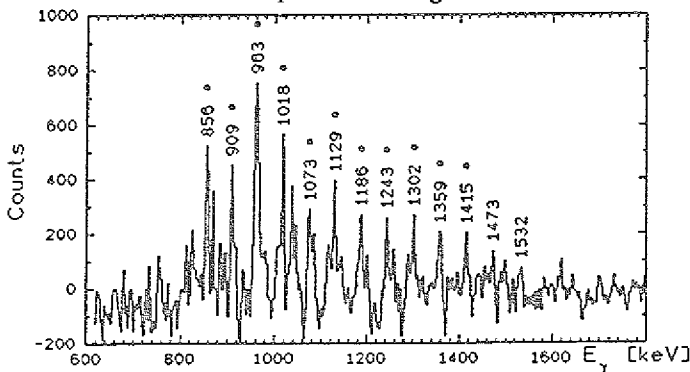


Figure 1: The excited SD band observed in ^{146}Gd .

The lines used as gates are marked with asterisks. From individual coincidence spectra we could prove that each transition is in coincidence with the other members of the band. The intensities of the strongest transitions in the excited SD band are about 0.7 % of the $3^- \rightarrow 0^+$ ground state transition of 1579 keV in ^{146}Gd [8]. The relative intensities of the transitions in the excited SD band show the same behaviour as those in the yrast SD band, namely a slow increase of intensity with decreasing transition energy, followed by a region where the intensities stay roughly constant and finally decrease rapidly corresponding to a depopulation of the band within one or two transitions at a rotational frequency of $\hbar\omega = 0.42$ MeV. In the summed coincidence spectrum we observed γ -rays of ^{146}Gd below the 29^+ state at 12.884 MeV that is populated very strongly [8]. From this we concluded that the spin of the lowest observed member of the excited SD band is $I > 29$.

The same band has been identified independently by Janzen et al. [4] using the $^{122}\text{Sn}(^{29}\text{Si},5n)^{146}\text{Gd}$ reaction. However the assignment of ^{146}Gd was only tentative. In the present work the assignment of the nucleus was made from an analysis of coincidences between fast SD transitions and known low-lying slow transitions in ^{145}Gd and ^{146}Gd . Using the data obtained in the backed-target experiment, summed coincidence spectra were produced by gating on uncontaminated transitions of the SD bands. Care has been taken of the fact that the fast SD transitions are Doppler shifted so that they have different locations for different detector angles. The resulting spectra are shown in fig. 2 for gates placed on transitions of the yrast SD band (a) and the excited SD band (b).

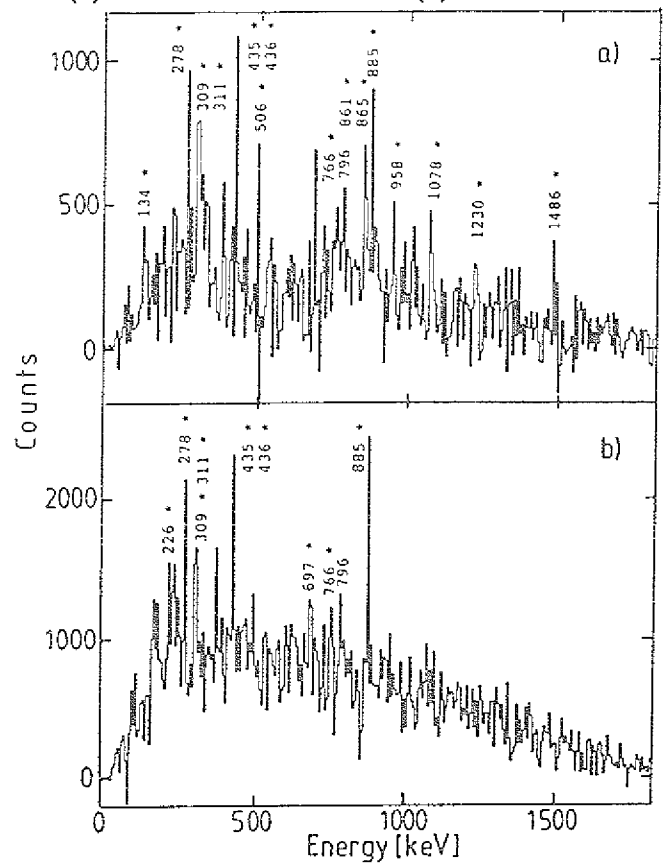


Figure 2: Summed coincidence spectra of slow transitions obtained by gating on fast transitions in the yrast (a) and excited (b) SD bands of ^{146}Gd .

In both spectra known transitions of ^{146}Gd [8] from the ground state up to states above the isomer can be identified. They are marked by asterisks. From the neighbouring nucleus ^{145}Gd only one known transition of 796 keV can be identified. We made a check of the reliability of our method by shifting

our gates in both directions. The summed spectra, produced using these false gates showed an intensity pattern which is completely different from the one seen in fig. 2. The total intensity of the transitions from ^{146}Gd decreased, while the intensities from transitions in ^{145}Gd became remarkably large in these false spectra. The second effect can be understood from the fact that the gates are moved into strong transitions of ^{145}Gd . From this result we are convinced, that our method is reliable in assigning SD bands to nuclei. Therefore we have assigned both SD bands to ^{146}Gd .

The dynamic moment of inertia $\theta^{(2)}$ derived for the second SD band in ^{146}Gd is shown in fig. 3 in comparison to that of the first band.

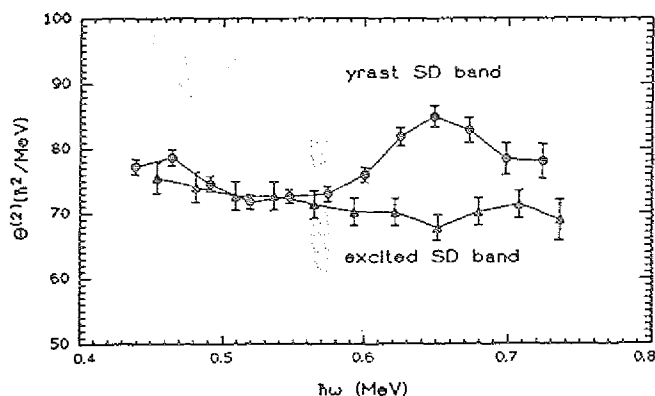


Figure 3: Dynamic moments of inertia versus rotational frequency for SD bands in ^{146}Gd .

For the excited SD band the dynamic moment of inertia changes smoothly with frequency and shows no indication of the band crossing observed in the yrast SD band at $\hbar\omega = 0.65$ MeV.

The transition energies of the second SD band in ^{146}Gd do not lie half way between those of the first one. This may indicate that the SD bands are not signature partners as expected from the configuration assignment $\pi 6^2/7^{1/2}6^1$ given in our previous paper [7].

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The Two- to One-Phonon E3 Transition Strength in ^{148}Gd

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Attempts to identify two-phonon octupole excitations in nuclei have proven to be exceedingly difficult. One reason for this is that in deexcitation of the two-phonon states unspecific low-multipolarity radiation can dominate such that the configuration-specific E3 transition cannot be observed. Some time ago the first case of a two-phonon octupole state in an even-A nucleus was firmly identified [1,2] in $^{148}\text{Gd}_{84}$ through the observation of a $12^+ \rightarrow 9^- \rightarrow 6^+$ stretched double E3 cascade, where the 6^+ basis state is the well established 6^+ yrast level formed by the maximum-spin coupling of the two valence neutrons in the $f_{7/2}$ orbit. From the 9^- level half life measured as 17ns in these experiments significant collectivity was established for the $9^- \rightarrow 6^+$ one- to zero-phonon transition, but the $12^+ \rightarrow 9^-$ two- to one-phonon E3 strength could not be determined since 97% of the 12^+ level decay proceeds through an E1 transition, which suggested a level half life of about 100ps. We have now determined this half life in a plunger measurement with the Nordball detector array consisting of twenty Compton-suppressed Ge-detectors positioned in four conical rings of five detectors each, viewing the target at 37° and 79° to the beam direction in the forward and backward hemisphere.

The measurement is quite difficult since several ^{148}Gd levels above the 12^+ double-octupole state have half lives of comparable magnitude. This requires delicate selection of the production reaction to maximize the side-feeding close to the 12^+ state. The choice of the $^{133}\text{Cs}(^{19}\text{F},4n)$ reaction at 73.5 MeV beam energy was based on such considerations. After passage of the supporting gold foil the beam energy is essentially at the Coulomb barrier. The schematic ^{148}Gd level scheme of Fig. 1 gives the feeding intensity distribution in the plunger measurement.

The target was prepared as a 1.3 mg/cm^2 thick layer by evaporation of undissociated natural CsJ on an $800 \text{ } \mu\text{g/cm}^2$ gold backing, and a 6 mg/cm^2 gold foil was used as stopper. Both were stretched in normal plunger manner. Data were taken at 10 distances ranging from 20 to 3000 μm .

The plunger data were recorded in the two Nordball rings at 37° to the beam. For coincidence gating all 20 Nordball detectors were used. Setting gates on the ^{148}Gd transitions emitted in decay of the 17 ns 9^- isomer gave clean spectra for the ^{148}Gd γ -rays in the region of interest, free from transitions in ^{142}Sm and other reaction channels. These ^{148}Gd data were then evaluated in the manner of singles plunger data.

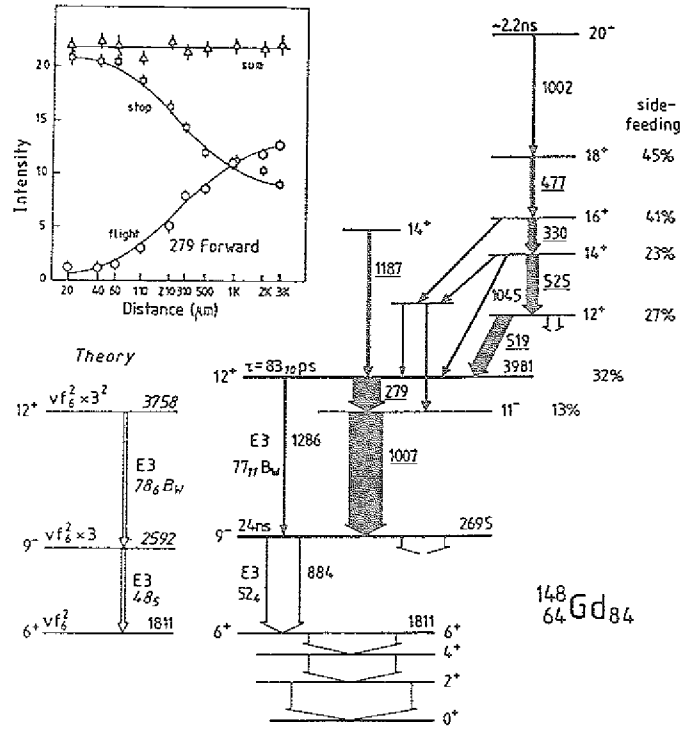


Fig. 1: Schematic partial level scheme of ^{148}Gd as observed in the present experiments. The filled γ -transitions were included in the decay time fit, individual timing data were obtained for the underlined γ 's. Transitions without energies are symbolic and represent the inclusion in the fit of complex known weak branches. Theoretical results are from ref. [2]. Insert: Fit to the data points for the 279 keV 12^+ to 11^- transition measured in the 5 forward detectors at 37° to the beam direction.

In the fit we used the measured timing data and intensities for the seven transitions marked in Fig. 1. Known decay branchings involving weak transitions (shown without energies) were included in the fit by giving their measured intensities, but they were too weak to extract meaningful timing data. The side feeding intensity to each level is represented by two independent components with different effective feeding times which were fitted as free parameters. In total, the fit had 39 parameters compared to 147 data points. The data sets of the two independent measurements from the forward and backward detector rings were analyzed independently and gave the results

$$\tau_{1430}(12^+, 3.98 \text{ MeV}) = 85.0(49) \text{ ps, with } \chi^2 = 3.6, \text{ and}$$

$$\tau_{370}(12^+, 3.98 \text{ MeV}) = 82.1(42) \text{ ps, with } \chi^2 = 2.8,$$

where the statistical errors are quoted. A characteristic property of the data are slow side-feeding components resulting from the fit for most of the levels. This is in accord with other plunger studies, where also side-feeding times comparable to the typical level half lives under study are found. We therefore investigated the stability of the fit against variations of the side-feeding times which probably represent the largest contribution to the systematic error. From the two above determinations we adopt

$$\tau(12^+, 3.98 \text{ MeV}) = 83(10) \text{ ps}$$

where the increased error should cover systematic uncertainties.

For more precise determination of the 279 and 1286 keV branching ratio a separate measurement was made with a backed target at 76 MeV beam energy. The resulting 2.8(2)% 1286 keV E3 decay branch gives

$$B(E3, 12^+ \rightarrow 9^-, 1286 \text{ keV}) = 1.00(14) \cdot 10^5 e^2 \text{ fm}^6 = 77(11) B_W$$

This is the fastest nuclear E3 transition observed so far and firmly characterises the 3.981 MeV 12^+ state of ^{148}Gd as a two-phonon octupole excitation.

The above backed-target data also gave an improved 884 keV E3 branch of the 9^- one-phonon state, of 39(3)%, where the earlier studies [2,3] had given rather different values. With the adopted 16.6(3) ns 9^- half life the transition strength

$$B(E3, 9^- \rightarrow 6^+, 884 \text{ keV}) = 0.70(3) \cdot 10^5 e^2 \text{ fm}^6 = 52(4) B_W$$

results slightly smaller and more accurate than given in Ref. [2].

The ^{148}Gd 12^+ two-phonon state deviates markedly from harmonic vibration where it should have twice the energy and double the E3 strength of the one-phonon state. Such anharmonicities are expected in nuclei and are the subject of current intense theoretical discussion. A probable origin for the ^{148}Gd case is related to the particular microscopic composition of the 3^- state in that nucleus. Here, the two in-shell antialigned 3^- particle-hole excitations involving the high- j orbital of the next-higher oscillator shell lowered by the spin-orbit force, $\pi h_{11/2}$ $\pi d_{5/2}$ and $\nu i_{13/2}$ $\nu f_{7/2}$, will form major components of the octupole phonon. It is clear that the Pauli principle will prevent the stretched double-excitation of these

components and thus shift the maximum-spin two-phonon state upwards in energy and reduce its E3 strength, which both is observed in experiment.

This effect and various other ones causing anharmonicity have been considered in theoretical studies for closed-shell nuclei. But at present straight-forward theoretical results for the specific case of a two-valence-particle nucleus like ^{148}Gd are not easily accessible. In a previous study [4] we have used experimental data of ^{147}Tb on the $\pi h_{11/2} \times 3^-$ septet to calculate the anharmonicity of the two-phonon 6^+ state of ^{146}Gd , but for ^{148}Gd that result is not applicable because it ignores the strong $(\nu f_{7/2} i_{13/2})$ 3^- component in the two-particle nucleus. We have however calculated [2] the 12^+ two-phonon properties using an indirect method, viz. treating it as a system of two $f_{7/2}$ valence neutrons coupled to the ^{146}Gd 6^+ two-phonon state, where we also took into account the coupling to the i_{13} neutron orbital through the large matrix element $m(i, f \times 3) = -0.8 \text{ MeV}$ known from data in ^{147}Gd . This parameterfree calculation predicts the ^{148}Gd 12^+ state at 3.58 MeV, within 0.2 MeV of its observed energy. The $12^+ \rightarrow 9^-$ two- to one-phonon E3 transition strength is calculated as $78(6) B_W$, in good agreement with the $77(11) B_W$ measured in the present experiment.

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Energy Inversion of the $f_{7/2}$ and $h_{9/2}$ Neutrons in Yrast Configurations of $^{154}\text{Yb}_{84}$

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The $N=84$ nucleus ^{154}Yb was investigated using the $^{106}\text{Cd}(^{54}\text{Fe}, \alpha 2p)$ reaction at 4.7 MeV/u beam energy, where 9% of the fusion-evaporation cross section leads to this exit channel. The $\gamma\gamma$ coincidence data gave the ^{154}Yb high spin level scheme of Fig. 1. A new yrast isomer was found at 4.479 MeV excitation energy and its 18.4(15)ns half life was extracted from the $t_{\gamma\text{RF}}$ data for the 622, 782, and 869 keV transitions (Fig. 2). The half life of the previously known¹⁾ lower-lying 8^+ isomer was redetermined using the $t_{\gamma\gamma}$ data of the three strongest transitions below and above it. The resulting 28(2)ns 8^+ half life is about twice shorter than the value from the earlier measurement, where only one isomeric state was known. No apparent delay was recognized in the $t_{\gamma\text{RF}}$ distribution of the higher-lying γ -rays.

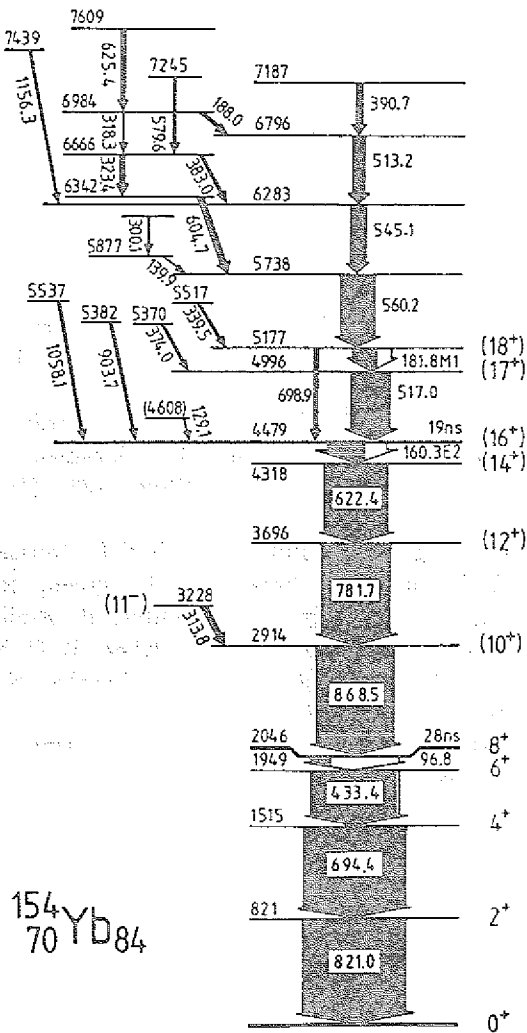


Fig. 1: High spin states in $^{154}\text{Yb}_{84}$ with transition intensities as observed in the $^{106}\text{Cd}(^{54}\text{Fe}, \alpha 2p)$ reaction at 4.7 MeV/ μ beam energy.

In our experiment we had also planned to identify new exit channels through measurement of neutron-gamma coincidences. The Osiris Ring was therefore placed perpendicular to the beam axis, which prevented DCO ratio measurements for spin determinations. But up to 8^+ the ^{154}Yb level spins are firm from β -decay results²⁾, and the tentative 10^+ to 16^+ assignments are based on energy systematics of the $N=84$ isotones. Our data give E2 for the 161 keV $16^+ \rightarrow 14^+$ transition from intensity balance.

An unexpected result was found above the 16^+ state where the main cascade suggests 17^+ and 18^+ for the two next higher states. This is inferred from the measured M1 nature of the 182 keV $18^+ \rightarrow 17^+$ transition, and from the observation of an $18^+ \rightarrow 16^+$ cross over γ -ray. At present no spin assignments can be made for other states in ^{154}Yb .

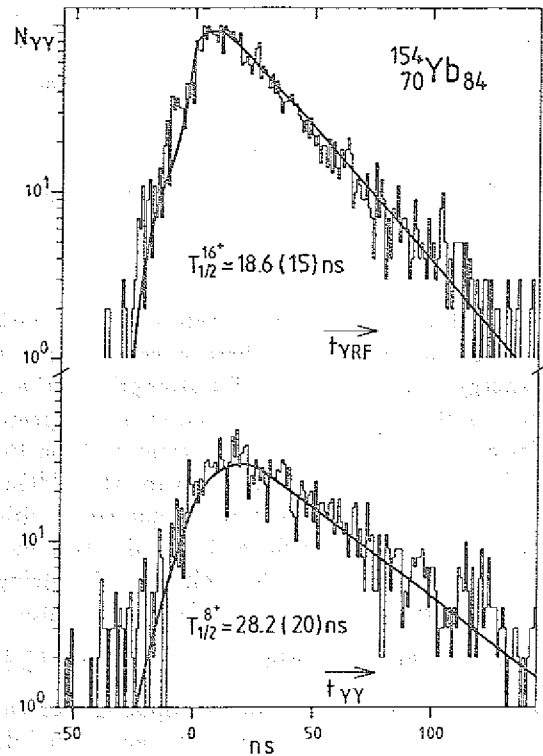


Fig. 2: Isomer half lives in ^{154}Yb . The upper curve displays $t_{\gamma\text{RF}}$ -data for the 622, 782, and 869 keV lines which were accepted only when in coincidence with other intense ^{154}Yb yrast transitions. The data points in the lower curve display $t_{\gamma\gamma}$ for the 9 possible pairs of intense γ -transitions with energy above 350 keV that have the 8^+ isomer between them. Both curves analyzed by deconvolution with the pertinent prompt response function.

The 16^+ and 14^+ energy separation decreases significantly in ^{154}Yb compared to the lower Z neighbours, and also its $B(E2)$ value of $3.6(2)B_W$ is much lower than in ^{152}Er ($\sim 10.5B_W$) and ^{150}Dy ($9.6(11)B_W$)^{3 4)}. These results, together with the above observed 17^+ and 18^+ states strongly suggest $(\pi h_{11/2}^6 \nu f_{7/2} h_{9/2})$ character for these three highest assigned levels. Manybody shell model calculations with empirical nucleon-nucleon interactions strongly support this conclusion (Fig. 3). They clearly reveal a significant

energy drop for the $(\pi h_{11/2}^6 \nu f_{7/2} h_{9/2})16^+$ (and 17^+) states while the 18^+ level stays at constant energy when proton pairs are added to ^{150}Dy . A similar effect is the well established⁵⁾ inversion of the $(\pi h_{11/2}^6 \nu h_{9/2})10^+$ and 9^+ states in the odd-odd $N=83$ isotones. It is related to the strongly attractive $(\pi h_{11/2}^6 \nu h_{9/2})1^+$ interaction that is blocked by the Pauli principle in the $(10^+$ and 18^+ coupling. In ^{154}Yb the $(\pi h_{11/2}^6 \nu f_{7/2}^2)16^+$ state is predicted 355 keV above the 14^+ state, clearly higher than the observed 16_1^+ yrast level, while the predicted $(\pi h_{11/2}^6 \nu f_{7/2}^2)14^+$ to $(\pi h_{11/2}^6 \nu f_{7/2} h_{9/2})16^+$ splitting is in very close agreement with experiment. Our results have thus established the energy inversion of the $f_{7/2}$ and $h_{9/2}$ neutron orbits in ^{154}Yb at $I^\pi=16^+$. In the higher- Z $N=84$ isotones, the inversion occurs already at $25/2^-$ or 8^+ , and gives rise to the high-spin α -decay isomers⁶⁾.

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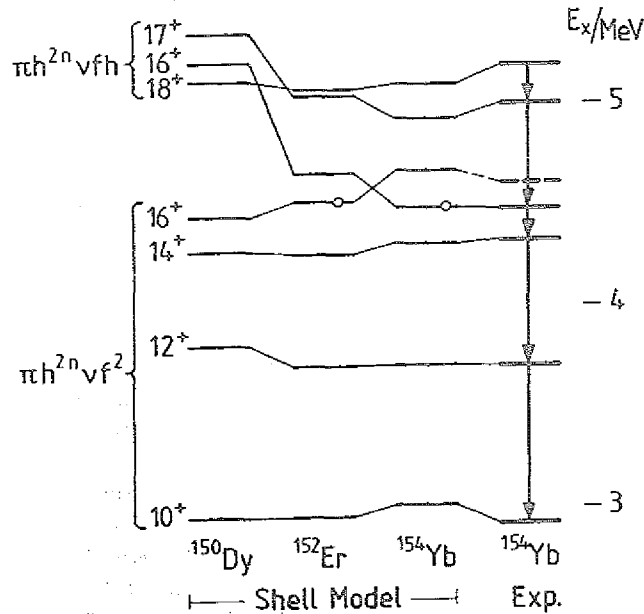


Fig. 3: The observed $I^\pi = 10^+$ to 18^+ ^{154}Yb yrast levels compared to the results of multi-particle shell model calculations with empirical twobody interactions. For ^{154}Yb and ^{152}Er the calculated energy levels are normalised (o) to the experimental 16_1^+ energies, the ^{150}Dy excitation energies are based on ground state masses.

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In previous works [1–3], following the γ – γ energy correlation data analysis, experimental evidence suggesting the existence of superdeformed (SD) structures in ^{180}Os was presented. A sequence of 9 γ –ray transitions in coincidence, with an average constant separation between them of ~ 36 keV was identified in the energy range of 900–1200 keV [2]. The deduced value for the dynamic moment of inertia was $\theta^{(2)} = 111 \hbar^2/\text{MeV}$ which corresponds to an axis ratio of $c/a \sim 1.7$ and a quadrupole deformation of $\beta_2 \sim 0.6$. These values would indicate a SD shape of the nucleus. The subsequent work reported here was performed to confirm the sequence of γ –ray transitions mentioned above and to measure the quadrupole moment associated with this sequence.

Two further experiments were carried out at the VICKSI accelerator of the Hahn–Meitner–Institut in Berlin. High-spin states in ^{180}Os were populated through the $^{150}\text{Nd}(^{36}\text{S},6n)$ reaction at 177 MeV. The emitted γ –radiation was measured with the OSIRIS spectrometer. To study fast transitions with minimum Doppler broadening, a ^{150}Nd foil of ~ 0.5 mg/cm² thickness was used in the first experiment whereas to obtain information about lifetimes using the DSA method, a ~ 1 mg/cm² ^{150}Nd target with a 7 mg/cm² ^{197}Au backing was used in the second one.

The results obtained from the γ – γ correlation matrix analysis of the thin target experiment ($\theta^{(2)}$, c/a and β_2) agree well with the previous one [1–3] using the $^{150}\text{Nd}(^{34}\text{S},4n)$ reaction. The search for discrete lines producing the ridge structures in the γ – γ correlation matrices was carried out using the method described in [4], because gating on individual lines proved not to be useful due to the low statistic of the data set. As a result of the search a sequence of 12 lines shown in fig. 1 was obtained. One should remark that the first 9 transitions have approximately the same γ –ray energies as for the sequence found in [2].

Coincidence events obtained from the thick target experiment were sorted in 16 γ – γ correlation matrices, each one characterized according to the fractional Doppler-shift (F) applied to it. They were sorted using high multiplicity ($M \geq 12$) and sum energy conditions, in order to enhance ^{180}Os in the final matrices. The results obtained from the analysis of these γ – γ energy correlation matrices are summarized below. A ridge structure is observed with an energy width of $2\Delta E_\gamma = 96 \pm 3$ keV for γ –rays with energies of 600–860 keV in the matrix with $F = 0$. Assuming a symmetric ellipsoid of prolate shape for the nucleus a value of $\theta^{(2)} = 84 \hbar^2/\text{MeV}$ is obtained which describes a

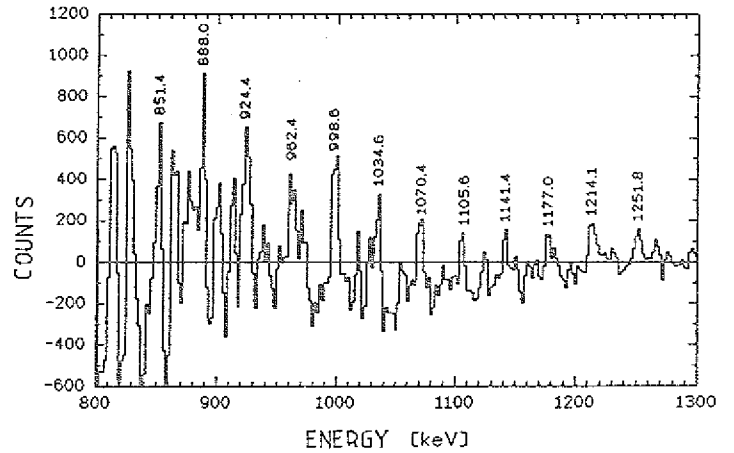


Figure 1: Correlation spectrum obtained using the method described in [4,5] for a constant energy separation of $\Delta E_\gamma = 36$ keV. The energy of the γ –ray transitions are given in keV.

slightly deformed shape of the nucleus and that corresponds mainly to the normal deformed states at moderately high-spin. The part of the ridge structure above 740 keV is found to have a fast and slow time component. A sharp ridge structure is seen for $F = 0$ as well as for $F = 0.7$ of the total Doppler-shift. From Monte–Carlo simulations follows that a Doppler-shift correction of $F = 0.7$ corresponds to a lifetime of 300 fs. Above 960 keV a ridge structure with a width of $2\Delta E_\gamma = 72 \pm 3$ keV occurs, from which a value of $\theta^{(2)} = 111 \hbar^2/\text{MeV}$ was obtained. Applying the DSA method, lifetime information was obtained and a further analysis of the Doppler–shift corrected matrices for the same energy range allowed us to obtain an estimate of the quadrupole moment of $Q_0 = 10 \pm 4$ eb, indicating a highly deformed structure of the nucleus which corresponds to a quadrupole deformation of $\beta_2 = 0.35$ and $c/a \sim 1.37$ (45 % and 23 %, respectively, smaller than the values determined from the thin–target experiment). In fig. 2 a plot of the experimental fraction of the total Doppler-shift vs. energy is compared with theoretical calculations for different values of the quadrupole moment.

Above 1 MeV a further ridge structure formed like a trumpet is observed in the 70 %–matrix (fig. 3) and in fig. 4 the deduced dynamic moment of inertia for different cuts perpendicular to the $E\gamma_1 = E\gamma_2$ diagonal are plotted vs. energy. This ridge structure is characterized by a gradual decrease of the moment of inertia [5] and might correspond to the predicted shape change from an almost prolate shape ($\gamma = -10^\circ$) to a strongly–triaxial shape ($\gamma = -90^\circ$) [6].

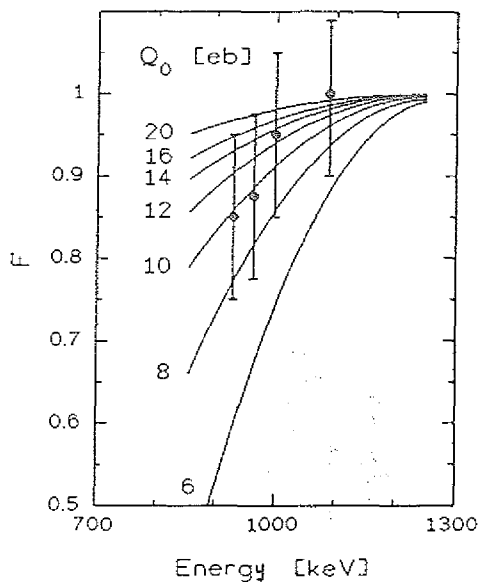


Figure 2: Experimental fractions of the Doppler-shift for some transitions of fig. 1. For comparison calculated values are given as solid lines for different values of the quadrupole moment.

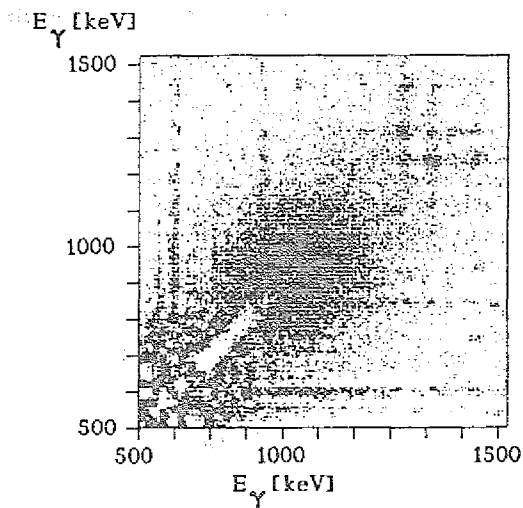


Figure 3: 70 % Doppler-shift corrected E_γ - E_γ correlation matrix.

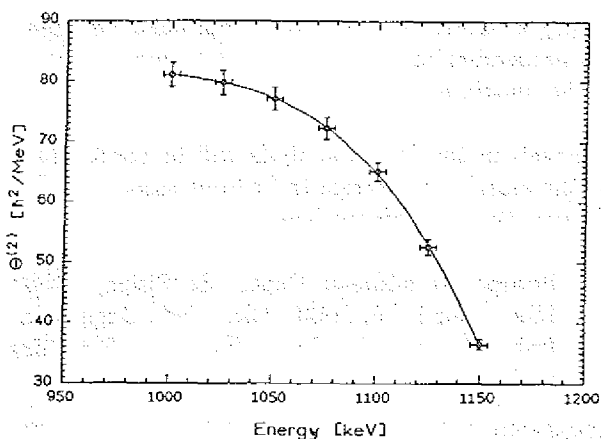


Figure 4: Variation of the dynamic moment of inertia ($\theta^{(2)}$) with energy for $E_\gamma > 1$ MeV, for the E_γ - E_γ correlation matrix (shown in fig. 3).

In fig. 5 total spin projection I_x values (obtained from the experiment as described in [5]) are plotted vs. rotational frequency. For comparison the calculated values for the expected strongly-prolate ($\gamma \sim 0^\circ$, $\beta \sim 0.4$), strongly-triaxial ($\gamma \sim -90^\circ$, $\beta \sim 0.15$) and triaxial shapes ($\gamma \sim -10^\circ$, $\beta \sim 0.2$) are also shown.

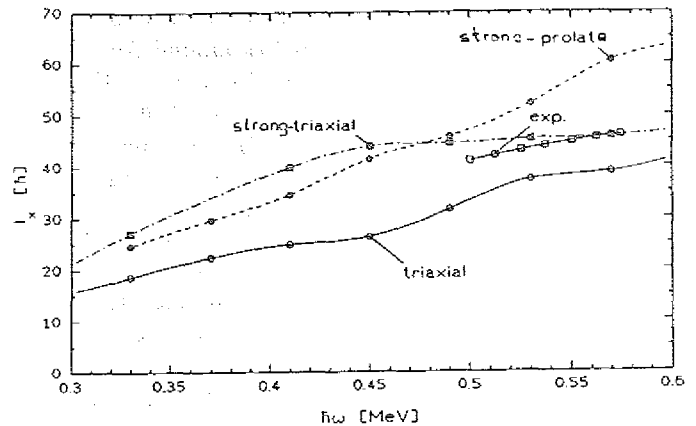


Figure 5: Total spin projection I_x as a function of the rotational frequency $\hbar\omega$. Theoretical values [6] for the assumed strongly-triaxial (---), strongly-prolate (---) and triaxial (—) shapes are also shown for comparison.

The agreement between the experimental and the strongly-triaxial values gives additional support to our interpretation that the states producing the trumpet-shaped ridges are mainly associated with a continuous change of shape, driving the nucleus to states of low collectivity with increasing energy.

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Several theoretical calculations have indicated a number of regions in which superdeformed (SD) structures $\beta_2 \geq 0.4$ $c/a \approx 2$ could become yrast at high-spins. One of these regions lies around $A=190$ ($Z=80$, $N=110$) and it was highlighted in [1] that the crossing occurs at spins close to $40 \hbar$. This prediction was soon followed by the experimental observation of several SD bands in Hg, Tl and Pb isotopes [2]. Further theoretical calculations [1,3] predict that SD structures should also exist in Au [$Z=79$] and Pt [$Z=78$] nuclei and that the SD minima should be deep enough to allow population through heavy-ion reactions. Since it has been shown that the SD structures in the Hg and Tl nuclei can be populated with rather light projectiles [4], it appears possible to exploit this fact to populate SD bands in the Au nuclei around $A=190$. High-spin states in $^{189,190,191}\text{Au}$ were populated through the $^{176}\text{Yb}(^{19}\text{F},xn)^{195-x}\text{Au}$ reaction at a beam energy of 107 MeV. The target consisted of 2 self-supporting ^{176}Yb foils (97.5 % enrichment) of $\sim 500 \mu\text{g}/\text{cm}^2$ thickness. The γ -rays were emitted from the recoiling Au in flight so that they were fully Doppler-shifted. The beam was provided by the VICKSI accelerator of the HMI, Berlin. The γ -radiation was detected with the OSIRIS multi-detector array, consisting of 12 Compton-suppressed Ge detectors and 48 BGO elements used as sum-energy and multiplicity filter. A total of $\sim 650 \cdot 10^6$ Ge-Ge-BGO events (after unfolding three- and four-fold events) were recorded and subsequently sorted into several two-dimensional $E\gamma$ - $E\gamma$ correlation matrices. During off-line sort of the data the γ -ray energies were corrected for Doppler shift.

Two-dimensional $E\gamma$ - $E\gamma$ correlation matrices for ^{190}Au were obtained by setting a two-dimensional gate in the sum-energy vs. fold spectrum. To find the proper shape of the gate, the entry distribution for the different exit channels of the $4n$, $5n$ and $6n$ reactions was determined by setting gates on strong discrete lines in the nuclei $^{189,190,191}\text{Au}$. One of the $E\gamma_1$ vs $E\gamma_2$ correlation matrix for ^{190}Au is shown in fig. 1.

In this mass region SD bands are expected in the 200–800 keV energy range. No ridge structures could be seen below ~ 600 keV, probably due to disturbing strong discrete lines.

However a weak ridge structure appears between 600 and 900 keV. In fig. 2 a cut perpendicular to the main diagonal $E\gamma_1=E\gamma_2$ for the energy range 600–800 keV shows that the average separation between the ridges is $2\Delta E\gamma = 64 \pm 4$ keV. The data analysis is still in progress (the experiment was made at the beginning of Sept. 1991) and our main

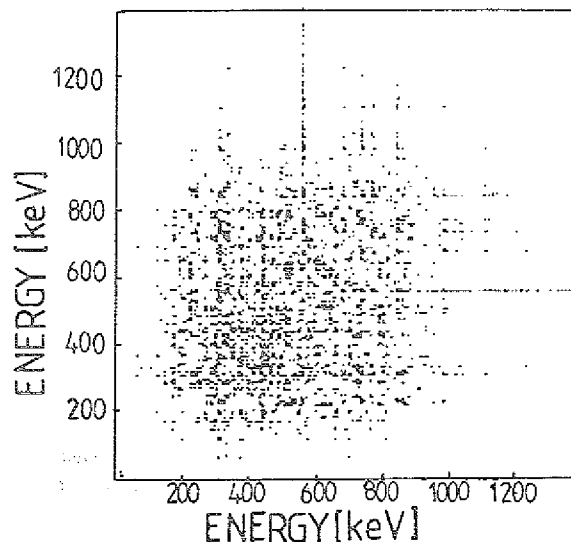


Figure 1: $E\gamma$ - $E\gamma$ correlation matrix for ^{190}Au . The coincidence events were sorted with a prompt time window (≤ 16 ns) and a 2-dim. gate corresponding to a large multiplicity and sum energy.

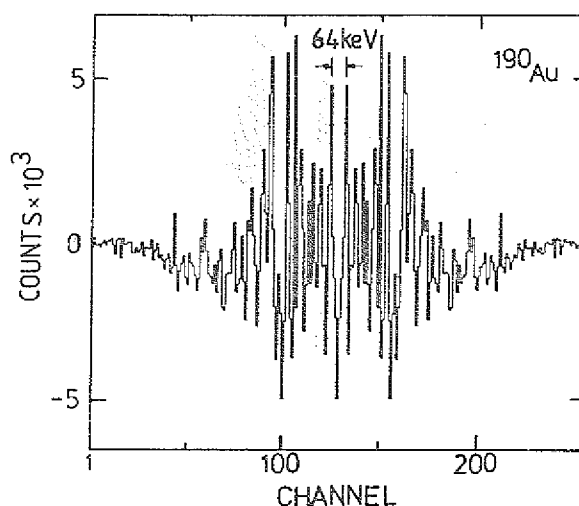


Figure 2: Correlation spectrum obtained through a cut perpendicular to the $E\gamma_1 = E\gamma_2$ main diagonal of the matrix of fig. 1.

emphasis in the future analysis will be concentrated on the search for discrete transitions associated with SD structures in this nucleus.

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High Resolution Measurement of K X-rays of Lanthanum Following the Collisions with 403 MeV N^{6+} Ions

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In an energetic collision between charged particles and target atoms the atom is left with large probability in a multiply ionized state. Due to the change in screening the corresponding X-ray transitions have energies which are somewhat higher than from a singly ionized atom and which are characteristic for the vacancy distributions. The corresponding X-ray spectra therefore exhibit a satellite structure on the high energy side of the diagram lines, which yields information about the production and relaxation mechanisms.

K-shell vacancy production by ion impact occurs predominantly in close collisions at distances even smaller than the K-shell radius. In contrast, L-shell vacancies are produced in a large impact parameter range up to several times the L-shell radius. The L-shell ionization probability changes only slowly in the impact parameter range where K-shell ionization occurs. Consequently, K X-ray satellite spectra can be used as a tool for the study of L ionization probabilities for impact parameters less than the K-shell radius, in particular in the case of fast collisions where the coincidence method between X-rays and scattered ions is not feasible.

As a last experiment of a systematic study we have investigated the K X-ray satellite spectra produced during the bombardment of lanthanum with 403 MeV N^{6+} [1, 2, 3]. The experimental spectra have been analyzed and compared with the theoretical predictions. We put special emphasis on the investigation of 3 subjects: 1) Comparison of different theoret-

tical models; 2) Angular momentum coupling; 3) Change of fluorescence yields.

The ionization probabilities have been calculated in the framework of the semiclassical separated atom approximation (SCA) using two models, based on two different potentials.

Model I provides relativistic calculations using a straight line path as trajectory and a variationally determined potential, which gives wavefunctions very similar to Hartree-Fock wave-functions. Model II uses screened relativistic electron wave-functions and the projectile passes on a Rutherford trajectory. The target recoil is taken into account. Model II predicts an L-shell ionization probability of about a factor 2 less than model I.

In Fig. 1 the measured $K\alpha_1$ spectrum is compared with the predictions of the two models. The wave-functions of the Hartree-Fock-like potential of model I, which takes into account the presence of the passive target electrons, predict L ionization probabilities, which reproduce the experimental spectra almost quantitatively. This is not the case for the hydrogenic wave-functions of Model II. An estimate shows that the recoil effect has negligible influence on the ionization probabilities. As the maximum contribution to the total Coulomb ionization cross section is given in first approximation for an impact parameter proportional to the projectile's velocity, it can be shown that the projectile's Rutherford trajectory is well described by a straight line path. Conse-

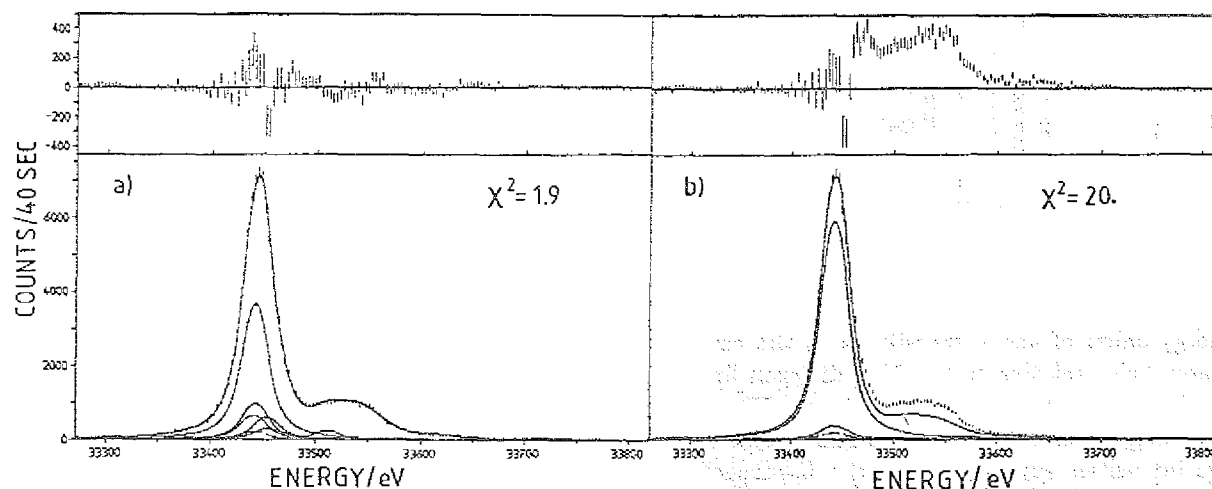


Fig. 1: Comparison of models I and II for the ionization probabilities in the framework of the SCA, including the angular momentum coupling for the L spectator holes. Capture and L subshell coupling have been omitted as they almost cancel a) model I, b) model II (see text).

quently, the origin of the difference between the two models has to be explained from the choice of the potentials.

Capture of L-shell target electrons into atomic orbits of the projectile and coupled-channels effects have been estimated. Their contributions have been calculated to be less than 5% of the ionization probabilities predicted by model I. Both contributions have similar values but differ in sign so that they almost cancel.

In a multiply ionized atom there exist generally at least two open atomic shells. As an example for a $K\alpha_1$ transition, in the initial state there is the 1s hole plus the shell(s) with the spectator hole(s) and in the final state the $2p_{3/2}$ hole plus the same spectators. Due to the coupling of the angular momenta the initial and final state will split, leading to a splitting of the corresponding X-ray transitions. In this example, the three satellite lines due to the presence of a L_I , L_2 or a L_3 hole will split into 9 transitions. The positions of these 9 lines relative to the diagram line are shown in Fig. 2 in units of the total reflection width of the $K\alpha_1$, which was measured with a 25 MeV proton beam to be 30 eV (FWHM). The splitting due to the angular momentum coupling redistributes the intensity of the transitions in a pronounced way. It is obvious that ignoring this splitting the experimental spectra cannot be reproduced.

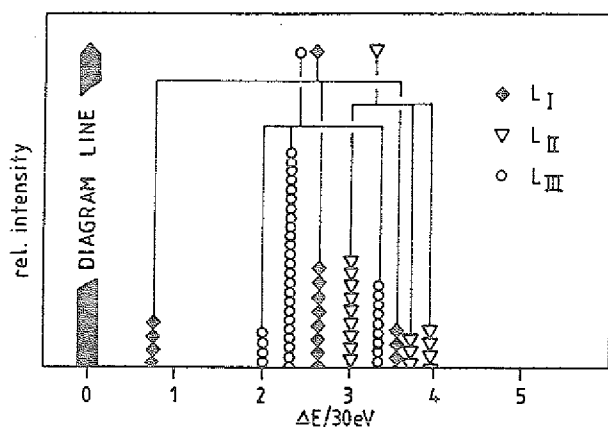


Fig. 2: Energy shifts of the 9 transitions in the case of L spectator holes relative to the $K\alpha_1$ diagram line in units of the total $K\alpha_1$ reflection width (FWHM = 30 eV). The heights of the lower bars correspond to the relative intensities supposing that the L subshell hole distribution is $L_I/L_{II}/L_{III} = 1/1/2$. The positions of the L spectator hole satellites, without taking into account the splitting due to angular momentum coupling are shown in the upper part of the figure.

In a multiply ionized atom the number and the distribution of the decay channels is changed as compared to the usual singly ionized situation. We estimated the change of the K fluorescence yields in the presence of L spectator holes. It could be shown that the total K fluorescence yield does not change by more than 3%, but the partial fluorescence yields change considerably (up to 50%). For example, in singly ionized La the $K\alpha_1$ partial yield is 53%, but in the presence of a L_{II} spectator hole this $K\alpha_1$ partial yield is increased to 64%. This leads to a drastic redistribution of the intensity of the corresponding satellite structure.

In conclusion, it could be demonstrated that each of the 3 subjects discussed play an important role for the correct understanding of the experimental satellite spectra.

We are very much indebted to P.A. Amundsen for providing us with the calculations of the ionization probabilities and to D.H. Jakubassa-Amundsen for performing the calculations of the L subshell coupling and the electron capture process. We are very grateful to D. Trautmann for supplying us with his code for the calculation of the ionization cross sections.

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Development of a Composite Ge Detector for EUROBALL

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Future γ -ray arrays are characterized by a significantly increased total detection efficiency. Present arrays, like OSIRIS, have a total photopeak efficiency of $P_{ph} < 1\%$ for 1.3 MeV γ -rays. For EUROBALL III the total photopeak efficiency should lie well above 10%. With such an array 4-fold coincidences can be detected with a much larger efficiency than 2-fold coincidences with present-day arrays. This implies that correlated photo events, as resulting from rotational bands, can be identified much better on a background of uncorrelated events. This results from the fact that the photo events are confined to a small volume element whereas the background events distribute evenly over the full multi-dimensional volume. This effect can be described by a background reduction factor R [1] which is for EUROBALL III about 3 orders of magnitude larger than, e.g., for OSIRIS.

In order to increase the total photopeak efficiency of the array a higher efficiency of the individual Ge detector is needed. It is unlikely that the volume of Ge detectors can be enlarged by another factor of 2 or more in the near future. Therefore very big detectors have to be composed from several crystals in a compact way. By adding all signals of the individual detectors the total-absorption efficiency and the peak-to-total ratio of a very large detector is achieved. The granularity of such a detector provides at the same time the small opening angle for the detection of γ -rays needed to reduce the resolution deterioration due to the Doppler effect. Such a composite detector is the cluster detector. The basic design of the cluster detector is shown schematically in Fig. 1. It consists of seven large hexagonal tapered Ge detectors which are individually encapsulated and closely packed in a common cryostat. The development of encapsulated Ge detectors was motivated by the demand to achieve a reliable operation and maintenance of multi-detector arrays.

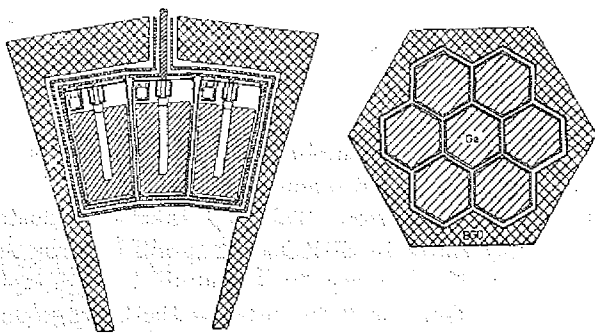


Figure 1: Basic design of the cluster detector.

It is well known that the intrinsic surface of Ge detectors is very sensitive to impurities of any kind. No method of surface passivation is known which gives a 100% probability of low leakage current when the detector is cooled to the temperature of liquid nitrogen. The main idea of encapsulated detectors is to decouple the vacuum of the individual Ge detector from the common cryogenic vacuum of an array. There are several advantages of encapsulated detectors: If there is a failure of one detector in an array it can be exchanged easily by the user; the detectors can be rearranged for different applications using different cryostats and the detectors can be easily annealed after neutron damage. The encapsulated Ge detectors can be removed from the cryostat and can be annealed in a standard oven in a few hours without risk.

The feasibility of the cluster detector has been demonstrated by a successful prototype development of a vacuum-tight encapsulated Ge detector by a collaboration of the KFA Jülich, the University of Cologne and Inter technique Strasbourg and of a hexagonal tapered Ge detector by a collaboration of Inter technique Strasbourg, the University of Cologne and the HMI Berlin. A prototype of the cluster detector is under construction and is expected to become available in summer of 1992. The project is funded by the German Minister of Research and Technology under contract No. 06 OK 143.

In order to optimize the cluster detector with respect to photopeak efficiency ϵ_{ph} and P/T ratio, Monte-Carlo simulation calculations for various geometries have been carried out using the code GEANT3. The P/T ratio describes which fraction of the total number of events (T) in a spectrum can be found in the photopeak (P). The hexagonal tapered detector is produced from a cylindrical Ge crystal which means that about one third of the Ge crystal has to be cut away. It is therefore worthwhile to check the effect of shaping if detectors are clustered since it is not obvious whether it is better to take cylindrical detectors of maximum volume and tolerate that the Compton-scattered γ -rays can escape through the resulting gaps or to accept a smaller detector volume but obtain an optimal packing. Ge detectors of hexagonal tapered, hexagonal untapered, conical, cylindrical and so-called hexaconical shape have been considered in the calculations. A hexaconical encapsulated detector is shown in Fig. 2. It has a hexagonal front face and a circular shape at the rear side of the Ge detector. This detector has a volume of 93% of the cylindrical detector but still allows a fairly compact packing.

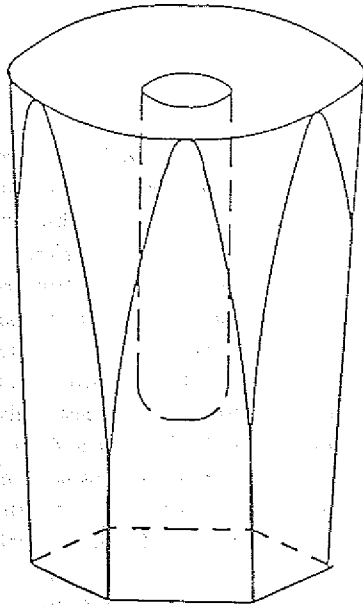


Figure 2: The hexaconical Ge detector.

The results of the calculations for various cluster geometries are given in Table 1.

Table 1: Results of Monte-Carlo simulation calculations for various cluster geometries for $E_\gamma = 1.33$ MeV and a fixed solid angle of $\Omega = 1.03$ %.

Shape of Ge det.	Volume cm ³	Source dist. cm	ϵ_{ph}	(P/T) _{uns.}	(P/T) _{suppr.}
Hexagonal tapered	1559	33.3	0.30	0.39	0.84
Hexaconical	2039	43.6	0.31	0.40	0.84
Hexagonal untapered	1852	45.5	0.29	0.40	0.83
Conical	1870	43.6	0.29	0.39	0.85
Cylindrical	2249	51.4	0.29	0.40	0.84

Each cluster consists of seven encapsulated Ge detectors which fill a fixed solid angle of $\Omega = 1.03$ % and which are surrounded by the same thickness of BGO. It can be seen from Table 1 that Cluster modules consisting of hexagonal or hexaconical Ge detectors have the largest efficiencies ϵ_{ph} while the P/T ratios are similar for all detector shapes. As for the hexaconical shape the largest amount of the Ge crystal is conserved, this geometry was chosen for the prototype of the first cluster detector. In addition, finite elements calculations for the stability of the detector capsule have shown that a hexaconical shape is very favorable and allows to reduce the wall thickness of the aluminium can to 0.7 mm. First hexaconical capsules have been produced and successfully tested under vacuum up to a temperature of 150°C. Seven hexaconical encapsulated Ge detectors are ordered from Intertechnique. Each detector will have a diameter of 70 mm at the circular rear end and a length of 78 mm giving an estimated relative efficiency of 65 % at 1.3 MeV.

The EUROBALL will be built in three phases. EUROBALL I consists of 45 individual Compton-suppressed Ge detectors. EUROBALL II consists of 30 such detectors and 24 clover detectors. A clover detector contains four small Ge detectors in a common cryostat surrounded by a BGO anti-compton shield [1]. EUROBALL III consists of 15 individual Compton-suppressed Ge detectors, 36 clover detectors and 15 cluster detectors. A side view of EUROBALL III is shown in fig. 3.

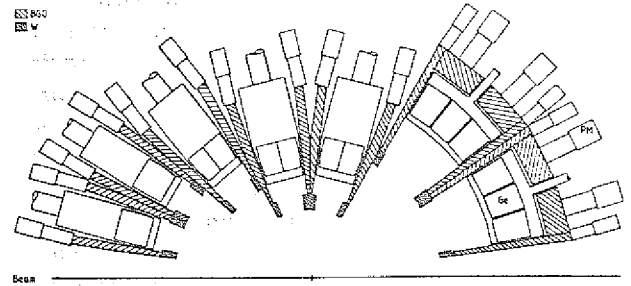


Figure 3: Side view of the 4π array EUROBALL III consisting of 15 cluster detectors, 15 individual Compton-suppressed Ge detectors and 36 clover detectors.

The cluster detectors form a sub-array of $\approx 1\pi$ solid angle. The cluster segment is depicted in fig. 4.

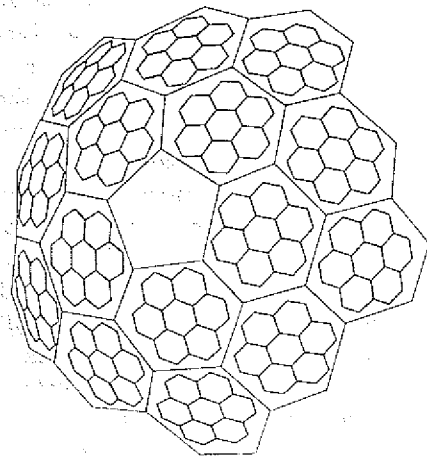


Figure 4: Geometry of the cluster sub-array consisting of 15 hexagonal faces surrounding a pentagonal face.

Table 2 gives a compilation of the features of different existing and planned arrays.

It can be seen from Table 2 that the cluster sub-array added to EUROGAM II will increase the sensitivity of the array for the study of high-spin states by a factor of ≈ 20 and that the background reduction factor R is ≈ 700 times larger for EUROBALL III as compared to the existing arrays.

Table 2: Features of arrays for $E_\gamma = 1.33$ MeV and $M_\gamma = 30$. The number of detector systems N and their energy resolution ΔE_γ and effective P/T ratio are given. Furthermore, the total photopeak efficiency P_{ph} and the background reduction factor R of the array are included.

Array	N	ΔE_γ	$(P/T)_{eff}$	P_{ph}	R
OSIRIS	12	7.3	0.43	0.007	1.3
NORDBALL	20	7.3	0.43	0.011	1.5
GASP	40	7.1	0.47	0.035	4.3
EUROBALL I	45	7.1	0.47	0.048	7.2
EUROBALL II	54 ^a	6.4	0.52	0.086	62
EUROBALL III	78 ^b	5.2	0.55	0.125	$1.1 \cdot 10^3$

a) EUROBALL II consists of 30 individual Compton-suppressed Ge detectors and 24 clover detectors.

b) EUROBALL III consists of 15 cluster detectors, 15 individual Compton-suppressed Ge detectors and 36 clover detectors.

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Within the polyhedra containing 12 regular pentagons and 60 hexagons, with each pentagon representing a five-fold symmetry axis [1], we searched for the optimal solution satisfying the following 2 requirements: 1) all the (irregular) hexagons to be identical, 2) each hexagon to be as most regular as possible, i.e., to have an inscribed circle tangent at least to 5 out of its 6 sides.

The centers of the 12 pentagons determine on the sphere an array of 20 equilateral spherical triangles ("fundamental triangles"), with $\alpha = 72^\circ$ and sides of $m \sim 64.43^\circ$ (fig. 1). The middle point B of each triangle belongs to three neighbouring hexagons. Let us fix one hexagon with vertices A_j ($j=1, \dots, 6$), where A_1A_2 is the side of the neighbouring pentagon and therefore $A_5 = B$. Let a be the apothem of the pentagons and P the common point of the plane containing the hexagon and the sphere. Fixing an orthogonal planar coordinate system with origin in P and being the y -axis orthogonal to A_1A_2 (fig. 2) let γ be the spherical distance between M (the center of the pentagon) and P (fig. 3).

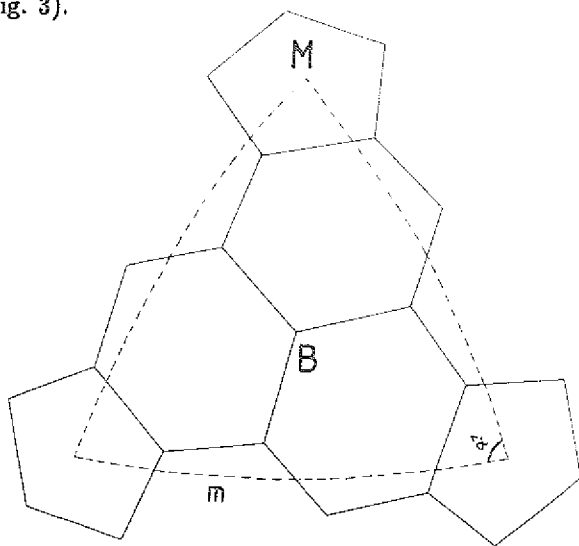


Figure 1: One of the 20 fundamental triangles on the sphere formed by the central points of the pentagons.

A circle in the xy -plane with center in P can fulfill the second requirement if it is simultaneously tangent to A_2A_3 and A_1A_6 (otherwise it would be tangent to at most 4 sides). Therefore its radius can be written as

$$r = a(1 + \cos \gamma) / (\tan^2 54^\circ + \cos^2 \gamma)^{1/2}.$$

The additional requirement for the circle to be inscribed in the hexagon ($r \leq a$) leads to the unreasonable condition $\gamma \geq m$. If we shift the center of the circle out of P , the normal to the plane containing the hexagon in this new point, would not go through the center of the sphere. Therefore we have choose, following [2], a cut-off

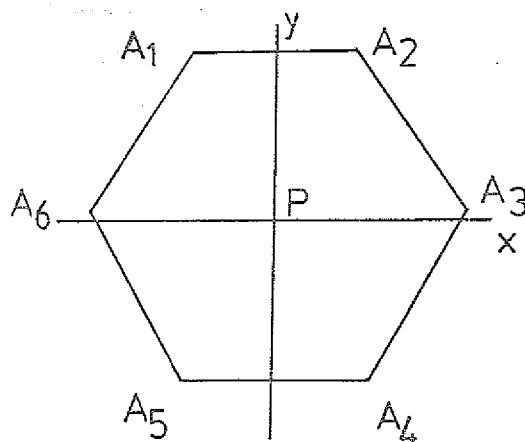


Figure 2: Projection of one hexagon into a plane tangent to the sphere in the point P .

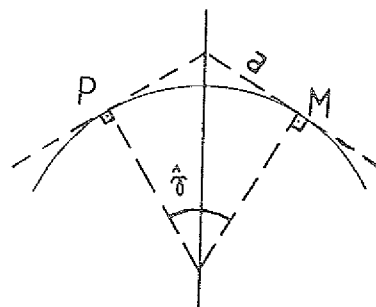


Figure 3: Planar cut of the sphere determined by the points P (tangential point of the plane containing one hexagon), M (centre of one pentagon, vertex of one of the fundamental triangles) and the centre of the sphere.

of the Archimedian pentagonal hexecontahedron (APH), in a way that the resulting hexagons have an inscribed circle tangent to their 6 sides. As a result the pentagons do not lie in a plane tangent to the sphere, but ~ 1.013 higher.

In order to calculate the dimensions of the array we covered the sphere with 60 spherical pentagons, each one having 4 angles of 120° and one of 72° (these are the vertices of the 20 "fundamental triangles"). Because of symmetry properties, the tangential point of the planes on which these pentagons have to be projected to obtain the APH (which are in addition the centers of the circles inscribed in the hexagons) can be easily found using simple geometrical relations.

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The Encapsulated Ge Detector

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The aim of the EUROBALL project [1] is the development of γ -ray detectors with substantially improved qualities, and the realization of large solid angle arrays, in particular of 4π spectrometers employing these detectors. The approach made by all detector developments which are currently considered is to increase the effective Ge volume. This is achieved when using a composite detector. The concept of a composite Ge detector which we developed is the CLUSTER detector described in [1,2]. In addition to the criteria which apply to the design of any state-of-the-art high purity Ge detector, any composite Ge detector has to fulfill two more requirements.

The dead layers between the single crystals of the composite detectors have to be minimized so that the composite detector will behave very different from the ideal large volume Ge detector. This close packing cannot be achieved using a conventional coaxial detector of cylindrical shape. In the case of the CLUSTER, seven crystals of hexagonal tapered shape are used. Building a prototype of a hexagonal tapered Ge detector, we could prove, that the properties of a Ge crystal do not suffer from the deviation from cylindrical shape [1,2].

Using the biggest as well as purest crystals, which can be manufactured currently, makes reliability a serious problem. The small probability for the failure of a single crystal when mounted inside its cryostat increases to an unacceptably high probability for a composite detector of, e.g., 7 crystals. The encapsulation of the individual crystal solves this problem. Each detector element can be tested individually after the encapsulation. Packing several encapsulated detectors into a common cryostat does not affect their function.

The technology of the encapsulation of Ge detectors without destroying their performance was developed in a cooperation between the Institut für Kernphysik and the Zentralabteilung Technologie of the KFA Jülich, the Institut für Kernphysik of the University of Cologne, the Hahn-Meitner-Institut, Berlin, and the French company Intertechnique, Strasbourg.

A first encapsulated coaxial Ge detector was delivered by Intertechnique in January 1991. The detector has a relative efficiency of 25 % and an energy resolution of 2.0 keV at 1.33 MeV using an ac-coupled preamplifier. The ac-coupling was chosen in order to operate the detector capsule on ground potential which facilitates the close packing and the cooling of several detectors in a common cryostat. The capsule is made from aluminum with a wall thickness of 0.8 mm. It is hermetically sealed by electron welding. The vacuum is maintained by

a getter material which is active up to the annealing temperature of 150°C and which is designed for a lifetime of more than 10 years. A special high voltage feedthrough was developed which can be welded into the aluminum capsule. The encapsulated detector was cycled many times between liquid nitrogen and room temperature during 10 months, removed from the cryostat and remounted again without deterioration of the performance. In October 1991 a neutron damage of the detector was induced by a dose of $2 \cdot 10^9$ n/cm² from a fusion-evaporation reaction at the FN-Tandem accelerator of the University of Cologne which deteriorated the energy resolution to 2.5 keV at 1.3 MeV with low-energy tailing of the peak shape. The encapsulated detector was then removed from the cryostat and annealed in an oven for 3 hours at 106°C. By this procedure the original resolution of 2.0 keV and a symmetrical peak shape could be restored.

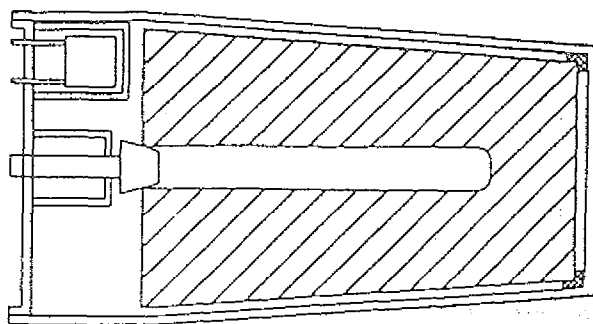


Figure 1: Sketch of the encapsulated Ge detector.

With the development an encapsulated detector the key technology for any design of composite detector became available.

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1. Upgrading to EUROBALL, Technical Proposal with Emphasis on Composite Encapsulated Ge Detectors, edited by J. Gerl and R.M. Lieder, Sept. 1991.
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A new design philosophy for processing of signals produced by high-resolution large volume semiconductor detectors is described. These detectors, to be used in the next generation of spectrometer arrays for nuclear research (i.e. EUROBALL, etc.), present a set of problems like charge trapping and ballistic deficit resolution degradation, low resolution at a high count rate, long term instability, etc. To solve these problems a new design approach has been developed. It includes three new advances:

- (i) Reconstruction of the charge distribution produced by any radiation event. A new digital algorithm has been developed which incorporates as a main process a simple and fast deconvolution.
- (ii) To transfer the problems from an analog to a digital world and to find for them an adequate digital solution, a totally digital signal processing system has been designed. The preamplifier signals are sampled at a very early stage, i.e. before any shaping which otherwise could introduce additional distortion of the charge distribution.
- (iii) To obtain an optimal performance with respect to the energy resolution and the throughput two separate channels with different shaping parameters have been introduced. These channels handle the major aspects of the noise (delta or serial and step or parallel) separately.

Reconstruction of the event charge

The basic elements of the semiconductor system are a detector followed by a charge sensitive preamplifier and a main (linear or shaping) amplifier. Any radiation event produces an amount of charge proportional to the absorbed energy. That charge results in a steplike waveform at the preamplifier output. In the system of our interest, i.e. a high resolution and high throughput system, preamplifiers with resistor discharge are mainly used, so the output step is followed by a decay with a very long time constant compared with the charge collection interval. The preamplifier output signal $U_p(t)$ is described by a convolution between the charge distribution function $g(t)$ and the preamplifier transfer function $f(t)$:

$$U_p(t) = \int_{-\infty}^{+\infty} g(\tau) f(t-\tau) d\tau \quad (1)$$

If the charge collection time is instantaneous $g(t)$ becomes a delta function and the eq. (1) can be rewritten as:

$$U_p(t) = G f(t) \quad (2)$$

where G is the total charge and is proportional to the absorbed energy.

The existing processing systems employ a differentiator, to extract G , followed by a set of integrators. They differ mainly in respect to how the integrations are realized (i.e. active integrators, active integrator plus gated integrator, weighted sum of the integrators outputs [1], etc.). These processing sequence works correctly only when the charge function is a delta one, which has been the condition for the convolution integral (eq. 1) to become a product (eq. 2). Still it works well if the charge collection time is very short compared with any shaping time constant used in the processing channel. In the case of large coaxial Ge detectors the electron-holes drifting time cannot be neglected. The interplay between that time and the shaping constant(s) leads to a degradation of the resolution, known as a ballistic deficit. The problem becomes even more severe for the detectors used for in-beam spectroscopy. The defects produced by the neutron irradiation are playing the role of trapping centers which introduce an additional delay in the charge collection time.

Under such circumstances the substitution of the convolution integral (eq. 1) with eq. 2 is no more tolerable. It means that the processing sequence should be reconsidered. One natural solution is to employ a deconvolution as a first process in a shaping channel. The effect of that approach is obvious — the total amount of charge is extracted from the preamplifier output signal regardless of its distribution on the time scale. As the origin of the ballistic deficit and the charge trapping losses is that time distribution, their influence is eliminated. To perform a deconvolution by analog means is practically impossible. Therefore, our approach involves as a first processing step the digitalization of the preamplifier signal. In the digital world the convolution integral becomes a sum:

$$U_p(i t_s) = \sum_{j=0}^{\infty} g(j t_s) f(i t_s - j t_s) \quad (3)$$

Where t_s is the sampling interval, i.e. the interval of the repetitive digitizing of the preamplifier function $U_p(t)$. If the time scale is normalized to t_s , than the eq. 3 becomes:

$$U_p(i) \Big|_{i=0}^{\infty} = \sum_{j=0}^{\infty} g(j) f(i-j) \quad (4)$$

In fact, eq. 4 is a set of equations to be solved with respect to $g(j)$. This task is difficult to be per-

formed in real time even by a big computer, but some facts help to simplify the process. First, the deconvolution can be performed in a window with a limited size M , which reduces the number of equations in (4) to M . The assumption is correct if the time interval Mt_s is bigger than the maximum possible event charge collection time. This requirement is easy to meet because the total distribution due to the electron-holes drifting and majority trapping delay does not exceed $1 \mu s$ even for very large Ge detectors. Second, for the energy analysis we are interested in the total charge in the window but not in its actual distribution. And third, the preamplifier transfer function $f(i)$ is known or easily definable in real conditions. Even more, for the resistor-discharged preamplifiers it is exponential. Some deviation of the transfer function from the pure exponent could exist due to the nonresistivity behavior of the discharging resistor at high frequencies. But that will not influence the final result, i.e. the total amount of charge, as soon as the window size is bigger than the settling time of the preamplifier response.

These three assumptions permit to solve the equations (4) by employing a simple recurrent procedure:

$$G_i = g(i) \Big|_{i=0}^i = U(i) + (1-k) \sum_{j=0}^{i-1} U(j); \quad (6)$$

where $k = \exp(-t_s/\tau_{pre})$ is the decay constant of the preamplifier transfer function for one sampling interval. At the time instant $i=M$ the accumulated charge G_i becomes G_M and is equal to the total

amount of charge in the time interval from 0 to Mt_s .

Two important features of the recurrence (6) should be pointed out:

- (i) The final result G_M is obtained after M consequent steps, while in each step only one multiplication-accumulation and one summing operation are needed. Each step corresponds to one new sampling result coming from the input digitizer, i.e. the recurrent process could run in real time.
- (ii) The recurrent procedure is easily expandable beyond the M boundary:

$$G_{L+M} = G(i) \Big|_{i=L}^{L+M} = U(i) - U(i-M) + (1-k) \sum_{j=i-M}^{i-1} U(j) \quad (7)$$

i.e. by introducing only one more operation (subtraction) per step the result of the deconvolution in a new time window is achieved. As the parameter L in eq. 7 could be any number it follows that the deconvolution process could run continuously. We call the new processing unit Moving Window Deconvolver (MWD). The block diagram of the unit is given in fig. 1 and it is a direct realization of the eq. 7.

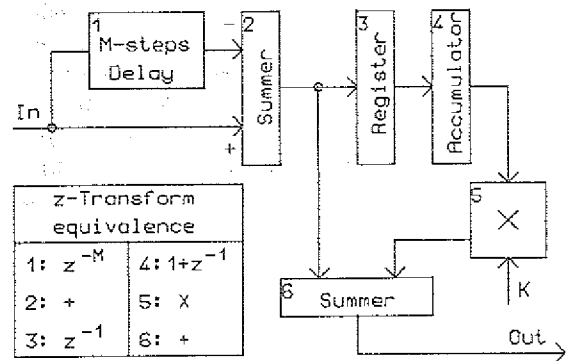


Figure 1: Block diagram of the Moving Window Deconvolver.

Noise shaping

Recently the noise analysis are carried out in the time domain [2,3]. The equivalent detector noise circuit includes two generators — for the serial and the parallel noise. The parallel noise is considered as a random sequence of delta charge pulses, and the serial — as a similar sequence followed by another one with reverse polarity. As the deconvolution in a window extracts all charges (signal and noise as well) belonging to that time interval the size of the window plays an important role. In analogy with the analog system it could be considered as a shaping constant of the digital system.

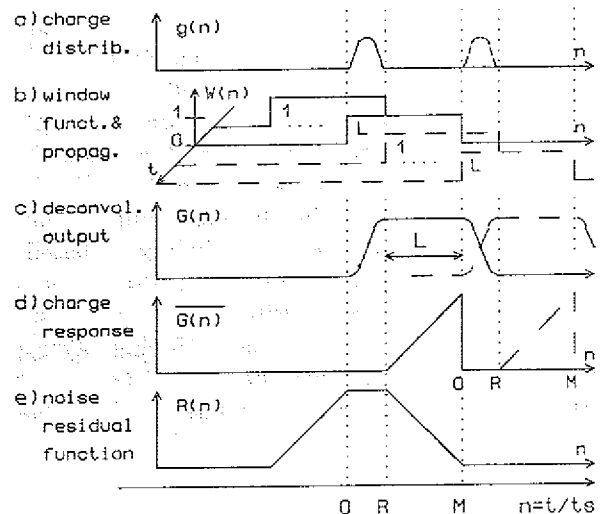


Figure 2: Development of the charge response and noise residual function.

If the window is by L sampling intervals bigger than the charge distribution produced by a single radiation event, than L consequent windows, shifted by one sampling interval, will cover that distribution. It means that L consequent results of deconvolution will represent one and the same event. Taking the mean value of these L results the signal to noise ratio increases. To obtain an analytical expression for the noise contributions the time domain method has been used. On fig. 2d the step

(charge) response of the system is shown produced by shifting $L=M-R$ times the deconvolution window (fig. 2b) and accumulating the results. The zero time reference corresponds to the beginning of the event and the zero to R time interval — to the maximum possible charge collection time. Shifting the delta noise charge, normalized to the event charge, from minus to plus infinity the noise residual function of the system is obtained (fig. 2e). As the residual function is a pure trapezoidal, the parallel and serial noise indexes are:

$$\overline{N_p^2} = \int \{R(t)\}^2 dt = \frac{2}{3} L t_s + R t_s \quad (8a)$$

$$\overline{N_s^2} = \int \{R'(t)\}^2 dt = \frac{2}{L t_s} \quad (8b)$$

Three remarks should be made in connection with eq. 8. First, due to the sampling process the residual function is not smooth but contains steps with a magnitude $1/L$. It is obvious that the parallel noise index will not be affected but the serial will increase. To eliminate that influence a low-pass filter has been placed in front of the sampling ADC with the time constant equal to one sampling interval. Second, the resolution of the sampling ADC should be higher than the signal to RMS noise ratio of the preamplifier output. If it is fulfilled the digitalization noise could be considered as a separate serial noise but the shape of the noise residual function will remain unchanged. The suppression of the digital noise will be discussed later. Third, the parallel noise has been considered as a sequence of charges with a constant rate. This assumption is not correct with respect to the low frequency interference, the detector "microphonic" effect and the temperature induced baseline fluctuations of the detector-preamplifier arrangement. To deal properly with these problems, a separate processing unit named Moving Average (MAU — see fig. 3) has been included in the system. The unit takes the results of the deconvolution in any S windows in which there are no radiation charges.

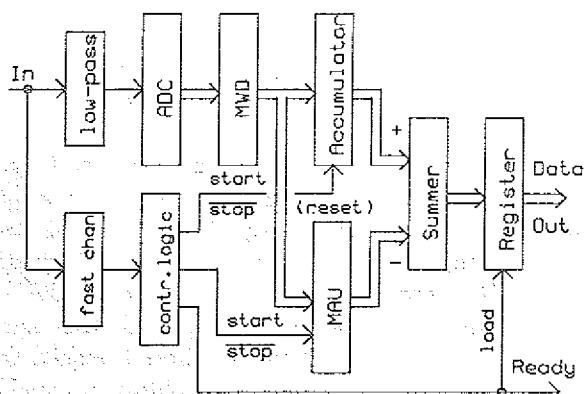


Figure 3: Block diagram of the whole digital pulse processor.

Thus, the unit contains at any moment the mean value of the parallel noise. As this value is subtracted from the total event charge, two terms have to be added to the eqs. 8a and 8b, respectively.

$$\overline{N_{pp}^2} = \overline{N_p^2} \frac{L}{S}; \quad \overline{N_{sp}^2} = \overline{N_s^2} \left(\frac{L}{S}\right)^3; \quad (8c)$$

If the number of windows S is at least 10 times bigger than the number of the charge windows L , than the impact of the MAU over the noise performance of the system is negligible, but still the unit suppresses perfectly the low frequency interferences.

The full pulse processing system is shown in a block form in fig. 3. It is best characterized by assuming 12-bit resolution, 25 MSPS sampling ADC (i.e. 40 nS sampling interval t_s), a MWD window size of 125 samples (i.e. 5 μ S), 100 L -type windows (i.e. 4 μ S effective shaping constant and 1 μ S protection time for full charge collection) and 1000 S -type windows (i.e. 40 μ S shaping constant for the low frequency interferences). The particular features of the design include:

- (i) The charge collection time does not influence the resolution of the system. This comes as a result of the new design philosophy which uses a deconvolution as a first process in the signal processing channel.
- (ii) The noise shaping function is a pure symmetrical triangular (in fact trapezoidal with a short flat top part). That shape is the best choice with respect to an optimal resolution — count rate performance [3,4]. As that shape cannot be realized by analog means in a pure form we consider it the second main achievement of the digital design. In the applications of our interest where the total processing time should be kept below 10 μ S, i.e. where the serial noise is dominant, the slight increase of the parallel noise index, due to the flat top part of the shaper, does not contribute significantly to the overall noise performance of the system.
- (iii) The low frequency disturbances (microphonic effect, 50 Hz interference, etc.) are practically eliminated thanks to the parallel moving average channel. Due to its long shaping constant, compared with the constant of the main channel, it does not virtually change the noise performance of the system.
- (iii) As the full processing is performed in a digital environment a number of advantages emerge: digital control of the processing parameters, reliability, on-line optimization and testing, etc. The final result, i.e. the radiation energy, is produced by the system directly in a digital form. As the system incorporates all necessary processing components we call it a Pulse Processing ADC (PPADC) to stress on its completeness. The utilization of the PPADC as a building unit of a new generation data acquisition system is discussed in ref. [5].

Digital noise

The analog to digital conversion introduces two new sources of error namely quantization error and sampling clock uncertainty error (i.e. clock jitter error). These errors are very low compared with detector parallel and serial noise, but to complete the picture of our digital processing system their contribution will be estimated. A natural way of their representation is to assume a sequence of delta noise charges with RMS amplitudes as follows:

$$(N_d)_{rms} = \frac{2^{-B}}{\sqrt{12}} t_s;$$

$$(N_{jFS})_{rms} = \frac{1}{\sqrt{12}} \frac{t_s}{\tau_{pre}} t_s; \quad (9)$$

where B is the resolution of the sampling ADC in bits. The clock jitter error has a linear amplitude dependence, therefore it is estimated for the worst case, i.e. for the full scale (FS) of the system.

As these errors are randomly positive or negative, they fit well to the serial noise model and their noise indexes can be determined solving the similar integral as in eq. 8b. The results are:

$$N_{dd} = \frac{2^{-B}}{\sqrt{6L}};$$

$$N_{djFS} = \frac{t_s}{\tau_{pre}} \frac{1}{\sqrt{6L}} \quad (10)$$

The quantization noise contribution depends on the resolution of the ADC whereas the clock jitter noise depends on the sampling frequency. To estimate their contribution we will take the actual design parameters (see above) for B, L, t_s and assume a preamplifier with a 40 μ S time constant. The calculated values are both below 40 ppm of the full processing scale. These results show that the digital noise has no practical influence on the performance of the system.

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Introduction

The next generation of high resolution spectrometer arrays for nuclear research like the EUROBALL spectrometer is demanding for a new type of data acquisition and electronic system. Standard solutions are, due to size and power consumption, no longer feasible. This is because in a typical high spin experiment, as described in the EUROBALL proposal [1], the EUROBALL array produces $9 \cdot 10^4$ events/s with an average size of 162 parameters/event. This gives an information rate of about $1.5 \cdot 10^7$ /s and a recording rate of 60 Mbyte/s assuming only 32 bit words are used. In addition to this high data rate one faces the problem that most of the signal processing channels are high resolution Ge-detector channels. The acquisition system has to account for the considerable processing time which is already needed for the high resolution signal processing itself.

The EUROBALL Electronics and Data Acquisition Design Group developed a new concept able to cope with these demands [2-5]. It uses VXIbus [6] (VMEbus Extension for Instrumentation) as the instrumentation system. This system houses the entire data acquisition system, the front-end electronics, the trigger and readout system, including the event builder and formatter. The front-end electronics, though on a high integration level, mainly involves conventional analog signal processing techniques. The very sophisticated trigger system is planned to provide a pipeline mode of operation by internally decoupling the different processing stages.

Parallel to these developments a quite different approach has been worked out by our group at KFA. This alternative concept is featuring extreme simplicity, since only one and only one type of module per detector is necessary to perform the whole signal processing and data acquisition task. Due to the full parallel mode of operation being designed it provides virtually zero system dead-time. In addition it overcomes all problems attributed to conventional analog signal processing, like ballistic deficit and charge carrier trapping effects, pile-up and long term stability etc., since it digitizes the detector signals at a very early stage and then performs all necessary signal processing in a digital environment.

Design considerations:

The new concept was developed on basis of the following main ideas:

1.) In contrast to typical high energy physics experiments with multi parameter detectors, where an efficient trigger plays an important role, in γ -ray spectroscopy experiments with multi detector arrays like EUROBALL the overwhelming majority of detected events, at least photopeak events, is of interest a priori. Therefore the system overhead due to complicated triggers should be reduced as much as possible in favour of maximum throughput. Maximum throughput is achieved by a fully parallel pipeline mode of operation where each individual detector channel operates completely independent from all others. In this way all kind of bottlenecks occurring in standard data acquisition systems can be avoided.

In our design complete parallelism is introduced through an absolute time stamp, derived from a system clock common to all channels, which is used as a tag for every event in any channel. Events marked by such an absolute time stamp can be acquired completely independent from each other even on physically different mass storage devices. Trigger logic decisions are implemented with highest possible flexibility since they are made off-line in such a system. Only a simple but versatile real-time majority logic circuit and individual ADC gates remain as optional on-line triggers. If necessary, they may provide for instance multiplicity selection and Compton suppression, which are the most important triggers for front-end data reduction in γ -ray spectroscopy.

2.) The γ -ray spectrometer arrays of the next generation like EUROBALL will provide a resolving power, which is improved by a factor 10^3 - 10^4 compared to existing devices. Large volume Ge-detectors or composite Ge-detectors giving an even larger combined volume like the CLUSTER detector play a key role in getting this improvement. However, large volume Ge-detectors are very sensitive to ballistic deficit and charge carrier trapping effects hence less resistant against neutron damage. Various methods to compensate for these effects and to maintain the detector performance under in-beam conditions as long as possible have been developed recently [7-9]. However, all of them involve analog circuits. Therefore they exhibit all the well known disadvantages related to analog electronics. The low noise immunity, temperature and long term stability, as well as the usually fixed and never ideal transfer functions of analog filters

are some of them. Moreover, the modules available up to now in general are difficult to adjust and applied to signals which are already imperfectly shaped by the preceding analog shaping amplifier stages. The latter also explains why even a careful adjustment never can yield a perfect result.

In our design a virtually true ballistic measurement of the detector signal is performed using a PPADC (Pulse Processing Analog to Digital Converter). The PPADC is an own development described in a separate contribution to this report [10]. It represents the main component of our data acquisition system since it combines most of the analog signal processing as well as the complete digitizing part normally used in conventional systems. The PPADC allows a virtually true ballistic measurement because it samples the detector signal immediately after the preamplifier with adequate resolution and sampling frequency. After the preamplifier the original charge distribution of the detector signal is convolved only by the preamplifier transfer function which is very well known and, since close to an ideal integrator, simple to define. In the digital environment given after the sampling process, it is therefore easy to extract the original charge distribution from the sampled preamplifier signal by applying a simple deconvolution algorithm. The final digital integration over the full charge distribution or charge collection time yields the total charge of the event as a result of a virtually true ballistic measurement. The residual function of the adaptive digital filtering performed by the system is optimized to give the highest possible noise-rate factor [11]. Due to the ballistic measurement being performed, our design of a data acquisition system based on the PPADC not only eliminates all problems connected to ballistic deficit and charge carrier trapping effects, which are of similar nature, it also provides automatic pile-up rejection and adaptive filtering, i.e. dynamic shaping automatically adjusted to each individual signal. For a given signal-to-noise ratio the latter allows an about 100 % higher peak throughput per channel than a conventional system with analog pulse processing. In addition it guarantees an excellent temperature and long term stability since the only remaining analog component in our system is the preamplifier.

3.) EUROBALL as well as other future arrays will consist of several hundred detectors. Most of them are Ge-detectors requiring high resolution signal processing and acquisition channels. In conventional electronics such a channel consists of at least 6 modules, namely spectroscopy amplifier, timing filter amplifier, constant fraction discriminator, time to amplitude converter and two ADC's for energy and time. In addition to the signal processing itself, various trigger logic modules, many analog and logic delays for timing adjustments, an event-builder and formatter etc. are necessary to build up the complete data acquisition system. Due to its complexity such a system cannot be setup

and controlled manually anymore. Even a computer controlled operation may be problematic because of the big variety of different modules. Despite the fact that for most of the modules computer controlled versions have not been developed yet, it can be expected that the service and maintenance with respect to software as well as hardware will be a difficult task.

The data acquisition system according to our design features extreme simplicity since it employs only one module per channel. Moreover, this module, the PPADC, is of exactly the same type for all high resolution signal processing channels. For the scintillator signals a simplified and cheap low resolution version of the PPADC may be used. The modification only means that the high resolution sampling ADC and the signal processor, which are the main components of the PPADC, have to be replaced by a standard Flash ADC and accumulator. Due to the digital signal processing techniques exclusively involved, full computer control of all signal processing parameters is a natural part of the system design. The effort to adjust these parameters is reduced to a minimum since many of the signal processing functions like baseline restoration, ballistic deficit correction, pile-up rejection etc. are performed automatically by hardware or by the pulse processing algorithm used. As the only task for the system setup the rather simple timing adjustment of the real-time multiplicity logic and the Compton suppression logic circuitries remain which are optional, however. The serial data read-out is performed via standard transputer links between each PPADC and the host computer(s). In this way high noise immunity and transfer speed is achieved while system overhead due to complicated bus systems and protocols is avoided. The system initialization, the DSP software download, the parameter setup and on-line control etc. can be done via any small bandwidth hardware standard like Ethernet or CAMAC.

For the final data recording a farm of Exabyte drives or some CREO optical tape drives may be used. If one wants to record in a typical EUROBALL experiment the full rate of valid Ge-detector events, i.e. all singles events after pile-up rejection and Compton suppression, the event size of 64 bit provided by the PPADC module and the expected total singles rate of 1.5 million events per second requires a recording bandwidth of 12 Mbyte/s. Taking into account the transfer speed of the two different types of mass storage devices one would need a farm of 24 Exabyte drives or 4 optical tape drives allowing a continuous data taking of 3 hours or 90 hours, respectively. However, the amount of required mass storage hardware can easily be reduced either by setting appropriate multiplicity thresholds on the individual channels or by intermediate hard disk buffering of the data to be recorded. Fig. 1 schematically shows the PPADC based data acquisition system design for EUROBALL.

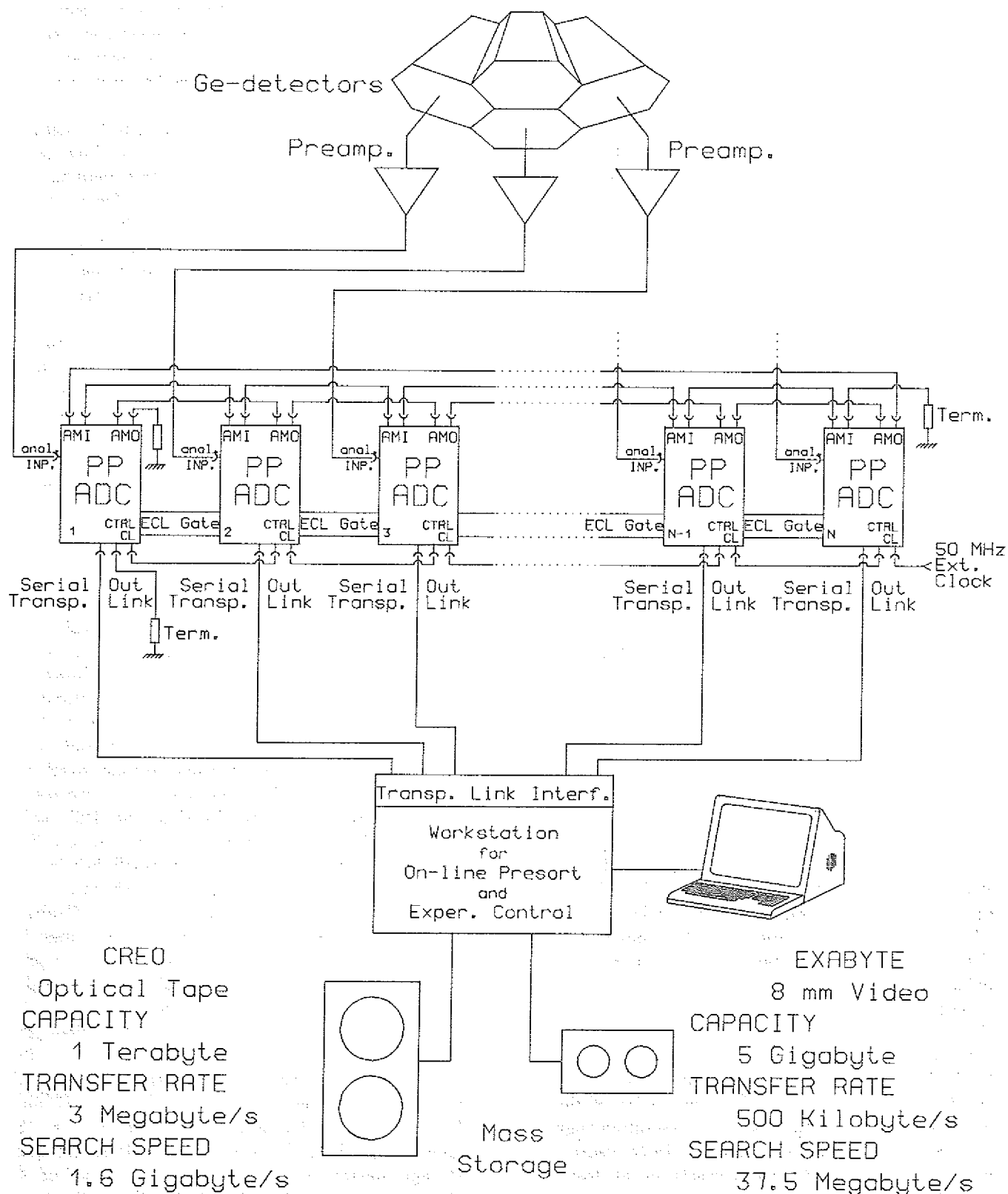


Figure 1: Schematic view of the PPADC based data acquisition system design for EUROBALL.

The PPADC data acquisition module

As outlined before there is only one type of module needed to perform the whole signal processing, the analog to digital conversion, the data acquisition, and the main part of the trigger logic tasks. Nevertheless this PPADC based data acquisition module has a rather simple internal structure since it takes full advantage of the digital signal processing introduced by the PPADC. As shown in the schematic diagram of Fig. 2 it consists of two main parts:

a) The right hand side represents the internal structure of the PPADC itself. The main components are the sampling ADC, the digital signal processor (DSP) and the moving window deconvolver (MWD) which is removed from the DSP algorithm and realized in hardware in this design.

when no signal is present. The switching between the two modes is controlled by the timing control unit (TCL). This unit also generates start and stop bits to tag the events and to check for pile-up conditions, and it enables the serial readout of the actual value of the counter (CNT) representing the absolute system time, into the event FIFO. In addition it checks for the occurrence of a veto signal during the time the FIFO is filled with the samples of one event. If a veto occurs and the hardware trigger mode is enabled, it resets the FIFO synchronously to the arrival of the stop bit. In this way all the samples belonging to the vetoed event are canceled. The hardware veto signal is provided by the gate control unit (GCL) which itself extracts it from the fast trigger logic described later.

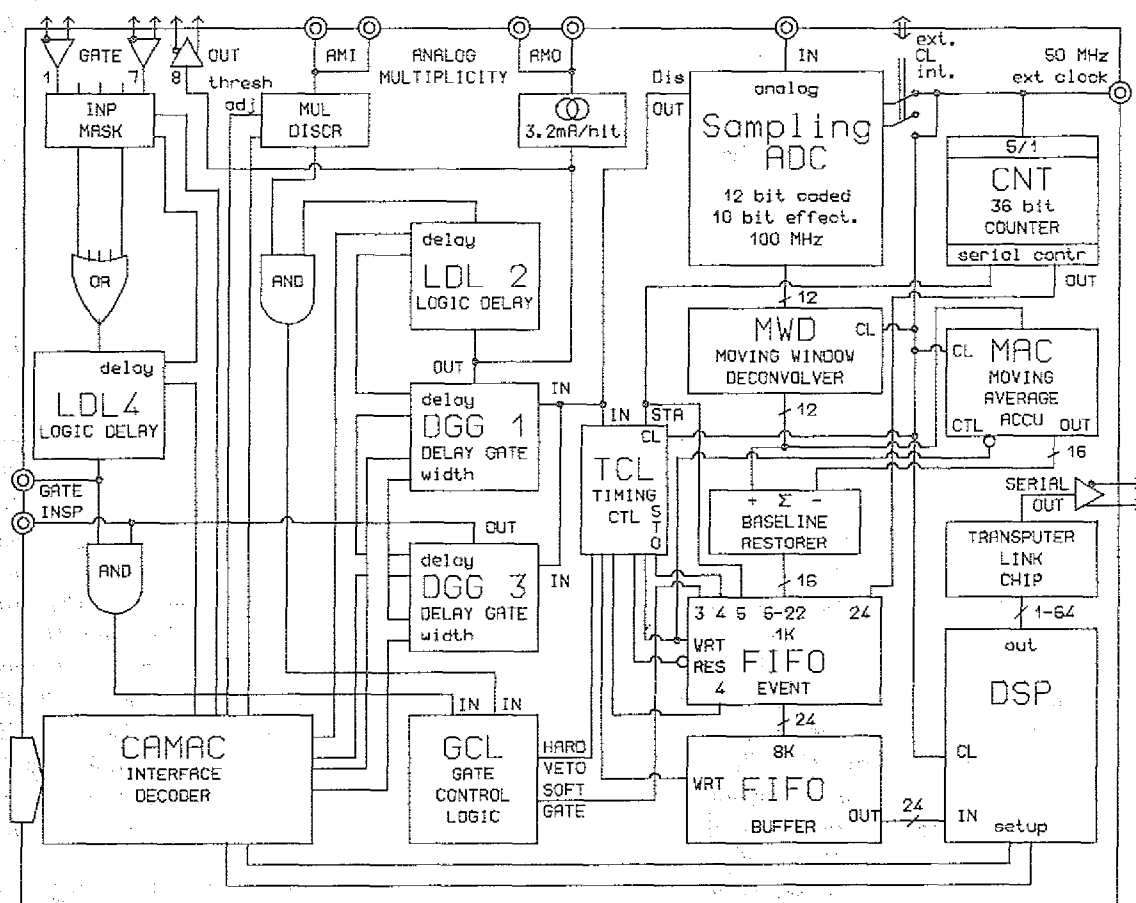


Figure 2: Block diagram for the internal structure of the PPADC data acquisition module.

The results of the moving window deconvolution of the sampled preamplifier signal which continuously and with sampling speed appear at the output of the MWD are fed via an automatic baseline restoration circuit into an event FIFO as soon as the presence of a signal has been detected by the fast discriminator of the sampling ADC. The baseline restoration circuit, represented by the moving average accumulator (MAC) and a subtractor, is continuously updated during the time

This technique avoids the necessity of a precise veto timing and provides high gating flexibility since the only requirement which remains is, that the veto signal has to appear within the time of typical 5 μ s, during which the samples of one event are collected.

The automatic hardware pile-up rejection works in a similar fashion. If a second pulse appears before a preset minimum number of samples is collected for the previous one a reset of the FIFO is generated

synchronously with the stop bit of the pile-up pulse. The hardware pile-up rejection can be switched off to enable a detailed pile-up inspection by the DSP process.

One or more additional bits of the event FIFO open the possibility to process gating tags together with each event. They can be used later by the DSP process or even off-line for putting software gate conditions.

The whole process described up to now is running in real-time with a sampling speed of 25–50 megasamples per second. Hence, except the unavoidable pulse integration time which is needed to obtain a sufficient signal-to-noise ratio, there is no additional dead-time introduced by the system. However, depending on the amount of further analysis the DSP has to perform, it may not be able to run at such a speed. Therefore a deep buffer FIFO is placed between the event FIFO and the DSP which derandomizes the incoming data stream. A derandomization at such an early stage is possible since after the event FIFO all tasks which have to be done in real-time are already performed. As a result the DSP can run completely asynchronous at its maximum speed. This means for instance: if five instructions are needed to analyze one sample, and the initial event rate is 30 Kcps with a mean number of 200 samples per event, then the DSP has to have a speed of just 30 Mips in order to make the system virtually dead-timeless.

The buffer FIFO not only improves the system performance with respect to the DSP dead-time. It also unloads all the following stages of the data acquisition. However, after the DSP the situation is much more relaxed, since in the standard setup the DSP algorithm converts the information, initially represented by ca. 200 samples of 16 bit per event, into one 64 bit word per event containing the system clock (36 bit), the total energy of the event (16 bit) and its time relative to the system clock (9 bit), and the software gate information (3 bit). The final data rate of 240 Kbyte/s obtained under these conditions for the assumed initial event rate of 30 Kcps can easily be handled by the transputer link interface and the current mass storage technology so that no further system dead-time is introduced.

b) The left hand side represents the hardware trigger logic part. It uses standard logic circuits and provides two kinds of hardware trigger options, multiplicity selection and external gating by up to seven different sources. All trigger logic decisions are made by fast real-time coincidences between the timing discriminator output of the sampling ADC and external fast ECL signals, thus providing a coincidence resolving time equal to that of conventional fast trigger logic modules. This and the before mentioned FIFO buffering of the digitized detector signal explains, why even quite late

triggers, at least on a conventional fast logics time-scale, will not introduce any dead-time, so that the whole trigger logic circuit can be considered as virtually dead-timeless as well.

The design philosophy of the multiplicity logic circuit has been adopted from the Le Croy Model 4532 Majority Logic Unit. The fast discriminator signal of the sampling ADC is first shaped in width by a delay and gate generator (DGG 1), to determine the multiplicity coincidence resolving time, and then converted to an analog multiplicity information by the 3.2 mA current source. The resulting analog multiplicity output AMO is used for daisy chaining the analog multiplicity information within all modules of the system by simply adding the corresponding currents. The logic delay of the DGG 1 module helps to compensate for the different timing properties of the various detector channels and for the cable delays. After passing the last module in the chain the total analog multiplicity information is fed back to the analog multiplicity input AMI of the module, which again is daisy chained with the analog multiplicity inputs of all other modules in the same order as the outputs. A multiplicity discriminator compares the total multiplicity input signal with the preset multiplicity threshold and generates the corresponding logical signal if it is exceeded. Finally, an autocoincidence circuit represented by a simple AND logic between the multiplicity discriminator output and the DGG 1 output checks whether the module by itself contributed to the incoming multiplicity event. If not, or if the multiplicity threshold was not exceeded, and the hardware multiplicity selection was enabled, the gate control logic (GCL) disables the readout of the actual event via the hardware veto line and the timing control unit. Alternatively, in the case that this hardware option was not enabled, the event may be tagged via the gate line with the actual multiplicity condition for further processing of the multiplicity information. The logic delay line (LDL 2) is used to compensate for the total output-to-input delay of the chain in order to provide the proper coincidence overlap.

By this design a hardwired, but nevertheless very flexible multiplicity logic circuit running in real-time, is established. Each module can operate under its individual hardware or software multiplicity conditions. Simply by setting the appropriate multiplicity thresholds one module for instance can run in singles mode whereas the neighbouring one processes events only if a five-fold coincidence condition is fulfilled. The setup of the multiplicity logic circuit is very simple, too. It is restricted to the adjustment of two delays, where one is determined by the total delay in the chain and therefore identical for all modules. The adjustments may be computer controlled if the appropriate interface, for instance CAMAC as in fig. 2, is provided.

The concept of the external gating logic is straightforward, too. Seven individual external ECL gate signals may be either accepted or masked off by the gate input mask. The masking feature is only included in order to simplify the cabling. One and the same cable can be used to distribute the necessary gating signals within one group of detectors, for instance a CLUSTER, by daisy chaining it from one module to the other. Those modules which need a certain gate signal, pick it up from the line, whereas the others mask it off. This method has the additional feature that the hardware gate conditions can be modified by software to a certain extent. The gate signals which pass the input mask are adjusted in time by the logic delay LDL 4. The validity of the gate condition is checked by a fast coincidence between the output of the delay and the discriminator signal of the module, which beforehand is shaped and timed by the delay and gate generator DGG 3 in order to fix the coincidence resolving time. Inspect outputs of both signals allow to control the adjustments by an oscilloscope. The output of the AND logic representing the gate coincidence again is fed into the gate control logic GCL where it can be decided, whether the signal will be used as a hardware gate or veto, or whether it will be processed as an event tag to provide a software gate option.

As it was discussed already for the multiplicity logic, after the fast real-time gate coincidence again the timing of the gate logic is uncritical and needs no further adjustment. This, together with the input mask feature, makes the setup of the gate logic an easy task. And it makes the choice of the gating conditions quite flexible, although they are virtually hardwired and therefore dead-timeless. At least for the main gating tasks in nuclear spectroscopy helping to reduce the front-end data rate, like Compton suppression, particle-gamma coincidences etc., the flexibility is by far sufficient. More elaborate on-line triggers would unnecessarily complicate the system and decrease the throughput. They can be set off-line in our system since for all events the absolute time of their occurrence is measured. However, the optional gating tags could for instance be used for the on-line experiment monitoring. The analyzer software could perform a fast on-line presort of the data according to these tags and increment the result in different gate spectra. In this way a simultaneous monitoring of gate spectra for the most important gate conditions can be done on-line.

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Elastic Transfer in $^{11}\text{B} + ^{12}\text{C}$ System at 344.5 MeV (Lab)

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The backward angle rise of the elastic and inelastic scattering cross sections observed for some heavy ion systems can be interpreted as being due to the large contribution from the elastic and inelastic transfer reactions. Investigations of such processes is a particularly good tool for the determination of spectroscopic amplitudes of the transferred particle.

In a study of this reaction for the $^{11}\text{B} + ^{12}\text{C}$ system in a broad energy range from 8 MeV to 41 MeV (CM) it was found that the value of the spectroscopic factor determined from model analysis of the experimental data deviated considerably from the theoretical value showing a remarkable strange energy dependence¹⁾. While at the lowest energies its value is close to the Cohen and Kurath theoretical prediction²⁾, it shows a rather steep rise to a value approximately twice as large in the vicinity of 10 MeV (CM) and then a much slower decrease again to a value slightly larger than theoretical one at 41 MeV, the highest energy investigated up to now (Figure 1). The spectroscopic factors for inelastic states $^{12}\text{C}(\text{g.s.}) = ^{11}\text{B}(2.12) + p$ and $^{12}\text{C}(4.44) = ^{11}\text{B}(\text{g.s.}) + p$ did not show such energy variation although their values exceed by ca. 60% the theoretical prediction.

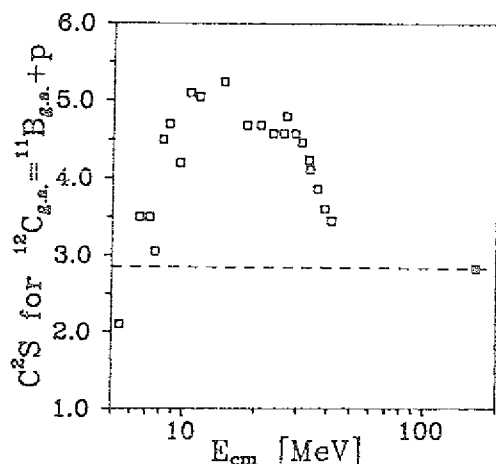


Figure 1:
Energy dependence of $^{12}\text{C}(\text{g.s.}) = ^{11}\text{B}(\text{g.s.}) + p$ spectroscopic factor. The squares are the data from ref. 1), full square - from the present work, dashed line - the theoretical Cohen-Kurath value.

In order to extend the data to much higher energies an experiment was performed at 344.5 MeV (lab) using the ^{12}C beam from the JULIC cyclotron. The detection system consisted of semiconductor counter dE-E telescopes allowing the charge

identification of ejectiles and thus the measurement of the angular distributions of elastically scattered ^{12}C ions as well as ^{11}B particles corresponding to the backward scattering region dominated by the elastic transfer. Figure 2 shows the obtained results for the elastic scattering.

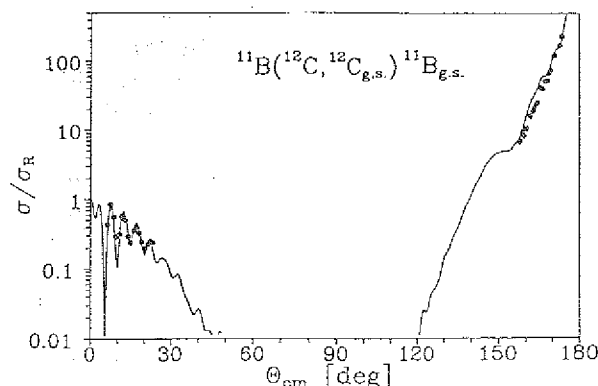


Figure 2:
Experimental angular distribution for elastic scattering $^{12}\text{C} + ^{11}\text{B}$. Solid line - a coherent sum of optical model and one proton transfer DWBA calculations

The analysis of the angular distributions in the elastic transfer region of elastic scattering was performed using the DWBA model with the waves in entrance and exit channels distorted by the optical potential with Wood-Saxon form factors found in the previous optical model analysis of $^{11}\text{B} + ^{12}\text{C}$ elastic scattering³⁾. The potentials binding the proton with the ^{11}B core in ^{12}C were assumed of Wood-Saxon shape with geometrical parameters $R = 1.25 \cdot A^{1/3}$ fm, $a = 0.65$ fm and the depths adjusted to obtain the appropriate binding energy. As can be seen from Fig. 1 the value of the proton spectroscopic factor at the high energy 167.8 MeV (C.M.) of the present work agrees well with the theoretical prediction.

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Inelastic Scattering to the Continuum in Collisions of ^{12}C ions with ^{12}C , ^{58}Ni and ^{90}Zr Nuclei at 344.4 MeV (LAB)

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In an extended study of nuclear reactions induced by 344.5 MeV ^{12}C ions, based on experiments performed formerly at the JULIC cyclotron, strong transitions of different reaction channels to the continuum were observed indicating the contribution of preequilibrium processes¹⁾.

Recently, theoretical approaches to a quantitative treatment of preequilibrium reactions have been developed based on the well established framework of direct reaction and compound nucleus theory. In this investigation an attempt was undertaken to apply the statistical multistep direct reactions (MSDR) developed in ref.²⁾ to our experimental data on inelastic scattering to the continuum of ^{12}C ions. This formalism is an extension to the continuum of methods which have been successfully used for describing the transitions to isolated states. Statistical assumptions and averaging procedures leading to incoherence of different reaction processes are used while preserving the coherence of the collective features of excitation and of relative motion. It is this property which makes the process calculable in terms of averaged direct multistep reaction amplitudes.

The formalism was applied to the inelastic scattering of ^{12}C ions on even-even ^{12}C , ^{58}Ni and ^{90}Zr nuclei. In these cases one can expect that random phase approximation (RPA) is able to provide a realistic description of the structure of these nuclei. In the present work the single particle wave functions of a shell model potential of Wood-Saxon form were used. The RPA transition densities have been approximated by the average density. The optical model potentials were taken from the previous analysis of elastic scattering data from the same experiment³⁾ with a slightly readjusted depths of the imaginary part. The form factors were calculated as derivatives of the optical model potentials and the deformation lengths are described microscopically with QRPA response functions. The transition strength functions include spreading widths parameterized according to analysis of damping of the giant resonances.

In this formalism the energy averaged double differential cross sections have been calculated for the reactions under investigation taking into account one-step and two-step processes, since the contributions from three- or more step-processes are negligible. The results are compared with

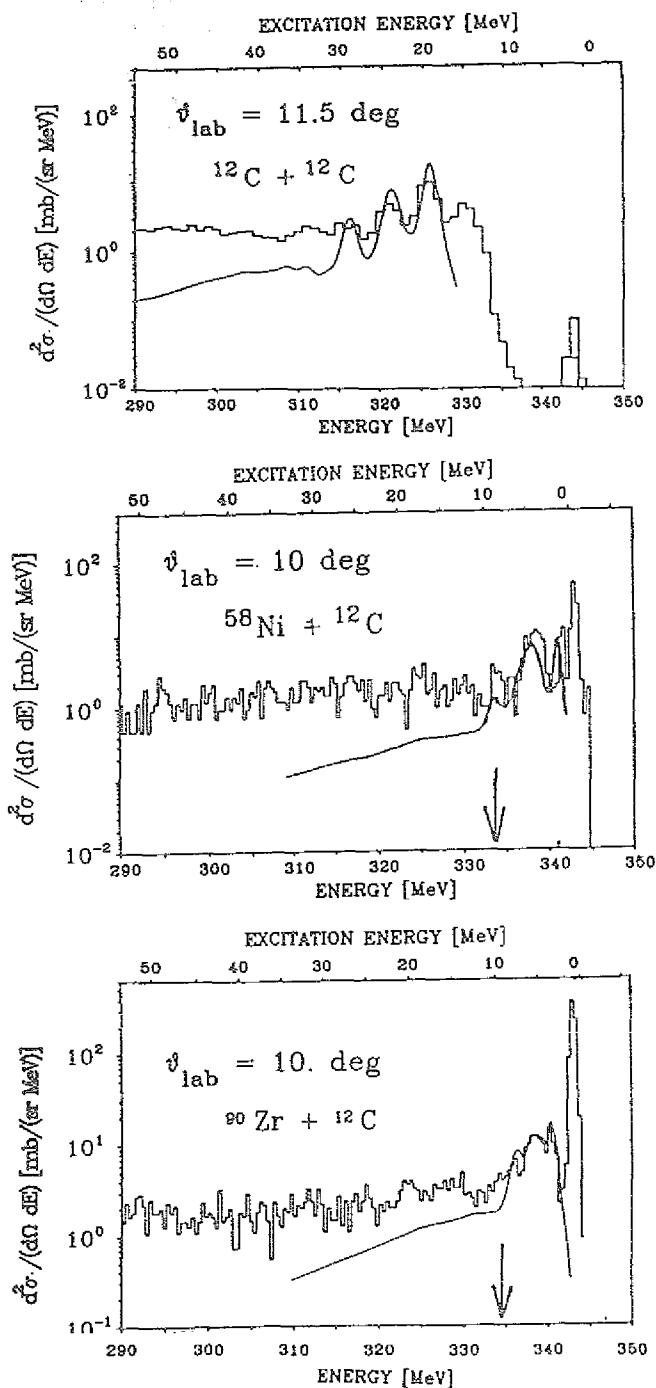


Figure 1:
Comparison of the theoretical spectra of inelastically scattered ^{12}C ions on ^{12}C , ^{58}Ni and ^{90}Zr nuclei calculated in MSDR formalism with experimental results. The arrows indicate energy at which the "side-feeding" starts

experimental spectra in Fig. 1. The calculations reproduce well the continuous part of spectra corresponding to the lower excitation energies up to 15, 8 and 6 MeV for ^{12}C , ^{58}Ni and ^{90}Zr , respectively. The higher excitation energy part of the spectra contains contributions from the multi-step excitations not taken into account in the present analysis e.g. the nucleon decay of the $(A+1)$ nuclides originating from nucleon transfer to the target or projectile nucleus of mass A . The threshold energies of this "side feeding" spectra agrees well with the energy at which the theoretical inelastic cross section calculated in the MSDR formalism starts to deviate strongly from the experimental data.

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Two Photon Correlations from $^{16}\text{O}+^{24}\text{Mg}$ Interactions at 33 MeV·A

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The space-time evolution of the interaction region between two heavy ions or a proton-antiproton system has been studied by measuring two particle correlation functions. These particles have been so far either nucleons, light ions or pions. A recent review is given in ref. [1]. These probes suffer from strong interactions in the final channel with their environment, i.e. an excited system of hot hadronic matter.

In this contribution we have chosen photon-photon correlations in order to study the space-time evolution of the $^{16}\text{O}+^{24}\text{Mg}$ -system at a bombarding energy of 33 MeV·A. A ring consisting of 50 lead glass modules was employed as shower detector. This ring was also used to study subthreshold π^0 -productions [2]. The two-photon correlation function R was thus measured as a function of the four momentum difference Q :

$$Q^2 = -(p_1 - p_2)^2 \quad (1)$$

with p_i the four momentum vectors of the two γ -quanta. In this case where the rest masses are zero Q reduces to the invariant mass. Uncorrelated events were generated through event mixing.

In Fig. 1 is the two-photon correlation function shown. It is normalized in the region $Q \geq 200$ MeV/ c^2 where no correlation is expected to exist.

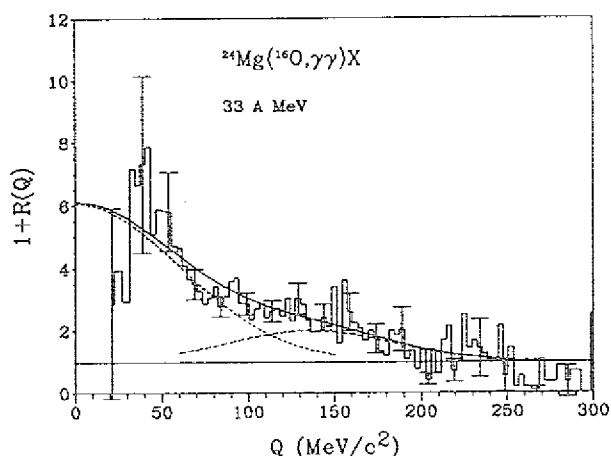


Figure 1: Correlation function (histogram with selected error bars) for two γ -emission. The solid curve is a fit with two Gaussians to the data (see text). The two components are also separately shown (dashed curves).

An enhancement around $Q = 135$ MeV/ c^2 is attributed to π^0 -decay. In order to study whether the enhancement for low Q -values stems from the Dalitz decay $\pi^0 \rightarrow e^+e^-\gamma$ we performed Monte Carlo simulations. These simulations indicate that the origin of the low Q -enhancement must have a different origin. Since our experiment was only sensitive to high energy photons we conclude that the enhancement may be due to Bose-Einstein quantum-statistics [2]. We have analysed the data further by assuming a Gaussian distributed source in space and time and by approximating the photons by plane waves. We then find for the correlation function

$$R(Q) = 1 + \lambda e^{-\lambda(Q/Q_0)^2} + p_0 e^{-(Q-135)^2/2\sigma^2} \quad (2)$$

with Q_0 the space-time extent of the source, λ is an empirical factor to adjust the height at $Q=0$. The last term takes the π^0 -decay into account. This fit yielded $Q_0 = 83.2$ MeV/ c .

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Fragment-Fragment Correlations at Fermi Velocity

Jayanti Rama Rao¹, H. Machner, and G. Buhr

Over the energy range from a few MeV per nucleon to a few GeV per nucleon, the nuclear reaction mechanism undergoes a gradual-transition from the predominance of mean-field phenomena and Pauli blocking effects at low energies to the dominance of individual nucleon-nucleon collisions and baryonic excitations at higher energies. An interesting milestone in the transitional region is the Fermi energy per nucleon (~ 30 MeV/A) where the nucleons in the target as well as the projectile have similar velocity distribution at their collision in the centre of mass system.

It is the aim of the present work to study the collisions of ^{16}O nuclei at 32 MeV/A with ^{197}Au nuclei in order to make an assessment of the reaction mechanism from an exclusive study of fragment-fragment correlations in the reaction plane. For this purpose we employed a bank of six counter telescopes on either side of the beam direction at relative angles between 20° and 80° . With a true-to-chance ratio better than 50, we studied the fragment-fragment correlations in a series of nine coincidence combinations. A preliminary comparison of singles and coincidence spectra (as in Fig. 1) shows a broad similarity in the spectral shapes leading to the inference that the fragmentation in collisions at Fermi velocity may

well be statistical to a large extent. This observation emphasizes the importance of meanfield effects and apparent validity of statistical thermodynamical models in the interpretation. Detailed analysis is in progress.

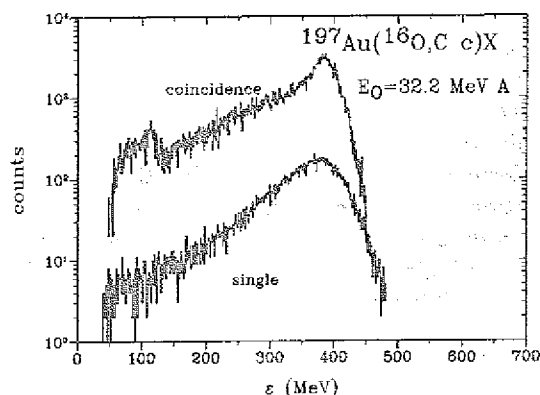


Figure 1: Singles and coincidence carbon spectra at 12° .

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H. Machner and M. Nolte

The energy regime between 10 to 100 MeV·A in heavy ion collisions is a transitional region. Different patterns in these reactions have been attributed to different underlying reaction mechanisms [1]. Much of such work was inspired from the participant-spectator picture originally applied to relativistic energies. Another interesting idea brought into the field is the assumption of a phase transition from liquid to vapor [2]. One sign of this transition should be multifragmentation. Also a power law was found for the intensity of fragments. In this study we have measured intermediate mass fragments (IMF's) emerging from the interactions of ^{32}S with ^{197}Au target nuclei at 33 MeV·A bombarding energy. The energies and charges of the fragments were measured with a 3-fold telescope consisting of an ionization chamber, a 300 μm thick Si-detector and a 5 mm Si(Li) detector. In addition we measured the time of flight over a flight path of approximately 1 m with a channel plate detector and the first silicon detector yielding fragment masses. Data were taken at a laboratory angle of 17° .

In Fig. 1 the mass distribution is shown ranging from 4 (α -particles) up to 56.

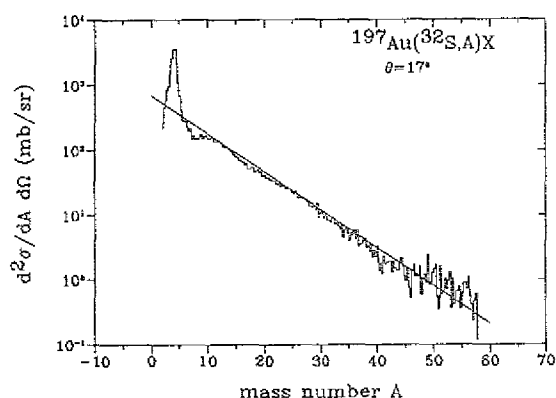


Figure 1: Mass distribution for the indicated reaction as obtained from the TOF measurement. Also shown is an exponential fit (straight line).

Similar to the charge distribution the mass distribution is better fitted by an exponential than by a power law. The exponent of such a fit follows the systematics of Ref. [3]. Because of the poor time resolution achieved in the experiment the isotopes were not separated. In order to extract isotopic yields we first sorted mass spectra for different elements. This mass spectra were then fitted by Gaussian where we varied intensities and assumed one width for each element which was also fitted.

In Fig. 2 the mass distributions for fragments ranging from helium to nitrogen are shown together with the fits.

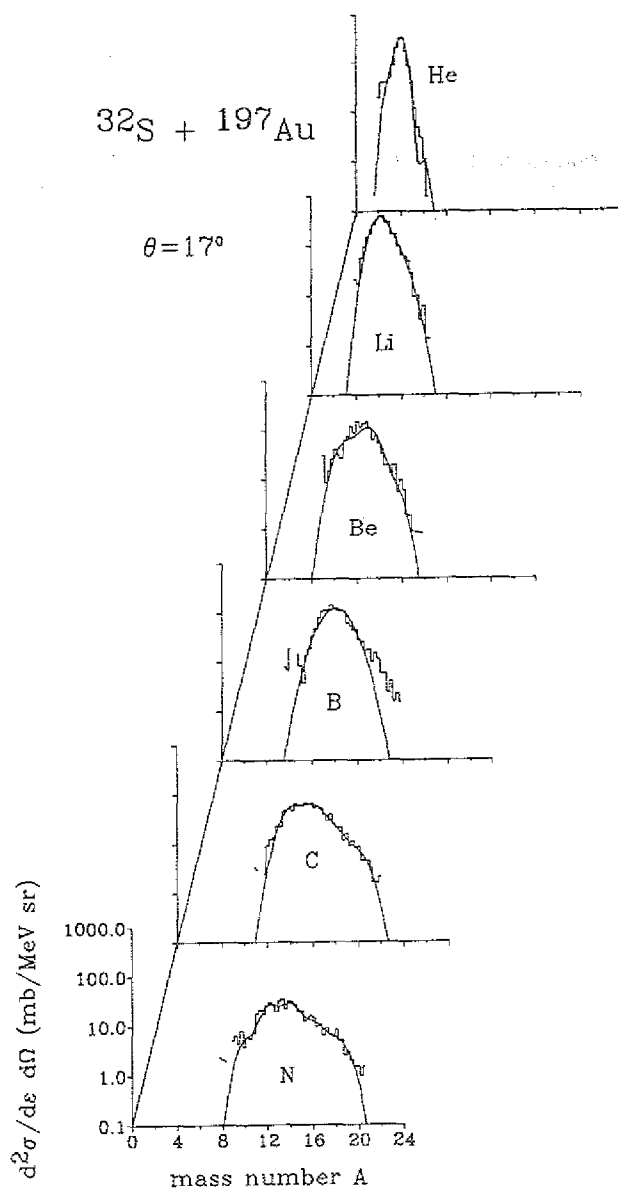


Figure 2: Isotope distributions for the indicated elements. Data are shown as histograms and fits as solid curves.

With the help of these fits isotopic cross sections were extracted and shown in Fig. 3. The most abundant isotope is ^4He . As expected from the mass yield curve (Fig. 1) there is a decrease in yield with increasing mass. For the lighter elements there is a pronounced structure in the abundances which

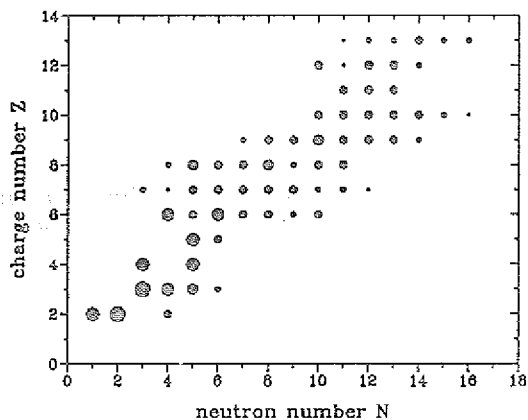


Figure 3: Extracted cross sections for IMF's. The symbols sizes correspond to cross sections in mb/sr: <5, 5-10, 10-50, 50-100, 100-500, and >500 which is only ^4He .

is not necessarily the one of occurrence in nature. There is a general trend to favour $Z = N$ isotopes. In summary we have measured isotopic yields of IMF's from an intermediate energy heavy ion reaction. The mass yield curve extends far over the projectile mass without any structure in this area. This may be a further hint to a unified reaction mechanism at these energies as found in a recent study of energy spectra [4].

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III. Theoretical Nuclear Physics

Within the meson exchange picture, we compute the electromagnetic form factor of the pion and show that it is sensitive to the elementary strong $\rho\pi\pi$ form factor. This sensitivity may provide a measure of the confinement region.

The starting point for our calculation is a QCD motivated effective meson Lagrangian such as that of Bando et al. [1]. A nice feature of this Lagrangian is that the photon coupling is included in a gauge invariant way, enabling us to compute, for example, the electromagnetic form factor of the pion. The important new ingredient for us is a diagram in which the photon transforms to a ρ -meson. In the usual vector dominance picture, the pion electromagnetic form factor would then be obtained from the photon transforming to a ρ which then decays to two pions. By using a Breit-Wigner propagator for the ρ to account for its width and adjusting the overall strength, one can obtain agreement with the data in the ρ resonance region but it is not possible to simultaneously fit the ρ resonance and the normalisation requirement $F_\pi(t=0) = 1$. However, in the meson exchange model of the $\pi\pi$ system [2], the $\pi\pi$ self energy corrections to the ρ means that the physical ρ propagator is more complicated than a simple Breit-Wigner.

If we write $\Gamma_{\gamma\rho}$ for the amplitude arising from the Lagrangian for the process $\gamma \rightarrow \rho$, d for the bare ρ propagator and f_B for the bare $\rho \rightarrow \pi\pi$ vertex, then, up to appropriate normalisations, the π electromagnetic form factor is written (in operator form) as

$$F_\pi = \Gamma_{\gamma\rho} d f_B + \Gamma_{\gamma\rho} d f_B G T, \quad (1)$$

where T is the coupled channel t -matrix [2] and G is the two-body propagator. The first term is just the diagram usually written down in the vector dominance picture, except that now the ρ propagator is simply the Feynman propagator for the bare ρ appearing in the Lagrangian. The second term is the final state interaction that dresses this bare ρ giving it its physical mass and width.

Since we take the $\gamma \rightarrow \rho$ interaction to be pointlike (i.e., it has no form factor associated with it) the only new parameter introduced is the $\gamma \rightarrow \rho$

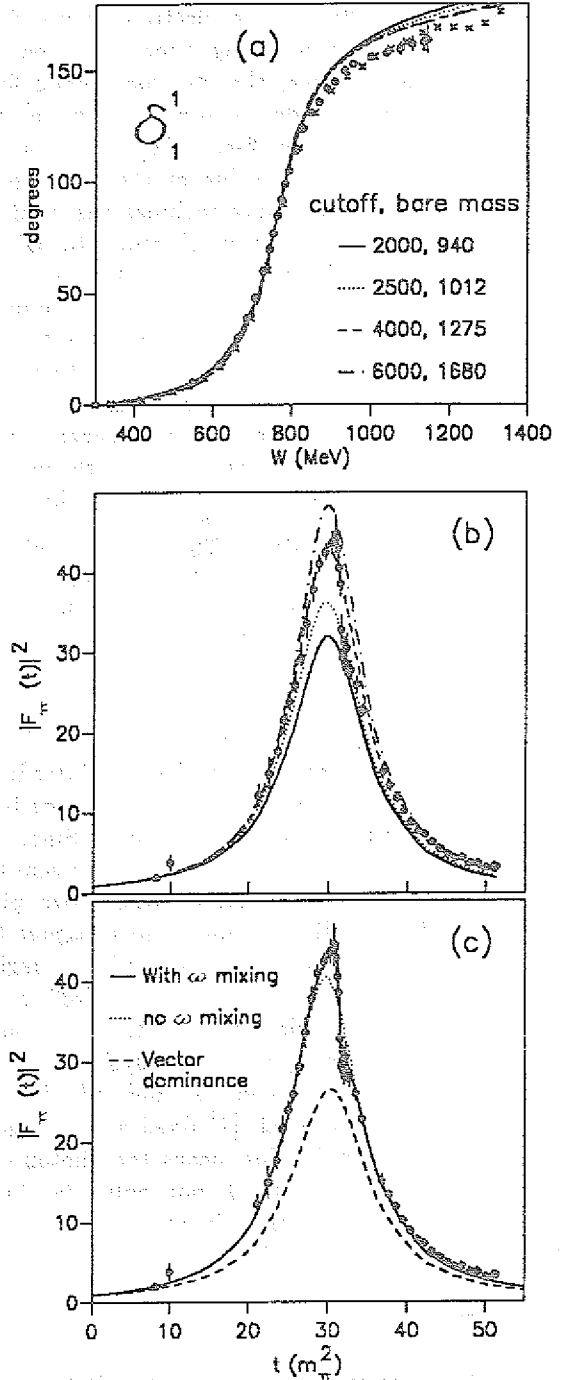


Fig. 1. (a) Dependence of the $\pi\pi$ phase on the $\pi\pi\rho$ cutoff mass. (b) Corresponding dependence of the electromagnetic form factor. (c) π electromagnetic form factor including ρ - ω mixing.

coupling constant. This is fixed by requiring $F_\pi(0) = 1$. However, it turns out that the $\pi\pi$ results are somewhat insensitive to the bare $\pi\pi\rho$ cutoff mass. This is illustrated in Fig. 1 (a), where we show the $\pi\pi$ phase shifts in the ρ channel for four values of the cutoff mass. In each case, the bare ρ mass is readjusted to fit the data but all other parameters are kept fixed. In Fig. 1 (b) we show the corresponding results for the electromagnetic form factor, where we see a strong dependence on the $\pi\pi\rho$ cutoff mass. Hence, the electromagnetic form factor essentially provides a measurement of the off-shell $\pi\pi\rho$ vertex. In Fig. 1 (c) we show the results using the optimal value of the cutoff mass (3300 MeV). The solid curve includes the addition of a ρ - ω mixing contribution obtained by adding to the calculation the usual Breit-Wigner vector (ω) dominance contribution. The coupling constants are fixed by known decays ($\omega \rightarrow \pi\pi$ and $\omega \rightarrow e^+e^-$), with only the phase adjusted to fit the data. If this ω contribution is neglected, we obtain the dashed curve. The dotted curve is a Breit-Wigner with the ρ mass and width and the strength adjusted to ensure $F_\pi(0) = 1$ (the naive vector dominance model). The full calculation results in an r.m.s. radius of 0.647 fm. This is much closer to the experimental value [3] of 0.66 ± 0.01 fm than the value 0.62 fm obtained using the vector dominance model.

It is an assumption of the meson exchange approach that the form factors that are introduced are a means of parametrising the underlying quark-gluon structure of the hadrons. Since the electromagnetic form factor seems to be so sensitive to the form factor at the vertex where two pions couple to a bare ρ , it is tempting to interpret this as a measure of the size of the confining region. Yielding to such temptation, the required cutoff mass of 3300 MeV would imply an r.m.s. radius of only 0.15 fm. However, in the analysis so far we have kept the parameter a defined in the Lagrangian of Bando et al. [1] fixed at the "vector dominance" value of 2. This means the photon does not couple directly to the π but only via the ρ . Relaxing this constraint introduces some ambiguity, making such a prediction less reliable.

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A recent calculation [1], based on dispersion integrals, indicated that the scalar radius of the nucleon is surprisingly large, resulting in a significant contribution to the nucleon σ -term from the pion cloud. Here, we show that a recent meson exchange model of the $\pi\pi$ system [2] leads to a similar conclusion. A remarkably unified picture results of $\pi\pi$ and $K\bar{K}$ scattering, the $N\bar{N} \rightarrow \pi\pi$ process and the pion, kaon and nucleon scalar form factors.

The nucleon σ term, $\sigma(t)$, can be defined in terms of the matrix element [1]

$$\sigma(t)\bar{u}(p')u(p) = \frac{m_u+m_d}{4m_N} \langle p' | (\bar{u}u + \bar{d}d) | p \rangle \quad (1)$$

A low energy theorem relates $\sigma(t)$ defined in this way, to the isospin even πN amplitude at the Cheng–Dashen point [1]. The interest in the σ term was generated when analysis of the πN data resulted in a value of 64 ± 8 MeV [3] yet naive estimates based on baryon mass differences gave $\sigma(0) \approx 30$ MeV. It was recently pointed out [1] that the strong s-wave $\pi\pi$ interaction produces a significant t dependence of $\sigma(t)$. That is, the usual assumption that $\sigma(2m_\pi^2) \approx \sigma(0)$ may not be valid.

To investigate this effect using the coupled channel $\pi\pi$ – $K\bar{K}$ model developed by the Jülich group [2], we consider $\sigma(t)$ to be given by (using operator notation)

$$\sigma = \sigma_B + \Gamma_B G\nu + \Gamma_B GTG\nu. \quad (2)$$

Here, σ_B represents the "bare nucleon" σ term. We do not attempt to calculate this contribution, concentrating instead on the t dependence of the remaining pieces. The next two terms involve the matrix element of $\bar{u}u + \bar{d}d$ between pion (or kaon) states. The bare part (i.e., without $\pi\pi(K\bar{K})$ rescattering which we put in explicitly) of this process is denoted by Γ_B . The pions (kaons) then connect to the nucleon directly or via multiple scattering. Also, G is the $\pi\pi(K\bar{K})$ propagator, T denotes the (off-shell) $\pi\pi$ – $K\bar{K}$ t -matrix [2] while the elementary $\pi\pi(K\bar{K}) \rightarrow N\bar{N}$ amplitude is denoted

by ν . The parameters of ν were adjusted so that $\nu + TG\nu$ reproduces the semi-empirical analysis [4] of the $\pi\pi \rightarrow N\bar{N}$ process in the unphysical region.

Embedded in Eq. (2) is the σ term of the pion (kaon), $\Gamma(t)$, which is given by

$$\Gamma = \Gamma_B + \Gamma_B GT. \quad (3)$$

For $t=0$ this can be expressed as derivatives of the pion (kaon) masses with respect to the quark

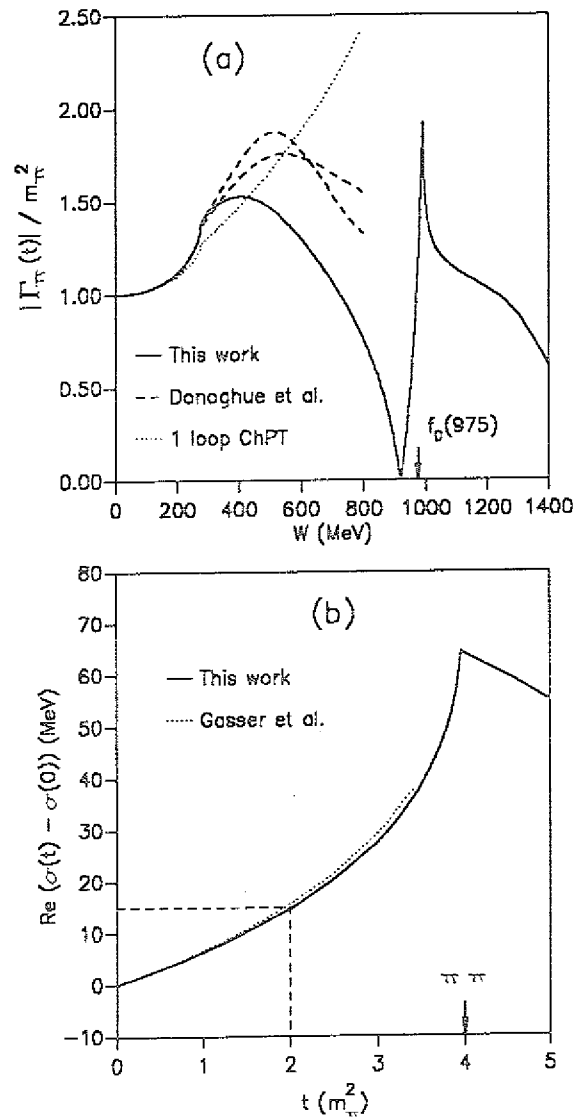


Fig. 1. (a) The pion scalar form factor. The dashed and dotted curves are from Ref. [5]. (b) t dependence of the nucleon σ term. The dotted curve is from Ref. [1].

masses [5]. We use this prediction to fix the two coupling constants that appear in Γ_B . This means we have no free parameters in the calculation.

The result for the pion form factor $\Gamma_\pi(t)$ is shown in the solid curve of Fig. 1(a). The effect of the $f_0(975)$ (which emerges in the model of Ref. [2] as a $K\bar{K}$ bound state) is clearly evident. The r.m.s. radius is 0.75 fm. For comparison, we also show the results of a one-loop chiral perturbation theory calculation [1,5] (dotted curve) and, in the dashed curve, a recent calculation by Donoghue, Gasser and Leutwyler [1,5].

Having fixed the coupling constants and confirmed that our calculation of $\Gamma(t)$ is sensible, we can proceed to the nucleon σ term. In Fig. 1(b) we show the difference $\sigma(t)-\sigma(0)$ in the vicinity of the Cheng-Dashen point (assuming the contribution of the "bare nucleon" σ term, σ_B is approximately a constant function of t). Our results are in almost exact agreement with the dispersion relation calculation of Gasser, Leutwyler and Sainio [1] (dotted curve) throughout the range of their calculation. (Note that embedded in our calculation is the same unitarity relation that drives their dispersion integrals.) In particular, we confirm their result that there is a strong t dependence of $\sigma(t)$ near the Cheng-Dashen point, giving rise to a 15 MeV contribution to $\sigma(2m_\pi^2)-\sigma(0)$ and a scalar r.m.s. radius of 1.22 fm.

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Effects of final-state interactions in the $\bar{p}p \rightarrow \pi\pi$ annihilation process

V. Mull, K. Holinde and J. Speth

Recently, we have presented a detailed study [1] of the annihilation of the nucleon—antinucleon ($\bar{N}N$) system into two mesons (all combinations of $\pi, \eta, \eta', \rho, \omega, \delta, a_1, a_2, f_2$ and K, K^*). The calculations have been done in a DWBA approach, using nucleon (N), delta-isobar (Δ) and hyperon (Λ, Σ, Y^*) exchange as transition mechanisms together with a realistic initial state ($\bar{p}p$) interaction. Although in most cases the trend of the data could be reproduced, some specific discrepancies occurred. One possible reason is that, so far, the final-state (meson—meson) interaction has been neglected. The main obstacle for a realistic inclusion of such effects is that, in general, the meson—meson interaction is largely unknown, especially in the required energy range above $\bar{N}N$ threshold. An exception constitutes the $\pi\pi$ channel, where experimental phase shifts and inelasticity parameters are known up to relative high energies.

In the present work [2] we have investigated the role of the final-state interaction in the $\bar{p}p \rightarrow \pi^+\pi^-$ process in detail using a microscopic model for the $\pi\pi$ interaction [3], which has been constructed by the Jülich group using likewise the conventional meson—exchange framework. Apart from providing a definite off-shell behavior, this model gives an excellent description of scattering observables and phase shifts for energies up to 1.4 GeV. We extrapolate it without any parameter modification to the energies ($\sqrt{s} \geq 1.85$ GeV) needed in the present analysis. The resulting phase shifts and inelasticities are in qualitative agreement with the empirical data, apart from those in the S-wave. It remains to be investigated whether the results will change if a model is used which describes the data quantitatively.

Fig. 1 shows that the inclusion of the $\pi\pi$ interaction has a large effect: it reduces the integrated $\bar{p}p \rightarrow \pi^+\pi^-$ cross section by a factor of about 2. A quantitative description can again be achieved by a suitable adjustment of form factor parameters in the basic transition process. The effect is even more dramatic in the differential observables, especially in the analyzing powers. Here, the pronounced angular structure at higher energies is obtained in

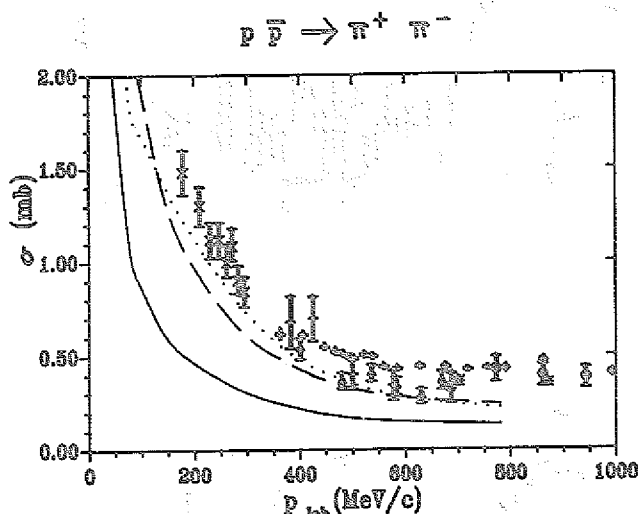


Fig. 1. Effects of the final state interaction on the total cross section. The dotted curve represents the calculation of Ref. 1, the dashed curve is the result with final-state interaction, but without any parameter readjustment, and the solid curve shows the result after a slight readjustment.

our model only after the inclusion of the $\pi\pi$ interaction. Moreover, it brings the model predictions nicely in line with experiment, especially in forward direction.

Furthermore, we have investigated the relative importance of partial waves. Surprisingly, the final-state interaction has a considerable effect not only in the S-wave but also in the D wave, in which the phase shifts are not far from 180° .

In summary, we have demonstrated that, in spite of the large momentum carried by the pions annihilating from the $\bar{N}N$ system, the effects of the $\pi\pi$ interaction is sizable and by no means negligible. Correspondingly, there is no reason to leave out interactions for the $\pi\rho, \rho\rho$ etc. channels occurring in the annihilation process. Therefore, reliable conclusions about the relevance of explicit quark—gluon dynamics can only be drawn after all final-state effects have been included in a realistic way.

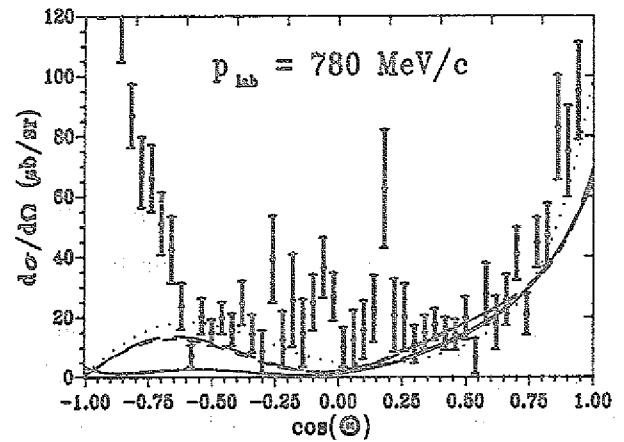
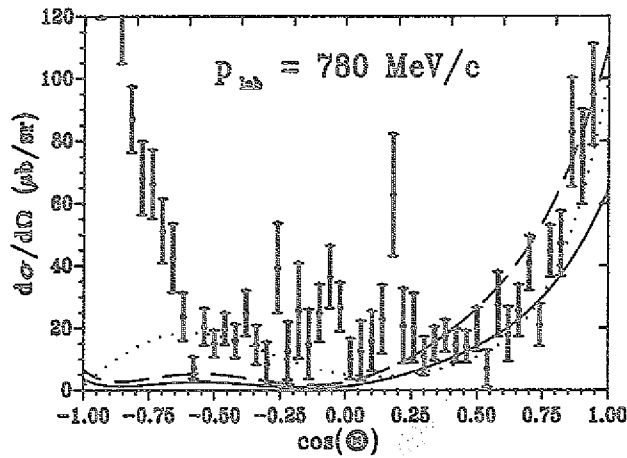


Fig. 2. Effects of the final-state interaction on differential cross section (a) and analyzing power (b). The description of the curves is the same as in Fig. 2.

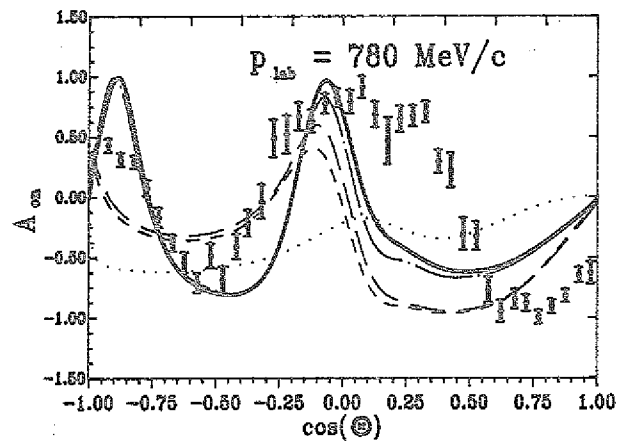


Fig. 3. Final-state interaction is included stepwise in different partial waves: dotted line: no FSI, short-dashed: FSI in s-waves, long-dashed: FSI in s+p, dash-dotted: FSI in s+p+d waves, solid: FSI in all partial waves.

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Can $N\bar{N}$ data tell us something about the nature of the NN repulsion?

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An interesting topic of current research is the question about the physical origin of the strong repulsion occurring between two nucleons at short distances. It is here that the quark-gluon substructure of hadrons is expected to manifest itself. Indeed if, in a meson-exchange model, the whole repulsion is parametrized in terms of ω -exchange (as done in the Bonn-potential [1]), the required strength $g_{NN\omega}^2(t=0)$ is about 10, i.e. roughly twice as large as the value obtained from SU(6) symmetry, which is $g_{NN\omega}^2 = 9 \cdot g_{NN\rho}^2 \approx 5$. Although these symmetry relations need not to be correct, they work surprisingly well for all other coupling constants, also in the strange sector [1,2]. This fact immediately suggests that the NN repulsion is only partly due to ω -exchange; the other part might be generated by different mechanisms, e.g. pomeron exchange [3], pair terms [4] or one-gluon-exchange between (antisymmetrized) quarks [5]. It is

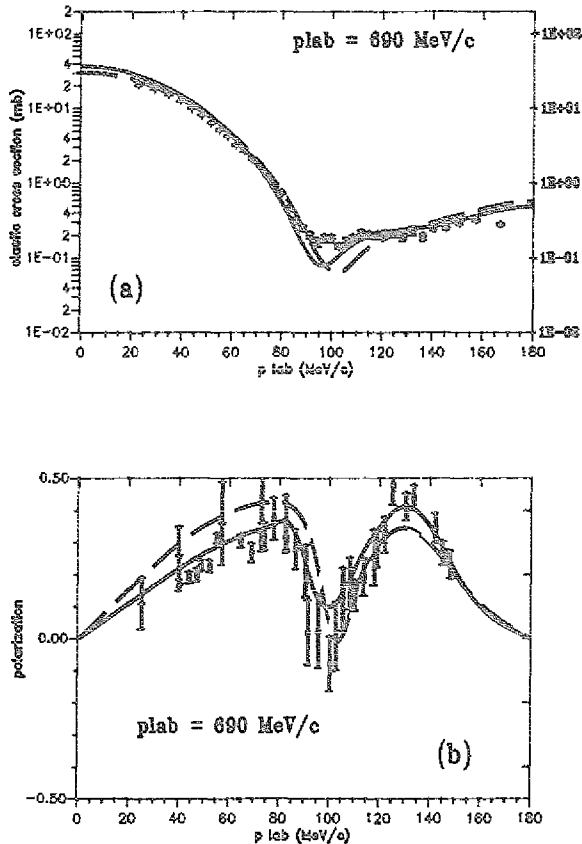


Fig. 1. Elastic differential cross section (a) and elastic polarization (b) for $g_{NN\omega}^2 = 20$ (solid) and $g_{NN\omega}^2 = 0$ (dashed). For references to the data, see ref.6.

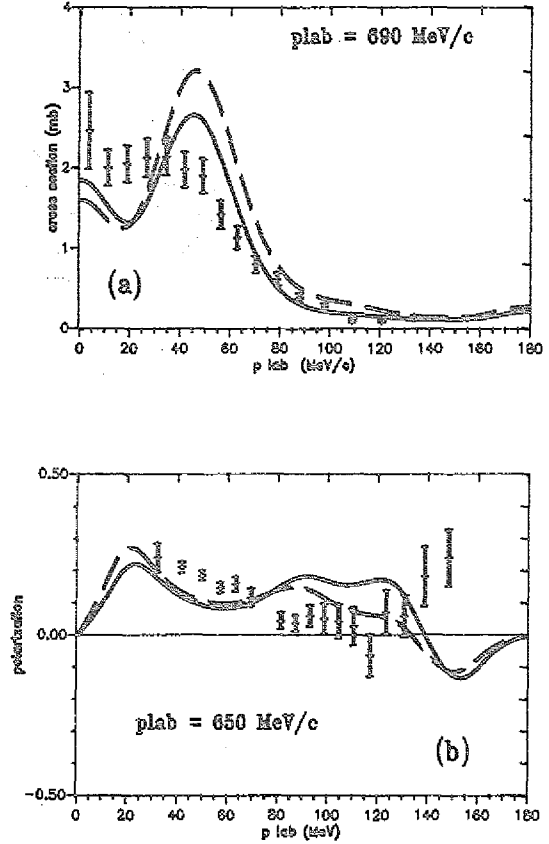


Fig. 2. The same as fig.1, for the charge-exchange channel.

interesting that all these "exotic" contributions stay repulsive when going from the NN to the $N\bar{N}$ system, whereas the ω exchange changes sign, i.e. different scenarios for the short-ranged NN repulsion can lead to quite a different (elastic) $N\bar{N}$ interaction in the inner region. Therefore, the $N\bar{N}$ system should provide an ideal testing ground in order to decide the question whether the whole NN repulsion is generated by ω -exchange only. In fact, historically, this was one of the main motivations for studying the $N\bar{N}$ system.

In the present work, we investigate the question by comparing the predictions of two $N\bar{N}$ interaction models which have the same (phenomenological) annihilation part but differ in the (short-ranged) elastic part: In the first model (A(BOX)[6]) the elastic part is obtained by a G-parity transformation of the (full) Bonn NN potential and thus

contains the strong attraction generated by the ω -exchange. In the second model, ω -exchange is eliminated completely; this would be realistic, at least for the inner region, if indeed half of the NN repulsion is generated by mechanisms which stay repulsive in the $N\bar{N}$ system.

It turns out that predictions for integrated cross sections as well as observables in the elastic channel (fig. 1) are quite similar. The reason is simply that the $N\bar{N}$ reaction is dominated by annihilation processes. Therefore, the data are rather insensitive to variations in the (short-ranged) elastic part. Some differences occur in the charge-exchange channel (fig. 2) which is possible since here the results are generated by the differences of two large isospin amplitudes. Still, in view of the present uncertainties in the data and, even more important, in the theoretical description of the annihilation, there is, in our opinion, little chance, at least in the near future, that the $N\bar{N}$ system can help to clarify the issue of the nature of the short-ranged repulsion.

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The Nucleon—Nucleon Potential in the Skyrme Model: Beyond the Product Approximation

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Calculations of the nucleon—nucleon potential in the Skyrme model use the product ansatz [1] to simplify the numerics. With this ansatz the two-skyrmion potential as a function of the relative angles of orientation between the skyrmions has a compact form, and the extraction of the nucleon—nucleon potential is straightforward. A major inadequacy is the lack of an intermediate range attraction in the central potential. Although many remedies have been proposed, this may not be due to a fault of the Skyrme model. The product ansatz, which is not a solution to the equations of motion, can only be considered accurate at large distances.

We go further by studying the two—skyrmion potential through exact numerical calculations [2]. This allows a more meaningful comparison with phenomenological nucleon—nucleon potentials. The general form for the two—skyrmion potential is

$$V(r, C) = \sum_{j=0}^{\infty} \sum_{m=0}^j V_{jm}(r) \frac{1}{2} [D_{m,m}^{(j)}(C) + D_{-m,-m}^{(j)}(C)] \quad (1)$$

where $C = A^\dagger A_2$ is an $SU(2)$ matrix given by $C = \exp(i\tau_3\alpha/2)\exp(i\tau_2\beta/2)\exp(i\tau_3\gamma/2)$ which determines the relative isospin orientation between two solitons. The nucleon—nucleon potential is then obtained by sandwiching Eq. (1) between the asymptotic two—nucleon states

$$\psi(N_1, N_2) = 2D_{i_1, -s_1}^{(1/2)}(A_1) D_{i_2, -s_2}^{(1/2)}(A_2) \quad (2)$$

Since $D_{m,m}^{(j)}(C) = \sum_n D_{m,n}^{(j)}(A_1^\dagger) D_{n,m}^{(j)}(A_2)$, one finds that terms with $j \leq 2$ can contribute only through intermediate states. The product ansatz generates a potential consisting only of terms with $j \leq 2$. Previous calculations have shown, however, that the coefficients V_{2m} are small with respect to V_{1m} and V_{00} , and that the potential derived in the product—ansatz is fit well by truncating at $j=1$. One need then perform calculations for only three different orientations C which is especially helpful. The calculations have been performed by discretizing the equations of motion on a $20 \times 20 \times 40$

spatial lattice and relaxing from an initial $B = 2$ configuration for fixed separation. We have studied configurations with orientations $C = 1, i\tau_2$, and $i\tau_3$. As compared to the product ansatz we see a better agreement between the exact calculations and the Paris potential at intermediate and short ranges than was found for the product—ansatz. In particular, the intermediate range central attraction, missing in the product—ansatz, is present here (Fig.1). The behavior at short—range for the present potential is quite different from that derived from the product ansatz and qualitatively similar to the Paris potential.

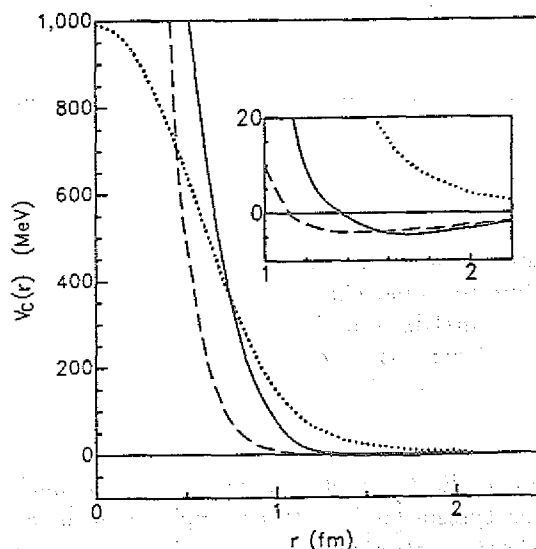


Fig. 1. The central part of the nucleon—nucleon interaction for the present calculation (solid line), the Paris Potential [3] (dashed line) and the product ansatz [1] (dotted line).

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Improved Boundary Condition, Coupled Channel Model: pp Predictions to 1 GeV

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In describing the nucleon-nucleon interaction beyond the 150 – 200 MeV cm kinetic energy (nucleon beam kinetic energy $T_L > 300 - 400$ MeV) inelasticity becomes very important. Pion production begins at $T_L = 280$ MeV, but is relatively weak until Δ production threshold, $T_L = 630$ MeV, is approached. However, the effect of inelasticity on the elastic channel begins well below the threshold because of the virtual role of the higher mass intermediate state. To describe the NN interactions in the range $0 \leq T_L \leq 1$ GeV a coupled channel system must be used. In that range the $\Delta\Delta$ and $NN^*(1440)$ channels are sometimes of importance, in addition to $N\Delta$ channels.

The Boundary Condition Model (BCM), by using the power of R-matrix theory, incorporates short-range degrees of freedom, quarks and gluons, through the energy-poles of the f-matrix (inverse R-matrix) [1]. At long range a meson exchange potential is used, consistent with the confinement property of QCD. Coupling to isobar channels is easily introduced by a potential matrix and the f-matrix in the context of coupled Schrödinger equations [2]. The decay width of the isobar is introduced through a mass distribution of channels [3].

In the applications of this model to the nucleon-nucleon system [1–3] with an emphasis on predicting the elastic scattering, the potentials connecting to isobars were simplified. In the nucleon-nucleon sector π , η , ρ , ω and two- π meson exchanges were included, sufficient to describe the interaction outside of the boundary range (~ 1 fm). For the transition potentials between the two-nucleon and the isobar sector, the one-pion exchange potential was used and, when needed, a phenomenological two-pion range potential to represent more massive meson exchanges. No potentials were used between isobar channels, but the diagonal f-matrix in the isobar channel sector was varied. It was expected that this short-range interaction would adequately represent the effect of meson exchanges for the indirect contribution to the elastic scattering. A good quantitative description of most of the nucleon-nucleon phase parameters was obtained [2,3]. The important exceptions were the detailed shapes of the inelasticity, η , in the 1D_2 and 3F_3 channels. The predicted strength and width of these structures is good indicating that they are threshold effects, without the need for significant contributions from quark degrees of freedom. In fact the BCM has been the only model

able to explain the strong onset of inelasticity in the 3F_3 channel. As $\eta(^1D_2)$ and $\eta(^3F_3)$ are dominant in producing the experimental structures in the polarized total cross sections $\Delta\sigma_L$ and $\Delta\sigma_T$, these structures are not well-represented when the shape of the inelasticities is not accurate.

We now include the OPE contribution to the diagonal $N\Delta$ potentials of both the direct $N\Delta \rightarrow N\Delta$, and exchange, $N\Delta \rightarrow N\Delta$, diagrams. Although $g_{\pi N\Delta}^2 \gg g_{\pi NN}g_{\pi\Delta\Delta}$, this is compensated by the much larger spin matrix elements of the direct, compared to the exchange diagram. Their contributions are comparable in magnitude and vary in sign. Consequently there is considerable variation in the sign and magnitude for the resulting spin-spin and tensor potentials.

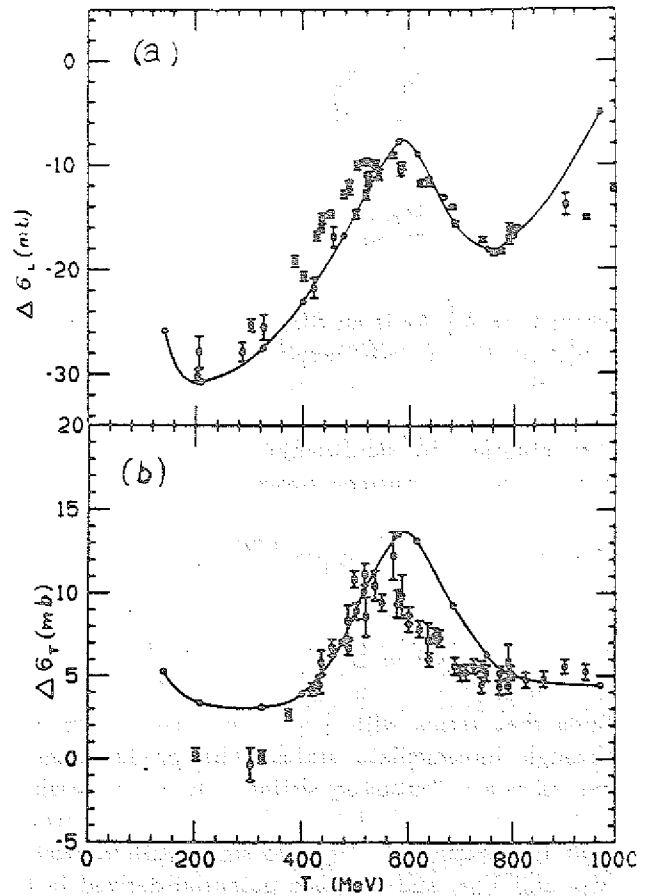


Fig. 1. Polarized total cross sections. The prediction of the improved model compared to the data for (a) $\Delta\sigma_L$ and (b) $\Delta\sigma_T$.

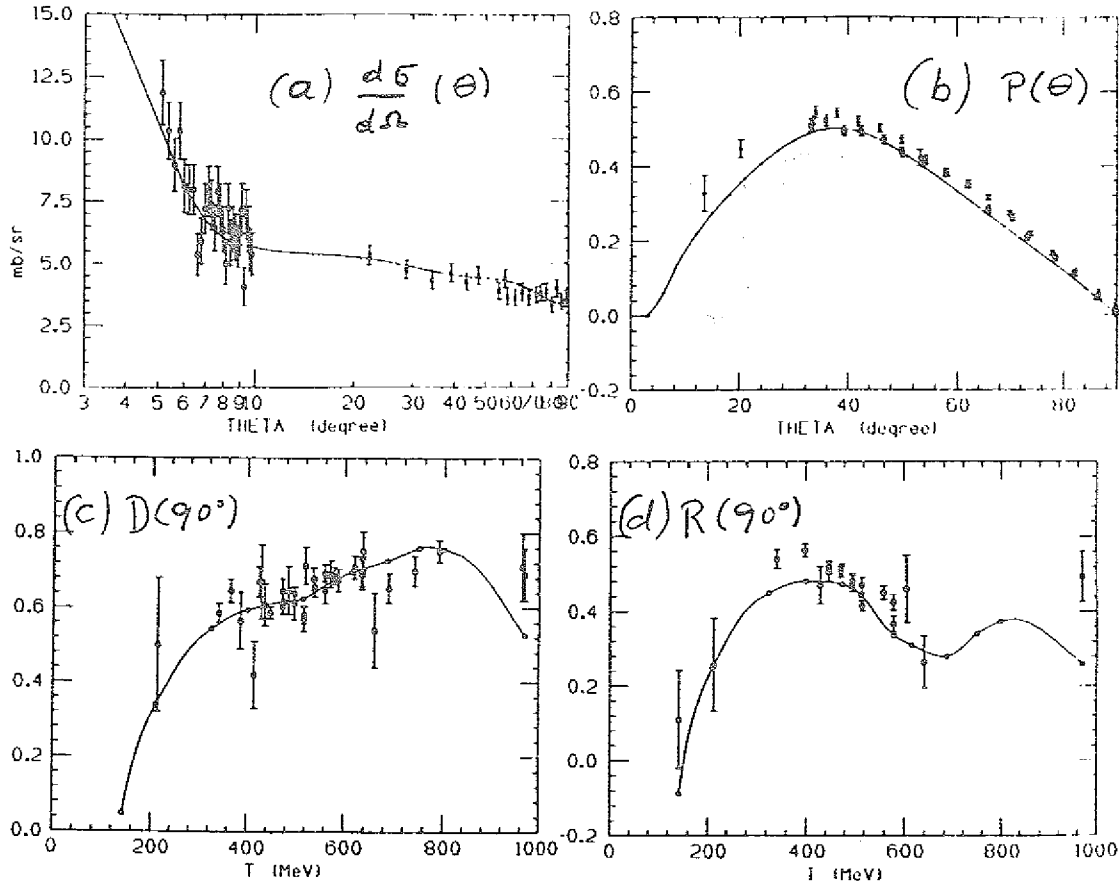


Fig. 2. Elastic pp cross sections. The predictions of the improved model compared with elastic pp data; (a) $d\sigma(\Theta)/d\Omega$ and (b) $P(\Theta)$ at $T_L = 515$ MeV; (c) $D(90^\circ)$ and (d) $R(90^\circ)$ excitation functions.

In addition to including the diagonal $N\Delta$, OPE, potentials in all channels (and correcting some coefficient errors in the previous work) several partial waves have been improved by the addition of higher orbital angular momentum or higher mass channels to which they are specially sensitive. In particular the $NN(^1S_0)$ system, which had previously been coupled to the $N\Delta(^5D_0)$ and the $\Delta\Delta(^1S_0)$, has now also been coupled to the $\Delta\Delta(^5D_0)$ channel.

In addition, the OPE potentials coupling the $N\Delta$ channel have been included. These additions were important in decreasing the two-pion $\Delta\Delta$ production to the experimental value [4], while maintaining the fit to δ and η . The $L = 2$, $\Delta\Delta$ channel is less inelastic than the $L = 0$, and the transition potentials convert some of the $\Delta\Delta$ production to $N\Delta$.

Similarly in the $NN(^3P_1)$ system the $N\Delta(^5F_1)$ and $\Delta\Delta(^3P_1)$ channels were added to the previous $N\Delta(^5P_1)$ components. This reduced the inelasticity for $T_L < 1$ GeV as required, without reducing the attraction needed for δ . In the $NN(^3P_0)$ system the $\Delta\Delta(^3P_0)$ was added to the $N\Delta(^3P_0)$ for the same reason.

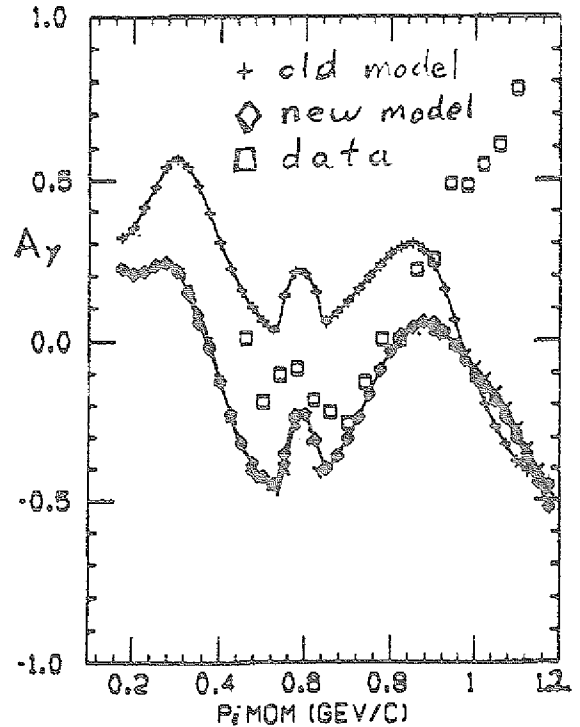


Fig. 3. $A_Y(pp \rightarrow pn\pi^+)$. Prediction of the improved model compared with $T_L = 800$ MeV production data.

The result of these improvements in the pp model is very satisfactory. Figure 1 shows the fit to $\Delta\sigma_L$ and $\Delta\sigma_T$. Figure 2 illustrates the good fit obtained to all pp elastic scattering data by showing four typical examples. Other $d\sigma/d\Omega$, P_N , R' , R , A , A' , D_T , D , C_{LL} , C_{NN} , A_{zx} and A_{xx} data $0 \leq T_L \leq 800$ MeV is similarly well-fitted. Figure 3 shows the improvement in fitting the A_y distribution of single pion production in pp collisions⁴⁾. The accuracy and completeness of the present model permits a new, more reliable extension to the higher energies ($T_L = 1.8 - 3$ GeV) at which the effects of exotic six quark resonances are manifested¹⁾.

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The $\Lambda N - \Sigma N$ Interaction with Isobar Coupling

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The confinement property of QCD leads one to expect that the long-range interaction between hadrons is mediated by hadron exchange. Experience shows that this is certainly so beyond ~ 1.5 fm inter-hadron separation, and lattice QCD calculations indicate such behavior. At short enough range the asymptotic freedom property of QCD implies that simple quark configurations will dominate over hadronic clusters of quarks, so that the hadron exchange interaction representation will be inadequate or at least much more complicated. A conventional, minimal allowance for the breakdown of the hadronic description is the use of form factors with "cut-off" parameters for each vertex. This approach, using meson exchange potentials in momentum space, has been used for several two-particle systems by the KFA-Jülich group, including a successful fit to the $\Lambda N - \Sigma N$ system [1]. In our view this is not the best description of the transition from the hadronic to the quark-gluon phase, and in particular does not allow for the effect of levels of the six-quark configurations.

Hybrid approaches in which hadronic degrees of freedom are replaced by quark degrees of freedom at small distances have been used with some success for the nucleon-nucleon system. In particular, the R-matrix method has been successful in producing quantitative fits to a broad range of two-nucleon data [2,3] and has also predicted six-quark "exotics" [4-6] for which there is some experimental evidence [5-7]. R-matrix theory provides powerful relations between states of the inner region and the S-matrix of the outer region, yielding quantitative results with relatively simple calculations. This allows the inclusion of the most important physics ingredients. R-matrix theory is formulated in configuration space, which is also appropriate to the expected rapid transition as a function of distance between the two limits of QCD [8].

This method is extended here to the $\Lambda N - \Sigma N$ system. We also recognize, from the experience with the NN system, that coupling to channels with isobar components has an important influence well below their thresholds. Furthermore, the exotic excitations are at higher energies than the lowest isobar thresholds, so that the energy dependence due to the threshold branch points must be treated accurately. The only way to do so is with an explicit coupled channel calculations, as we do here in coupling to channels with the $\Delta(1232)$ or the $\Sigma(1385)$.

In addition to the exchange of π, σ (for 2π correlated),

$\eta, \Lambda, \rho, \omega, K$ and K^* mesons used in Ref. [1], we have included the exchange of η' and ϕ mesons. The transition amplitudes between each two-baryon pair were expanded to bilinear order in the three-momentum transfer and the average of initial and final three-momenta. The Fourier transforms of these produced potentials with central, spin-spin, tensor, spin-orbit, $i(\vec{S}_1 - \vec{S}_2) \cdot \vec{L}$ and $(\vec{S}_1 \times \vec{S}_2) \cdot \vec{L}$ components. The last two of these operators connect singlet and triplet spin states, their coefficients vanishing when the baryons of leg 1 have the same mass as the baryons of leg 2. Partial wave amplitudes were calculated for all the $J = 0, 1$ and 2 channels.

The separation radius, r_0 , was fixed at 1.1 fm, in accordance with the conditions imposed by the perturbative QCD Hamiltonian and the R-matrix method [4,7]. We varied only the energy-independent boundary conditions f_{ij}^0 in

$$r_0 \psi_1(r_0) = \sum_i f_{ij} \psi_j(r_0)$$

$$f_{ij} = f_{if}^0 + \sum_k \frac{\rho_{ij}^k}{\omega - \omega_k}$$

where ω is the barycetic mass and the ω_k and ρ_{ij}^k are determined by the internal quark configurations. The resulting fit to all the available data is very good. The Λp scattering length are predicted to be $\alpha_s = -2.53$ fm and $\alpha_t = -2.03$ fm, within the wide experimental limits¹⁾. The inelastic $\Sigma^- p$ capture ratio to $\Sigma^0 n$ compared to $\Sigma^0 n$ and Λn combined) is predicted to be 0.483 in agreement with the experimental value 0.474 ± 0.016 ¹⁾. Figure 1 compares total and differential cross section predictions with the available data (the $\Sigma^- p \rightarrow \Sigma^0 n$ cross section, not shown, also fits well). Only the fit to the Λp elastic cross-reaction in the vicinity of the ΣN threshold ($P_\Lambda > 400$ MeV/c) is poor in magnitude, but it is consistent with the threshold peak seen in the data. Figure 2 shows the large isobar threshold effect on the inelasticity η for a partial wave coupling to an S-wave isobar channel. The effect decreases with the angular momentum of the isobar channel. The effect of the coupling on the real phases, δ , is to add increasing attraction up to inelastic threshold, but the discontinuity at threshold is small.

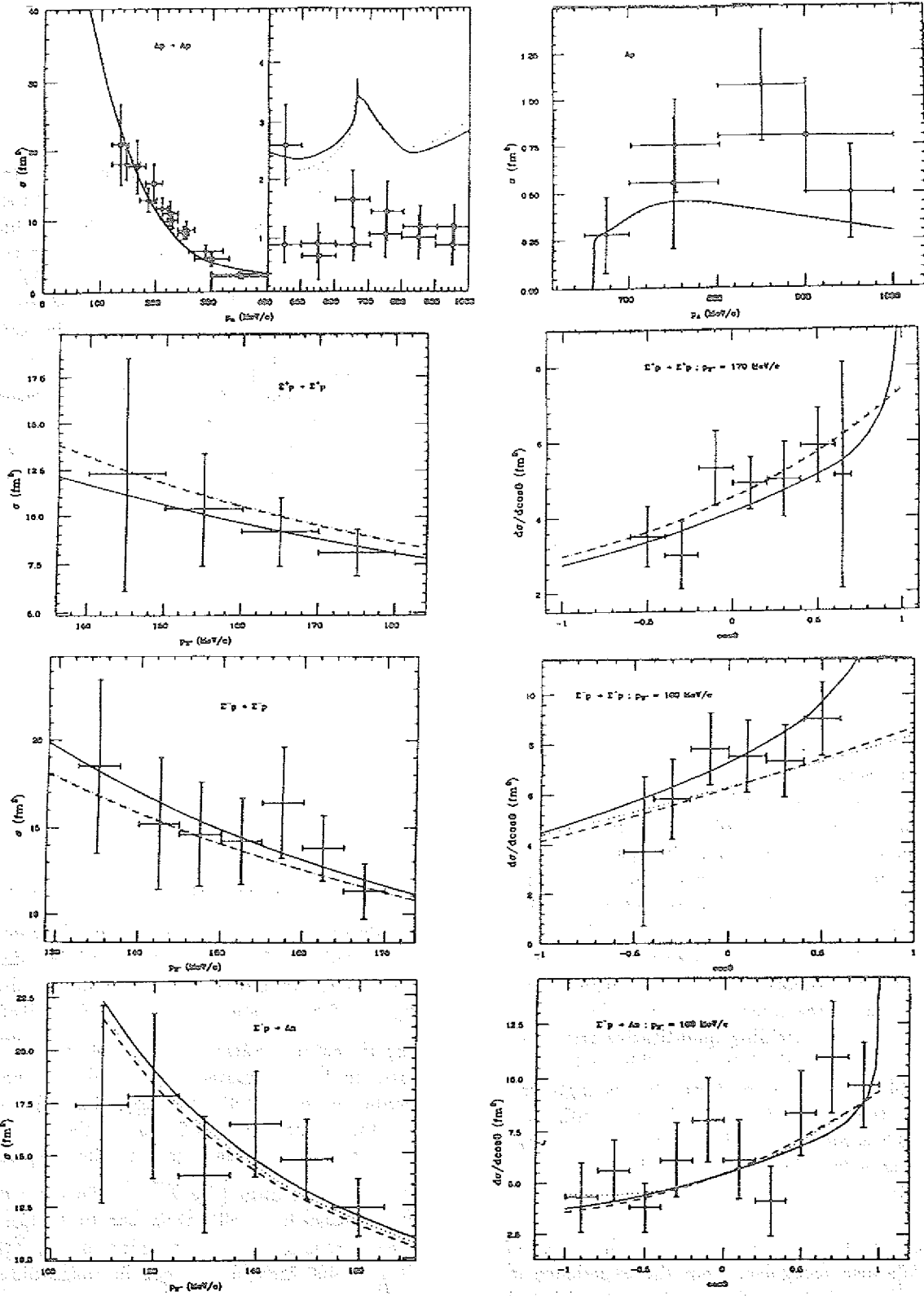


Fig. 1. Comparison of model predictions with available total and differential cross section data. The full curves represent the model with isobar coupling and includes the Coulomb interaction where relevant. The dashed curves omit the Coulomb interaction. The dotted curves represent a model which omits isobar coupling and the Coulomb interactions.

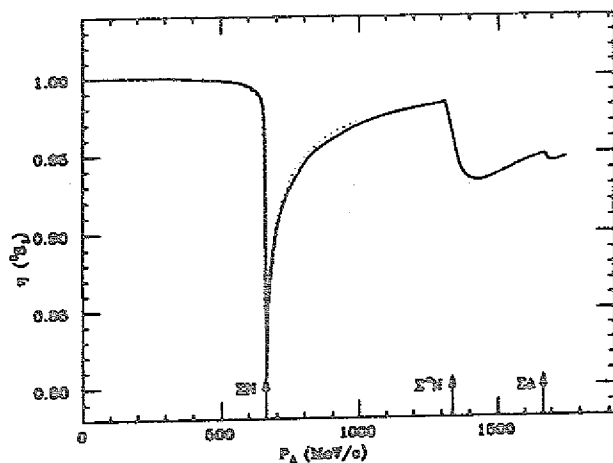


Fig. 2. The inelasticity parameter $\eta(^3S_1)$ in Λp elastic scattering. The effects of the thresholds for the Σ^*N and $\Sigma\Delta$ channels are evident, as well as the stronger effect of the ΣN threshold.

The position and residues of the f -matrix poles due to the six-quark configurations have been presented earlier [7]. Exotic inelastic resonances of width > 50 MeV are predicted to begin just above the $\Sigma\Delta$ threshold. Detailed predictions for the behavior of observables at these resonances, as well as at the isobar production thresholds, await the determination of f -matrix parameters for the higher angular momentum partial waves that are significant at these energies.

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The structure of hadrons in a dense, hot nuclear environment is of considerable current interest. In several respects the role of mesons has come into focus: 1. it is believed that baryons consist of a small quark core surrounded by a meson cloud. While the core is more or less inert until the critical temperature and density for quark-gluon plasma formation is approached, cloud polarization due to the nuclear medium may play a significant role under more moderate conditions, 2. the nucleon-nucleon interaction at large and intermediate distances is mediated by one-pion and correlated two-pion exchange and modifications of the pion propagator may lead to new perspectives for the nuclear equation of state [2], 3. the rho-meson is coupled strongly to the two-pion continuum and will hence be modified in the medium. It may thus serve as a monitor of density and temperature in relativistic heavy-ion collisions through measurement of e^+e^- -yields [3].

In a first study we have examined the role of two-pion correlations in cold nuclear matter by considering in-medium $\pi\pi$ -scattering using a meson-exchange model for the free T-matrix [5].

This model treats $\pi\pi$ - and $K\bar{K}$ scattering as coupled channels and accounts well for the scattering data up to 1.4 GeV center-of-mass energy. The T-matrix is given by

$$T^{ij}(q, q', E) = V^{ij}(q, q', E) + \sum_i \int \frac{d^3k}{(2\pi)^3} V^{ii'}(q, k, E) D_{i'}(k, E) T^{i'j}(k, q', E). \quad (1)$$

where $D_i = D_{\pi\pi(K\bar{K})}$ denotes the free $\pi\pi(K\bar{K})$ propagator which is chosen to be either of a time-ordered (TO) form [5] or a Blankenbecker-Sugar (BS) form [9]. The quasipotentials V^{ij} are constructed from s- and t-channel meson exchange using SU(3)-symmetry relations among the coupling constants and phenomenological form factors at the vertices as described in ref. [5]. Fits to phase shifts and inelasticities are very satisfactory and it is found that correlated $\pi\pi$ and $K\bar{K}$ pairs play a crucial role in explaining the data, particularly below 1 GeV. In this region the cross section is completely saturated by s- and p-waves.

Due to modifications of the pion propagator, in-medium $\pi\pi$ -scattering in the low partial waves is strongly affected especially near threshold. Changes in the kaon propagator as well as in the s- and t-channel mesons are less significant. In matter, the pion acquires a density-dependent self energy Π , such that the propagator is given by

$$D_{\pi}(k, \omega; \rho) = (\omega^2 - k^2 - m_{\pi}^2 - \Pi(k, \omega; \rho))^{-1}. \quad (2)$$

The most important contributions are attributed to p-wave pion-nucleon scattering exciting virtual particle-hole and isobar-hole states. Since these processes provide the strongest dependence on k and ω we treat them explicitly. All other processes can be incorporated as a renormalization factor z_{π} which multiplies the πNN - and $\pi N\Delta$ vertices and can be reliably estimated in a gas approximation [1]. In evaluating the particle-hole and isobar-hole contributions to Π is essential to incorporate short-range correlations which we parameterize by the Migdal parameters [6] g'_{NN} , $g'_{N\Delta}$ and $g'_{\Delta\Delta}$. Their values are determined from Gamow-Teller resonance systematics, hadron-nucleus excitations of the isobar and double-charge exchange reactions [7]. Then Π is given by

$$\Pi = k^2 z_{\pi}^2 \frac{(\chi_N^0 + \chi_{\Delta}^0 + \chi_N^0 \chi_{\Delta}^0 (2g'_{N\Delta} - g'_{NN} - g'_{\Delta\Delta}))}{((1 - g'_{NN} \chi_N^0)(1 - g'_{\Delta\Delta} \chi_{\Delta}^0) - g'^2_{N\Delta} \chi_N^0 \chi_{\Delta}^0)} \quad (3)$$

where $\chi_N^0(\chi_{\Delta}^0)$ are the usual nucleon (isobar) Lindhard functions [8]. To simplify the $\pi\pi$ scattering calculations we assume that the nucleons are static and neglect the isobar width. Thus, broadening effects due to Fermi motion, two-particle two-hole correlations and delta decay are not taken into account. Then the Lindhard functions have a very simple (k, ω) -dependence with

$$\chi_N^0 = \frac{\alpha_N}{\omega^2 - (\omega_k^N)^2}; \quad \alpha_N = 2\omega_k^N \rho \left(\frac{f_{\pi} \Gamma_{\pi}}{m_{\pi}} \right)^2$$

$$\chi_{\Delta}^0 = \frac{\alpha_{\Delta}}{\omega^2 - (\omega_k^{\Delta})^2}; \quad \alpha_{\Delta} = \frac{32}{9} \omega_k^{\Delta} \rho \left(\frac{f_{\pi} \Gamma_{\pi}}{m_{\pi}} \right)^2 \quad (4)$$

where $\omega_k^N = \sqrt{M_N^2 + k^2}$ and $\omega_k^\Delta = \sqrt{M_\Delta^2 + k^2}$ and Γ_π denotes the πNN -form factor. Given these expressions D_π can be constructed from the three-branch eigenvalue equation

$$\mathcal{H}^2 \begin{pmatrix} \Phi_N \\ \Phi_\pi \\ \Phi_\Delta \end{pmatrix} = \Omega^2 \begin{pmatrix} \Phi_N \\ \Phi_\pi \\ \Phi_\Delta \end{pmatrix} \quad (5)$$

where

$$\mathcal{H}^2 = \begin{pmatrix} (\omega_k^N)^2 + g_{NN}^2 \alpha_N & z_\pi \sqrt{\alpha_N} k & g_{N\Delta}^2 \alpha_N \\ z_\pi \sqrt{\alpha_N} k & k^2 + m_\pi^2 & z_\pi \sqrt{\alpha_\Delta} k \\ g_{N\Delta}^2 \alpha_\Delta & z_\pi \sqrt{\alpha_\Delta} k & (\omega_k^\Delta)^2 + g_{\Delta\Delta}^2 \alpha_\Delta \end{pmatrix} \quad (6)$$

as

$$D_\pi(k, \omega; \rho) = \sum_i |\Phi_i(k, \rho)|^2 / [\omega^2 - \Omega_i^2(k; \rho) + i\eta] \quad (7)$$

Due to mixing with nucleon- and delta-hole excitations the pion is no longer an eigenstate and the probability $|\phi_\pi(k; \rho)|^2$ to find a pion of momentum k at density ρ is a strong function of density. Up to ρ the pion mixes dominantly with the delta-hole branch which exchange their structure around $k \sim 2.5 m_\pi$. The lower branch then becomes the "pisobar" and is experimentally accessible through the spin-longitudinal reponse function in the isobar region. As the density increases the mixing of the pion branch with the nucleon-hole branch becomes very strong with $\phi_N(k, \rho)$ approaching unity as ρ approaches ρ_c , the critical density for pion condensation.

It is now straightforward to determine the $\pi\pi$ -scattering observables in matter. As long as we neglect the medium dependence of the quasipotentials V_{ij} and of the kaon propagator we just have to replace the vacuum $\pi\pi$ propagator by the dressed one. For two pions with relative momentum k and CM energy E the dressed propagator is given by

$$D_{\pi\pi}(k, E; \rho) = \sum_i \frac{\Phi_i^2(E + 2\Omega_i)}{E^2 - 4\Omega_i^2} + \sum_{i < j} \frac{2\Phi_i \Phi_j (E + \Omega_i + \Omega_j)}{E^2 - (\Omega_i + \Omega_j)^2} \quad (8)$$

in the TO-form, or

$$D_{\pi\pi}(k, E; \rho) = \sum_i \frac{\Phi_i^2 4\Omega_i}{E^2 - 4\Omega_i^2} + \sum_{i < j} \frac{4\Phi_i \Phi_j (\Omega_i + \Omega_j)}{E^2 - (\Omega_i + \Omega_j)^2} \quad (9)$$

in the BS-form. For both choices the scattering results are similar. In Fig. 1 we show results for the imaginary part of the s-wave and p-wave T-matrix at several densities. The density dependence in the s-wave is large especially near

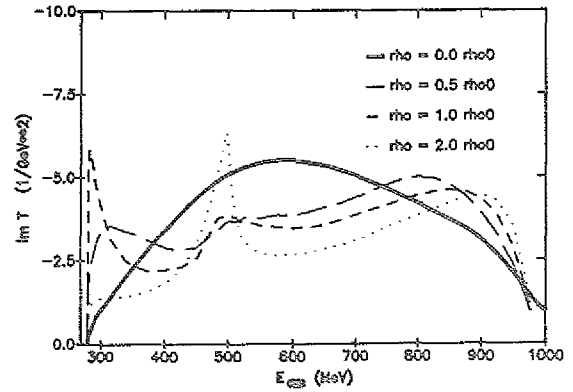


Fig. 1. Imaginary part of the T-Matrix in the $J=0; I=0$ channel for some densities of nuclear matter: $\rho=0$ (solid), $0.5 \rho_0$ (long-dashed), ρ_0 (short-dashed) and $2 \rho_0$ (dotted).

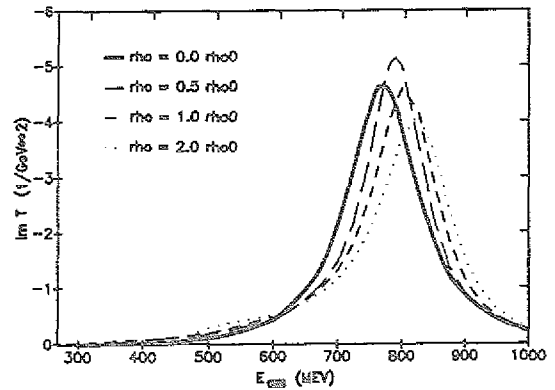


Fig. 2. Same as Fig. 1, for the $J=1; I=1$ channel.

threshold. This behavior is explained by the in-medium reduction of the pion group velocity $d\omega/dk$ for small momenta which enlarges the phase space for scattering. Bound states occur for $d\omega/dk=0$ and are found above ρ_0 . Including damping of the pion propagator these will acquire a finite width, however. The modification of the p-wave (the ρ -meson) is not quite as dramatic. As density increases the resonance position is shifted to slightly higher energies. These results should prove interesting in relation to the medium modification of mesons and baryons under extreme conditions.

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Recently [1] we have presented a microscopic model for the nucleon-antinucleon ($N\bar{N}$) interaction, which is based on meson-baryon dynamics. The elastic part is the G-parity transform of the (full) Bonn NN potential [2]. Annihilation is accounted for by microscopic baryon ($N, \Delta, \Lambda, \Sigma, Y^*$) exchange processes with two-meson intermediate states involving $\pi, \rho, \omega, \sigma, \delta, K$ and K^* . Open form factor parameters in those diagrams have been adjusted to the $N\bar{N}$ cross section, and predictions have been made for differential cross sections and spin observables. Whereas the resulting total and elastic cross sections are in good agreement with empirical data, the corresponding charge-exchange cross section turns out to be too large; consequently, the annihilation cross section is about 10% too low. Also, sizable discrepancies occur in differential cross section and spin observables.

The reason for the unsatisfactory description lies in the fact that the annihilation processes so far included contribute only 30% to the empirical annihilation. Thus, in order to obtain sufficient annihilation the parameters have been artificially enhanced. Obviously, such a procedure then leads to an unrealistic spin- and isospin dependence of the annihilation interaction.

In the present work, we improve our calculations in the following way: First, we include further two-meson annihilation channels involving also η, a_1, a_2 and f_2 . Second, we now take all two-meson annihilation channels with their realistic size, i.e. adjust the open parameters independently to the empirical known branching ratios of those channels in the annihilation of the $N\bar{N}$ system [3].

The considered two-meson annihilation channels now amount to about 50% of the total observed annihilation. In order to take into account the other half, consisting of further two-meson channels but also of channels involving three or more uncorrelated mesons (which are very difficult to describe in the meson-exchange framework) we now add a simple optical potential of Gaussian form, i.e.

$$V_{\text{opt}} = (U_0 + iW_0) e^{-\frac{r^2}{2r_0^2}}$$

with $U_0 = -0.5$ GeV, $W_0 = -0.5$ GeV, $r_0 = 0.4$ fm

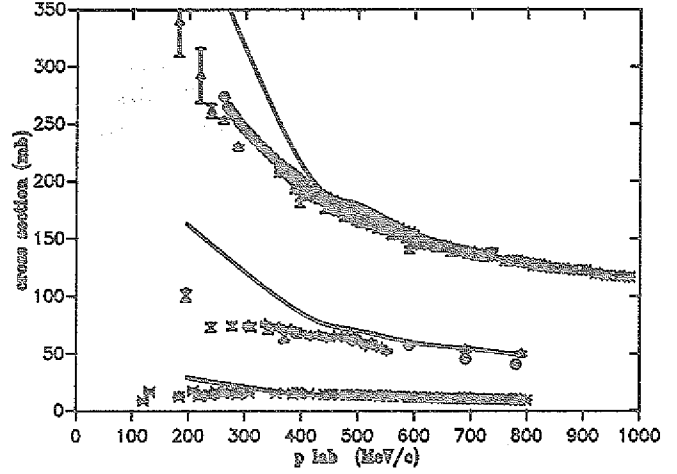


Fig. 1. Total, integrated elastic and integrated charge-exchange cross section, as obtained with our new model. The references to the data can be found in [1].

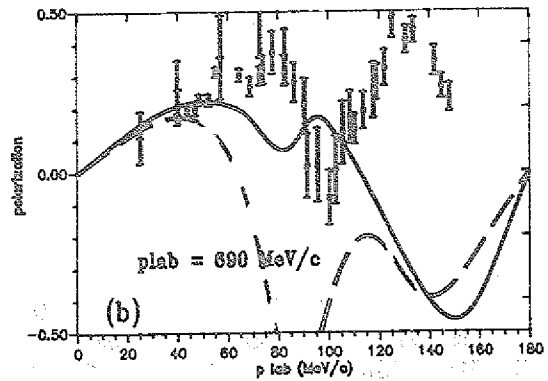
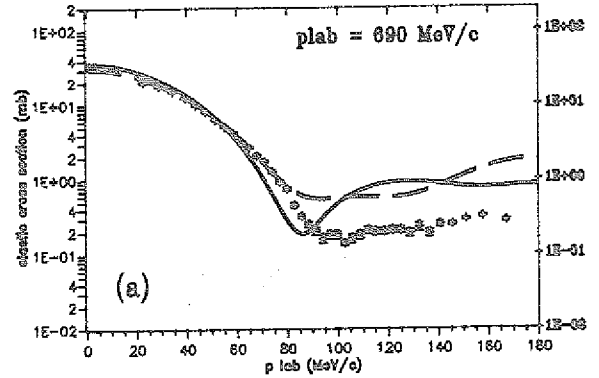


Fig. 2. Elastic cross section and polarization at $p_{\text{lab}} = 690$ MeV/c. Results obtained in our new study (solid) in comparison with the microscopic model presented in [1] (dashed). The references to the data can be found in [1].

adjusted to the integrated NN cross sections.

Results are shown in fig. 1. A good description of total and elastic cross sections has been achieved; concerning the charge-exchange cross section, the prediction is much improved compared to our former result of [1].

Fig. 2 shows the result for the differential cross section and polarization (at one specific lab momentum) in the elastic channel. The description is now of good quality. In the charge-exchange channel (fig. 3), the description is far improved in comparison with our former study. This reaction is very sensitive to details of the model so that minor changes will bring these predictions fully in line with experiment.

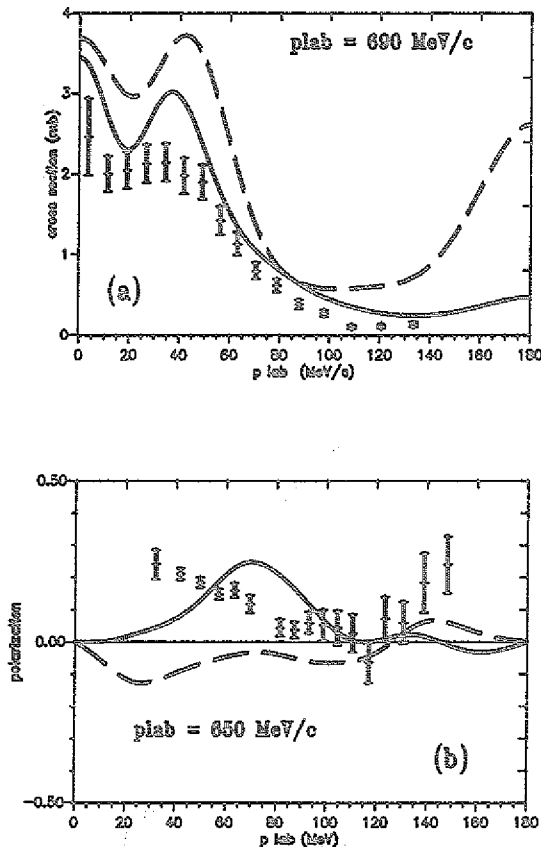


Fig. 3. Same as fig. 2, for the charge-exchange channel.

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$N\bar{N} \rightarrow \pi\pi$ scattering amplitude in the pseudophysical region in the TOPT framework

H. Kim, J. Durso*, K. Holinde, and J. Speth

The correlated 2π -exchange has a dominant contribution to the physical NN interaction in higher angular momentum partial waves which correspond to the long and intermediate range of the NN-interaction. In the full Bonn potential, the fictitious " σ " meson exchange was used as the correlated $\pi\pi$ S-wave contribution which provides about 2/3 contribution of the total 2π -exchange [1]. Although this fictitious " σ " meson can help us shape to the $\pi\pi$ S-wave interaction in NN potential, it is nothing but a parametrisation of the physical $\pi\pi$ S-wave exchange because there is no resonance in S-wave in $\pi\pi$ interaction. We need to make this S-wave correlated $\pi\pi$ rescattering clear.

As a first step to construct the physical correlated 2π -exchange NN potential, we have studied the pseudophysical $N\bar{N} \rightarrow \pi\pi$ scattering amplitude below the nucleon threshold. The square of this amplitude in each s-, p-wave gives the correlated 2π -exchange NN-potential via the dispersion integral with the full relativistic unitarity relation. The $N\bar{N} \rightarrow \pi\pi$ amplitude is just connected to $\pi\pi$ amplitude by crossing symmetry. We have calculated the $N\bar{N} \rightarrow \pi\pi$ amplitude in the time ordered perturbation theory in the pseudophysical region $4\mu_\pi^2 < t < 50\mu_\pi^2$. We can include $\pi\pi$ -rescattering effect in a unitary way in TOPT.

$$T(N\bar{N} \rightarrow \pi\pi) =$$

$$V_{\text{Born}}(N\bar{N} \rightarrow \pi\pi) \left[1 + \frac{1}{Z - 2\omega_k} \tau(\pi\pi \rightarrow \pi\pi) \right] \quad (1)$$

V_{Born} is explicitly real and contains all informations concerning $N\bar{N} \rightarrow \pi\pi$. Since the Born terms are governed by N, Δ exchange, we have included only N, Δ . The Δ -exchange makes a significant contribution to $J = 0$ $N\bar{N} \rightarrow \pi\pi$ at relatively small t values. But this Δ -exchange spoils $J = 1$ $N\bar{N} \rightarrow \pi\pi$ amplitude corresponding ρ meson resonance. In order to avoid this problem, we should introduce the offshell correction in the Lagrangian for $N\Delta\pi$ vertex and $N\bar{N} \rightarrow \pi\pi$ amplitude or simply neglect Δ -exchange as J. Durso, et al. did [2]. In the case that we introduce the off-shell term in the Lagrangian, we need large cut-off masses.

The factor in the bracket in Eq.(1) can be understood as the wave operator which creates the physical $\pi\pi$ wave function when it acts on the $\pi\pi$ plane wave function. We have used the $\pi\pi$ interaction calculated by D. Lohse, et al. [3], in which the $\pi\pi$ phase shifts were computed in the coupled formalism with $K\bar{K}$ channel and were compared to the experimental data with good agreement.

In order to compare our $N\bar{N} \rightarrow \pi\pi$ amplitude to the empirical data by H  hler, et al. [4], we connect our $T(N\bar{N} \rightarrow \pi\pi)$ to the scattering amplitude defined by Frazer and Fulco [5].

$$f(t)_+^J = \frac{-pE_p}{(pq)^J} \pi \sqrt{E_p \omega_q} T(t)_+^J \quad (2)$$

$$f(t)_-^J = \frac{-p}{(pq)^J} \pi \sqrt{E_p \omega_q} T(t)_-^J \quad (3)$$

Fig. 1 compares our Ref_+^0 with the empirical Ref_+^0 . In the small t region ($4\mu_\pi^2 - 20\mu_\pi^2$) our Ref_+^0 agrees well with the empirical data. From Fig. 2 finds that $\text{Im}f_+^0$ is somehow different from the empirical data above $15\mu_\pi^2$. However, since the

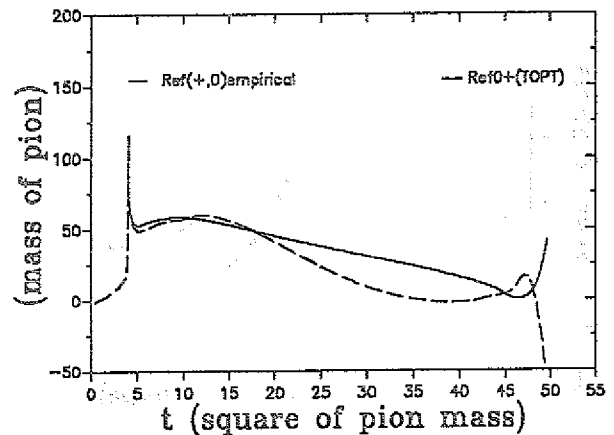


Fig. 1. The solid line is the real part of the amplitude in s-wave. It compares with the empirical data (dashed line) by H  hler.

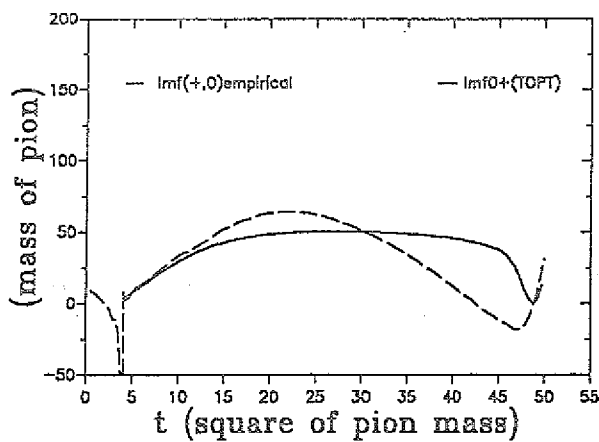


Fig. 2. The solid line is the imaginary part of the amplitude in s-wave. It compares with the empirical data (dashed line) by Höhler.

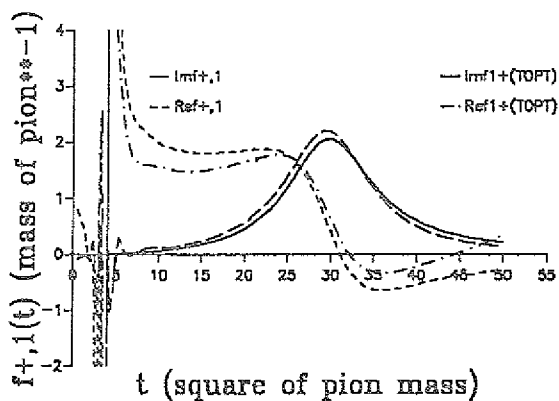


Fig. 3a. These figures are the p-wave helicity plus amplitude. The dot-dashed line is the real part of the amplitude. The solid line is the imaginary part of the amplitude.

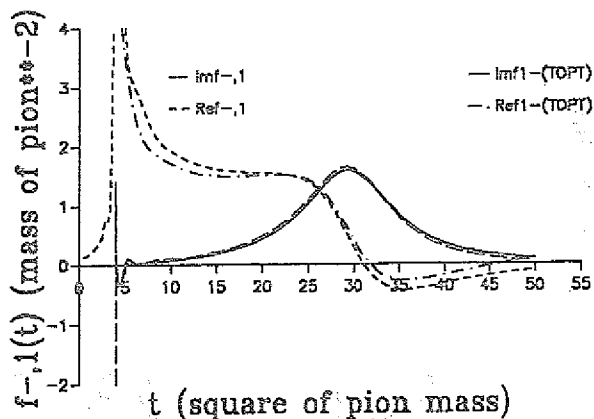


Fig. 3b. These figures are the p-wave helicity minus amplitude. The dot-dashed line is the real part of the amplitude. The solid line is the imaginary part of the amplitude.

modulus of the scattering amplitude is much more important to determine the $\pi\pi$ correlation in the NN-potential, that kind of small differences does not make any serious difficulty. The comparison of the p-wave is given in Fig. 3.

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Model calculations will be confronted with the accurate new data from LEAR on $d\sigma/d\Omega$ $\bar{p}p$ elastic and charge exchange $\bar{p}p \rightarrow \bar{n}n$ scattering at $p_{lab} = 700$ MeV/c as well as the analyzing power of these two reactions. Some parameters of the models will be changed, to show how the observables change with variations of these parameters. The hope is to extract information on the spin-dependent meson exchange potential (MEP) and use this new information together with the nucleon-nucleon scattering data to further restrict the large distance parts of MEP, because the charge-exchange reaction, given by the difference between two large isospin amplitudes, is a very sensitive probe.

From the static baryon properties (magnetic moments, rms radii, masses etc.) and the success of chiral symmetry in low energy hadronic physics, a (primitive) model of the nucleon's structure has emerged: The nucleon has a core of confined quarks (core rms radius of 0.5–0.6 fm). This core is surrounded by an extended meson cloud which is distributed according to the requirements of chiral symmetry (proton rms radius 0.9 fm). A consequence of this model is that NN scattering for large impact parameters should be given by MEP. In fact, for large NN partial waves ($L \geq 2$) and $T_{lab} \leq 300$ MeV, the meson exchange potential successfully describes the NN phase shifts without any vertex form factors and with the expected SU(3) coupling constant $g_{\omega}^2/4\pi \approx 4.5$ [1].

The MEP applied to the NN S- and P-waves requires vertex form factors and/or short distance cut-offs. However, quark model calculations successfully reproduce the energy behaviour of the NN S-wave phase-shifts [2]. For NN S-wave scattering the two nucleons feel the meson exchange forces at large separations but the antisymmetry of the overlapping six quark wave functions is crucial in order to get the observed hard-core-like S-wave phase shifts for increasing NN energy. This short distance NN repulsion includes the standard SU(3) value for g_{ω} . However, in most MEP (including OBEP) models a value $g_{\omega}^2/4\pi = 10$ –12 is required if the NN repulsion is ascribed to ω -exchange alone (which includes necessary form factors and/or short distance parametrizations) [3,4,5].

In light of the above model presentation, we think

it is justified to reduce the g_{ω} value from the usual OBEP value to its SU(3) value. The quark degrees of freedom are not expected to give any sizeable repulsion or attraction in $\bar{N}N$. Therefore what will be presented below is an exploratory calculation to demonstrate the (in)sensitivity of observables to a smaller value of g_{ω} . A more careful investigation is being pursued [6].

In the following calculations we will use an annihilation potential with a radial dependence derived from confined quark model wave functions where all intermediate quark states have been summed over. This annihilation range is basically determined by the overlap of the two extended quark cores. The strength of this interaction is fitted to the total cross section [7]. We will also use another simplistic model for annihilation in this calculation, where it is assumed that the annihilation can be simulated by a "black sphere" like boundary condition where we only have incoming flux and no reflections. This boundary is of the order of 0.5 fm in order to fit the total $\bar{N}N$ cross section [8]. In both cases, due to the strong attraction of the MEP, the $\bar{N}N$ wave function is pulled into the annihilation region, but we still expect as in NN that the MEP description is valid for $\bar{N}N$ partial waves $L > 2$ at low energies.

The figure shows first exploratory calculations using the original value for $g_{\omega}^2/4\pi$ and a value which is half the value of the two different OBEP models used (the Bryan-Phillips [9] and the Nijmegen model D [4]) connected by a shaded region in the figure. No large effects are seen in the forward $d\sigma/d\omega$ for $\bar{p}p$ elastic and charge exchange scattering with a reduced value of g_{ω} . The backward scattering cross sections are very sensitive to form factors or short distance models (including the annihilation model). As seen the rise in the analyzing powers of the two scattering processes at small angles are relatively insensitive to the change in g_{ω} .

However, $A_0(\bar{p}p \rightarrow \bar{n}n)$ changes very much with changes in g_{ω} including the value of the first forward maximum, and larger angle scattering.

Some MEP model parameters which influence NN

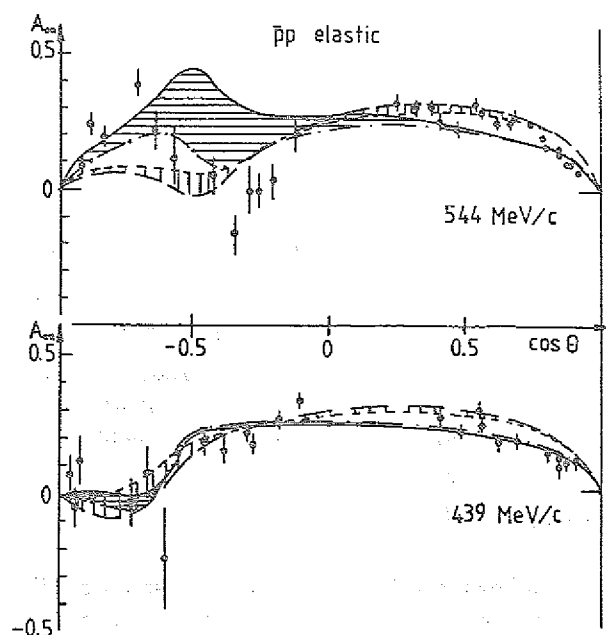


Fig. 1 The analyzing power for elastic $\bar{p}p$ scattering at $p_{lab} = 544$ and 439 MeV/c. Experimental points from [10].

and $\bar{N}N$ should be carefully reexamined. The questions raised goes to the core of our theoretical understanding of the MEP description for NN and $\bar{N}N$, and especially the $V_{2\pi}$. Furthermore, this author has argued that the energy behaviour of the NN S-wave phase shifts are one example where quark degrees of freedom manifest itself in low energy physics. This implies that the NN repulsion is not described by ω exchange alone and as a consequence the $\bar{N}N$ potential is much less attractive than what OBEP calculations have assumed up to now. A refit to the NN data including the new precise $\bar{N}N$ data at low $\vec{q}$ is necessary. The meson exchange description cannot be valid for large $\vec{q}$, where not only MEP vertex form factors (size of the hadrons) but also annihilation becomes very important. Fortunately, we only need to describe the global properties of annihilation in $\bar{N}N$ scattering processes, but our models for annihilation are still too primitive to describe $d\sigma/d\Omega$ and A_{00} for high $\vec{q}$ and $p_{lab} < 1$ GeV/c.

As stated the charge exchange reaction is a unique testing ground for the MEP models since this reaction is given by the differences of two large isospin amplitudes. Experimentally it is still not clear if for the very forward $d\sigma/d\Omega$ ($\bar{p}p \rightarrow \bar{n}n$) we have a shoulder or a minimum. Some models give a shoulder, others a minimum. As discussed by Phillips [9] this behavior of the cross section is tied

to the long range spin-structure of the MEP. What is correct? Measurements of the very forward $\bar{p}p \rightarrow \bar{n}n$ cross section is necessary to settle this question.

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The virtual pion cloud plays a crucial role in the explanation of the nucleon's long-range structure. It was recognized already in the 1950's, that the process $n \rightarrow \pi^- p$ naturally explains the negative tail of the neutron's charge distribution. The concept of the pionic cloud may be easily extended by taking into account effects of additional mesons. In deep-inelastic lepton-nucleon scattering, the virtual off-shell photon can scatter from either a meson ($\pi, \rho, \omega, \sigma, K, K^*, \eta$) or the recoil baryon (nucleon, delta, hyperons). The contribution of quarks from a (pseudo)scalar meson can be written as a convolution of the meson structure function and its momentum distribution $f_M^i(y)$ in the nucleon [1].

$$xq_f^{\text{Sull},i}(x, Q^2) = Z \int_x^1 dy f_M^i(y) x/y q_{f,M}^i(x/y, Q^2), \quad (1)$$

where $q_{f,M}^i(x/y, Q^2)$ is the quark distribution function in the meson M_i . A similar equation holds for the recoil baryons. The only difference in comparison to ref. [1] is the presence of the factor Z in front of the right hand side of Eq.(1) which takes into account the proper renormalization of the valence quarks. The functions $f_M^i(y)$ are momentum distributions of the meson M_i in the nucleon. They are functions of the coupling constants and form factors. Explicit expressions have already been presented elsewhere [2]. However, the way we treat vector mesons is quite different. Now we will sketch briefly the essence of the present approach. For this purpose let us consider the reaction $N(e, e') \tilde{X} B$, where the virtual γ^* interacts with the vector meson M_V and turns it into debris $\tilde{X}$, leaving the baryon B as a spectator of the reaction. The square of the invariant amplitude for the total process, averaged over spins in the initial channel and summed over the spins of particles in the final channel, may be written down in the schematic way

$$|M|^2 = L_{\mu\mu'} H_{\alpha\alpha'} W^{\mu\alpha\mu'\alpha'} \quad (2)$$

In this formula the tensor $L_{\mu\mu'}$ is connected with

the electromagnetic vertex $e \rightarrow e' + \gamma^*$, the tensor $H_{\alpha\alpha'}$ with the low energy vertex $N \rightarrow B + M_V$ and $W^{\mu\alpha\mu'\alpha'}$ with the $\gamma^* + M_V \rightarrow \tilde{X}$ process.

Since in general the deep-inelastic structure functions for vector mesons are unknown experimentally, we are forced to replace the unpolarized deep-inelastic structure functions of the ρ meson by those for the pion. In the naive quark structure model, the quark content of the ρ meson is the same as that for the pion. In more sophisticated models the ρ meson may be treated as a mixture of the quark-antiquark pair and the two-pion component [3]. The structure function for K^* and other vector mesons may then be easily obtained by the assumption of SU(3) symmetry. Since the $H_{\alpha\alpha'}$ tensor can be treated as a generalized unnormalized density-matrix of virtual vector particles [4], the optimal Lorentz-invariant replacement is

$$H_{\alpha\alpha'} \rightarrow H = \sqrt{H} H^{\alpha\alpha'} \quad (3)$$

and

$$W^{\mu\alpha\mu'\alpha'} \rightarrow W^{\mu\mu'} \quad (4)$$

In this way we get a tensorial structure fully equivalent to that for (pseudo)scalar particles. Since we include both the vector and the tensor coupling for the low-energy vertex, the corresponding formulas are long and complicated and will be presented elsewhere [5].

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The effective quark dynamical Lagrangian of Nambu and Jona-Lasinio (NJL) [1] shares one important feature of QCD, chiral symmetry, which is spontaneously broken. It is therefore very well suited for the description of Goldstone bosons, like the pion, but because of its simplicity it is also quite attractive for the analysis of many other hadronic data. Since the NJL-model shows neither asymptotic freedom nor confinement, such calculations become more and more questionable with increasing time-like or space-like 4-momenta. Because of the missing confinement the a_1 meson, for instance, comes out to be unstable against the unphysical decay into free quarks [2]. Obviously, it is necessary to improve the model, if one wants to enlarge its range of validity.

In the NJL-model the quarks interact with each other directly via a point-like four-fermion coupling, whereas in QCD the interaction is of finite range (mediated by gluons). Since the explicit inclusion of gluonic degrees of freedom would destroy the simplicity of the model, we simulate the finite range effects by replacing the NJL-fourpoint interaction by a non-local Lagrangian:

$$\mathcal{L}_I = \int d^4x_1 \dots d^4x_4 \alpha(x; x_3, x_4; x_1, x_2) (\bar{\psi}(x_3) \gamma^\mu \lambda_k \psi(x_1)) (\bar{\psi}(x_4) \gamma_\mu \lambda_k \psi(x_2)) \quad (1)$$

The non-locality is given by the function

$$\alpha(x; x_3, x_4; x_1, x_2) = \int \frac{d^4k_1}{(2\pi)^4} \dots \frac{d^4k_4}{(2\pi)^4} e^{-i(k_1+k_2+k_3+k_4) \cdot x} \tilde{\alpha}(k_3, k_4; k_1, k_2) e^{i(k_1 \cdot x_1 + k_2 \cdot x_2 + k_3 \cdot x_3 + k_4 \cdot x_4)} \quad (2)$$

which guarantees the invariance of the model under translations. For the global internal symmetries it does, of course, not matter, whether the Lagrangian is local or non-local. In particular $\mathcal{L}_I$ is chirally symmetric.

Like in the original model, chiral symmetry is spontaneously broken. This generates a selfenergy Σ . In the Hartree-Fock approximation to the original NJL-model, Σ is interpreted as the

constituent quark mass M . The non-locality in coordinate space leads to momentum dependent vertices in momentum space (proportional to $\tilde{\alpha}$). As a consequence the selfenergy also becomes momentum dependent in our model. This makes the solution of the selfconsistency equation for Σ in general very difficult. Therefore we solve the problem — as a first step — in a schematic model with a separable interaction:

$$\tilde{\alpha}(k_3, k_4; k_1, k_2) = g_0 [f(k_3^2) f(k_4^2) f(k_1^2) f(k_2^2)]^{\frac{1}{2}} \quad (3)$$

This separability assumption allows to study the functional behaviour of the selfenergy analytically. We find

$$\Sigma(q^2) = -\frac{16}{9} g_0 f(q^2) (4 \text{infn}) \int \frac{d^4k}{(2\pi)^4} \frac{m_0 + \Sigma(k^2)}{k^2 - (m_0 + \Sigma(k^2))^2 + i\epsilon} f(k^2) \quad (4)$$

We can choose the overall constant g_0 such that

$$\Sigma(q^2) = f(q^2) \quad (5)$$

For the function $f(q^2)$ we choose

$$f^2(q^2) = \left(\frac{q^2 + a^2}{2} \right) + \sqrt{\left(\frac{q^2 + a^2}{2} \right)^2 + b^4} \quad (6)$$

with a and b being free parameters.

From eqs. (5) and (6) it follows, that the quarks never reach their mass shell, because $\Sigma^2(q^2)$ is always larger than q^2 . Consequently the unphysical emission of free quarks into the continuum is not possible in our model: All mesons are stable against quark-antiquark decay.

Table 1 shows some first results for $m_0 = 7$ MeV. All calculations have been done consistently with the same set of parameters. The coupling strength of the different channels have been determined from the interaction Lagrangian (1) by a Fierz transformation. In the local NJL-model this procedure leads to a too weak coupling constant in the vector channel. In our non-local model it works

very well. This could be interpreted as a justification for the parametrization (6).

Finally it should be pointed out again, that chiral invariance is not destroyed by making the interaction non-local. In the chiral limit $m_0 = 0$ it can be shown analytically that the pion mass becomes zero in our model, consistent with Goldstone's theorem.

After getting rid of the unphysical quark-antiquark decay channels of the mesons, the physical decay channels, like $\rho \rightarrow \pi\pi$, are required to describe the finite meson widths. This work is in progress.

	non-local model	experiment
$m_{\pi^0}/m_{\pi^\pm}$	139.2 MeV	135.0 / 139.6 MeV
m_ρ/m_ω	782 MeV	768 / 782 MeV
m_{a_1}/m_{f_1}	1290 MeV	1260 / 1282 MeV
$\langle \bar{u}u \rangle$	$-(225 \text{ MeV})^3$	$-(225 \text{ MeV})^3$
f_π	91 MeV	93 MeV

Table 1: Meson masses, the vacuum condensate and the pion decay constant, obtained within our non-local schematic model, compared with the corresponding experimental values. The theoretical calculations have been performed with the parameters $a = 235 \text{ MeV}$ and $b = 435 \text{ MeV}$ (see eq. (6)). The current quark mass is $m_0 = 7 \text{ MeV}$. The divergent integrals have been regularized using a sharp cutoff $\Lambda = 1200 \text{ MeV}$ in four dimensional Euclidean space. All mesons are stable against quark-antiquark decay.

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Renormalization of the ρ -meson mass due to tad poles

R. Tegen* and J. Speth

We consider here the contribution to the self-energy of the ρ -meson to lowest order in the strong coupling g which is given by the Lagrangian L of Bando et al.[1]; these contributions are depicted by the first two lines of fig.1.

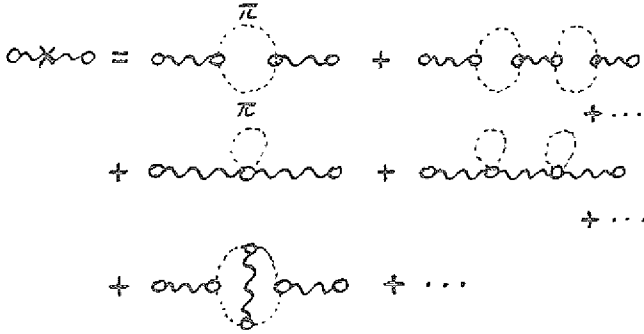


Fig.1

The ρ - $\pi\pi$ coupling is $-ig(p^\mu + p'^\mu)$ whereas the contact term (also named "tad pole" or "sea gull") $\rho\rho\pi\pi$ is $2ig^2 g^{\mu\nu}$, see fig.2

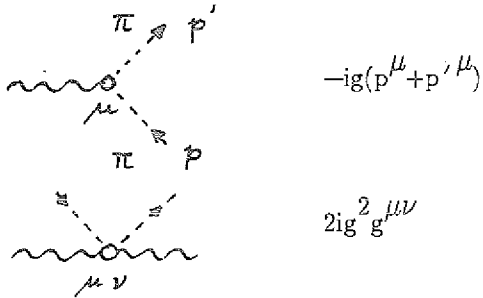


Fig.2

Note that the contact term $\rho\rho\pi\pi$ arises from the kinetic term in L which is *quadratic* in the covariant derivative ∂^μ ; local gauge invariance (gauge bosons are the ρ -mesons) introduces couplings of the gauge boson to the fields in L via minimal substitution, $\partial^\mu \rightarrow \partial^\mu - ig V^\mu$. The kinetic term of the fermion fields is *linear* in ∂^μ so that there are no fermionic contact terms (see: Compton scattering).

The bare ρ -propagator on the r.h.s. of fig.1 is given by

$$-i \frac{g^{\mu\nu} - k^\mu k^\nu / m_\rho^2}{k^2 - m_\rho^2 + i\epsilon} \quad (1)$$

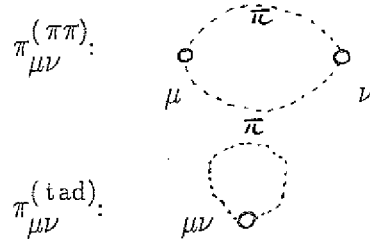
and the pion-propagator is $i/(k^2 - m_\pi^2 + i\epsilon)$, as usual. It is then easy to sum up the infinitely many contributions in fig.1 (geometric progression) with the result

$$-i \frac{g_{\mu\nu} - k_\mu k_\nu (1 - \pi(k^2)) / m_\rho^2}{k^2 - m_\rho^2 - k^2 \pi(k^2) + i\epsilon} \quad (2)$$

where $\pi(k^2)$ is defined by

$$\pi_{\mu\nu} = \pi_{\mu\nu}^{(\pi\pi)} + \pi_{\mu\nu}^{(\text{tad})} = i(k^2 g_{\mu\nu} - k_\mu k_\nu) \pi(k^2) \quad (3)$$

The tensors are given by the individual loops



or, specifically, by the (divergent) integrals

$$\pi_{\mu\nu}^{(\pi\pi)} = g^2 \int d^4 p \frac{(2p+k)_\mu (2p+k)_\nu}{((p+k)^2 - m_\pi^2 + i\epsilon)(p^2 - m_\pi^2 + i\epsilon)} \quad (4)$$

$$\pi_{\mu\nu}^{(\text{tad})} = -2g^2 \int d^4 p \frac{1}{p^2 - m_\pi^2 + i\epsilon} g_{\mu\nu} \quad (5)$$

To show the gauge invariance of the expression (3), i.e. $k^\mu \pi_{\mu\nu} = 0 = k^\nu \pi_{\mu\nu}$, we evaluate (4) and (5) in the dimensional regularization scheme. The idea (in a nut-shell) is to isolate the divergence of certain

integrals as poles in the variable dimension $n=4-\epsilon$ and make use of cancellations in the sum of individually diverging integrals; the remaining divergence is then absorbed into one physical observable (here the mass). Higher dimensional integrals have to be transformed to Euclidean space where relative angles can be given a well-defined meaning, even in higher than 3 dimensions. Then it is straightforward to calculate the following class of integrals

$$I(n; r, m; R^2) = \int \frac{d^n k}{(2\pi)^n} \frac{(k^2)^r}{(k^2 - R^2 + i\epsilon)^m}$$

using the symmetric integration rules and the n -dimensional γ -algebra as given in [2]. We obtain the result [3]

$$\pi(k^2) = -\left[\frac{g^2}{4\pi}\right] \frac{1}{12\pi} \Gamma(2-n/2) \quad (6)$$

We want to compare (6) with the covariant (or: η -) regularisation. This approach is due to Schwinger [4], see ref. [5] for details. We obtain

$$\pi_{\mu\nu}^{(\text{tad})} = 2g^2 g_{\mu\nu} i/(4\pi)^2 1/\eta E_2(\eta m_\pi^2) \quad (7)$$

where the divergent integral $\int_0^\infty ds/s^2 e^{-sm^2}$ has been regularized, in accordance with the Schwinger method, by introducing a lower limit cut-off,

$$\int_\eta^\infty ds/s^2 e^{-sm^2} = 1/\eta E_2(m_\pi^2/\eta)$$

with the exponential integral of order 2 [6]. In an analogous way we obtain $\pi_{\mu\nu}^{(\pi\pi)}$,

$$\begin{aligned} \pi_{\mu\nu}^{(\pi\pi)} = & (g^2/3) i/(4\pi)^2 \{ E_1(\eta |k^2/6 - m_\pi^2|) (k_\mu k_\nu - \\ & k^2 g_{\mu\nu}) + g_{\mu\nu} (k^2 E_1(\eta |k^2/6 - m_\pi^2|) \\ & + 6/\eta E_2(\eta (k^2/6 - m_\pi^2))) \} \end{aligned} \quad (8)$$

The sum of (7) and (8) is gauge invariant; this fact fixes η as a function of m_π^2 and k^2 , see [5]

$$y E_1(z) + \frac{6}{x} (E_2(z) + E_2(x)) = 0 \quad (9)$$

$$\text{with } z = \eta (k^2/6 - m_\pi^2), \quad x = \eta m_\pi^2, \quad y = k^2/m_\pi^2.$$

We therefore obtain the result

$$\pi(k^2) = -\frac{g^2}{4\pi} \frac{1}{12\pi} E_1(x | \frac{k^2}{6m_\pi^2} - 1 |) \quad (10)$$

In ref. 3 we have used a non-covariant regularization,

$$F^{(s)} = \frac{\Lambda^2 + m_\rho^2}{\Lambda^2 + \omega_p^2}, \quad \omega_p^2 = p^2 + m_\rho^2$$

This form factor enters twice in $\pi_{\mu\nu}$, one for each $\rho\pi\pi$ -vertex. To calculate $\pi_{\mu\nu}$ we split the $d^4 p$ into a dp_0 - and a $d^3 p$ -integration. Using Cauchy's theorem for the dp_0 -integration one obtains for the integral in $\pi_{\mu\nu}^{(\text{tad})}$,

$$\begin{aligned} & \int \frac{d^4 p}{(2\pi)^4} \frac{(\Lambda^2 + m_\rho^2)^2}{(p^2 - m_\pi^2 + i\epsilon)(\Lambda^2 + \omega_p^2)^2} \\ &= -\frac{i}{(2\pi)^2} \frac{(\Lambda^2 + m_\rho^2)^2}{2\Lambda^2} \\ & \quad \left\{ 1 - \frac{m_\pi^2}{2\Lambda \sqrt{\Lambda^2 + m_\rho^2}} \ln \frac{\sqrt{\Lambda^2 + m_\rho^2} + \Lambda}{\sqrt{\Lambda^2 + m_\rho^2} - \Lambda} \right\} \end{aligned} \quad (11)$$

Given Λ and $m_\rho^{(0)}$ as in ref [3] we obtain the equivalent x in (10) from a comparison of (7) with (11), which is a transcendental equation in x ,

$$E_2(x)/x = \frac{1+2zr}{z} \left(1 - \frac{z}{\sqrt{1+2zr}} \ln \frac{\sqrt{1+2zr}+1}{\sqrt{1+2zr}-1} \right) \quad (12)$$

with $r = m_\rho^{(0)2}/m_\pi^2$, $z = m_\pi^2/2\Lambda^2$. In this scheme

$\pi(k^2)$ is given by (10) where x is now a function of $(\Lambda, m_\rho^{(0)})$ according to eq.(12). For small x , $x \ll 1$,

one can approximate $E_2(x) \approx e^{-x} - x E_1(x) \approx$

$1 + x \ln x - x(1-C) - (3/2)x^2 + O(x^3)$, where the last relation holds because of the approximation $E_1(x)$

$\approx -x - \ln x - C$ with $-C = E_1(1) - 1 = -0.781$ Euler's

number [26].

The pole of the propagator (2) defines the renormalized mass $m_{\rho}^2(1-\pi(m_{\rho}^2)) = m_{\rho}^{(0)2}$.

Two remarks are in order here:

The tad-pole contribution (5) to the tensor $\pi_{\mu\nu}$ (3) cancels the corresponding term in $\pi_{\mu\nu}^{(\pi\pi)}$ such that the sum is gauge invariant as in (3). If the individual contribution of the pion loop to the ρ -meson mass is repulsive, then the tad-pole contribution will always be attractive such that a (partial) cancellation of their respective contributions to $\pi(k^2)$ occurs. We demonstrate this for

$$g^2/4\pi=2.8, \text{ taken from } \Gamma_{\rho \rightarrow \pi\pi} = \frac{2}{3} \frac{p^3}{m_{\rho}^2} \frac{g^2}{4\pi} = 153$$

MeV, ($p^2 = \frac{1}{4}(m_{\rho}^2 - 4m_{\pi}^2)$), and $\Lambda \approx 850$ MeV. For these values the resultant $\pi(m_{\rho}^2)$ is very small, of order -0.1 , so that the physical mass is close to the bare mass. Higher values of Λ have been discussed in [3].

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Geometrical picture of coherent $\gamma\text{--}\gamma$ and $\gamma\text{--}g$ processes in ultrarelativistic heavy ion collisions

G. Baur

Coherent particle production via the $\gamma\text{--}\gamma$ mechanism occurs in disruptive as well as non-disruptive relativistic heavy ion collisions [1,2]. Due to the enhancement factor of Z^4 , relativistic heavy ion colliders may compete with e^+e^- colliders in the study of $\gamma\text{--}\gamma$ physics. To reduce the strong background, non-disruptive collisions must be selected.

In disruptive collisions, the dilepton production via coherent $\gamma\text{--}\gamma$ processes is calculated as a function of the impact parameter. Cross-sections are large, comparable to those expected from thermal and Drell-Yan mechanisms. However, the very low $p_\perp$ -values will allow a discrimination. The characteristic features of the processes are discussed and the importance of the mechanism at SPS, RHIC and LHC energies is assessed.

In addition to the $\gamma\text{--}\gamma$ fusion process, $\gamma+g$ fusion, where the photon is coherently produced can also play a role [3], e.g. leading to the production of $c\bar{c}$ and $b\bar{b}$ pairs. It was observed [3] that the production of $c\bar{c}$ pairs via $\gamma\text{--}g$ fusion at RHIC could be of interest to obtain information on the gluon structure function in nuclei. The restriction imposed by strong absorption can be conveniently taken care of in the present geometrical approach [1,2]. (The photon is produced "outside" of one of the nuclei, the gluon is produced "inside" the other, and the nuclei do not interact strongly with each other). More detailed calculations of the $\gamma\text{--}g$ fusion process, especially $c\bar{c}$ production at RHIC energies is in progress.

The intense electromagnetic interactions of colliding heavy ions results in a new type of background: copious electron-positron pair production. This is of relevance e.g. for RHIC [4] and LHC (Pb-Pb collisions) [5]. A rather simple and accurate formula for single electron (or positron) production at 90° is found:

$$\frac{d\sigma(\text{QED})}{d\epsilon d\Omega} = (Z\alpha)^4 \frac{8}{3\pi^2} \left(\frac{p}{\epsilon}\right) \frac{1}{\epsilon^3} \left(\ln \frac{0.681\gamma m_e}{\epsilon}\right)^2 \quad (1)$$

where ϵ and p denote the energy and momentum of the observed lepton, Z is the nuclear charge, γ the

Lorentz-factor. This formula is valid for $\epsilon \gtrsim \gamma m_e$, i.e. for lepton energies up to the GeV region (for LHC, with $\gamma = 3400$). Calculations are done [6] for all angles, in order to assess the importance of these QED electrons for the detector design at LHC [5].

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Lepton pairs play a very important role in the diagnostics of the hot hadronic matter formed in relativistic heavy ion collisions. There are various physical sources which contribute, like the Drell–Yan mechanism, preequilibrium pairs or thermal emission of lepton pairs from the quark–gluon–plasma, and which have been and will be studied, see e.g. Refs. [1,2].

The mechanism of coherent $\gamma\text{--}\gamma$ lepton pair production in relativistic heavy ion collisions is of continuing interest. In central, disruptive collisions with impact parameter $b < R_1 + R_2$ coherent pairs are a potential background to the other mechanisms, while in non-disruptive collisions $b > R_1 + R_2$ the $\gamma\text{--}\gamma$ process can be studied in its pure form.

The semi classical approach [3] provides a direct treatment of the impact parameter dependence of the process. Double differential cross sections as well as impact parameter dependent creation probability distributions can be calculated in a straight forward way [4] for energies of existing and future colliders.

In the following we restrict ourselves on the ^{208}Pb – ^{208}Pb system. Fig.1 shows the results for the differential cross sections for coherent e^+e^- production in ^{208}Pb – ^{208}Pb collisions at rapidity $Y=0$ as a function of the invariant mass M at SPS, RHIC and LHC. Below 0.5 GeV all differential cross sections are of the order of b/GeV to mb/GeV . Above 1 GeV the cross section for SPS breaks down while the others still contribute. This is a result of the general fact that for increasing Lorentz factor the available invariant mass regions do also increase.

In Fig.1a the differential cross section for disruptive collisions at SPS, RHIC and LHC is shown. For low invariant masses all cross sections tend to the same values. The cross sections for LHC and RHIC do not differ significantly below 1 GeV. This is a result of the saturation property for disruptive collisions for increasing Lorentz factors.

As expected, for higher invariant masses the SPS cross section is several orders of magnitude lower.

Fig.1b shows the differential cross sections for non-disruptive collisions. All cross sections are enhanced nearly one order of magnitude in com-

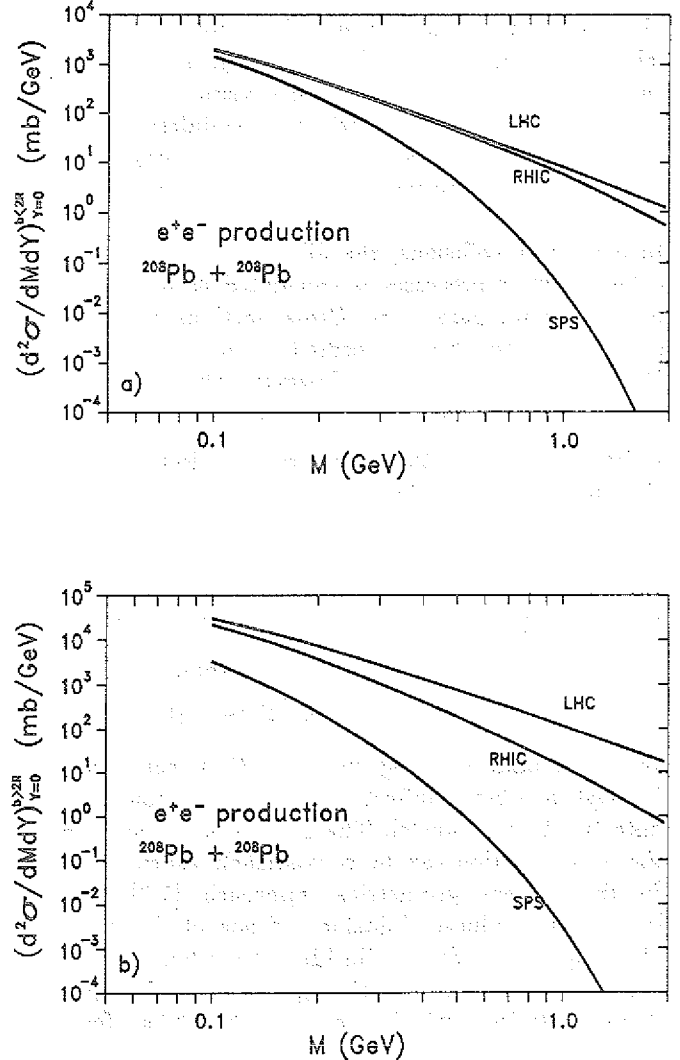
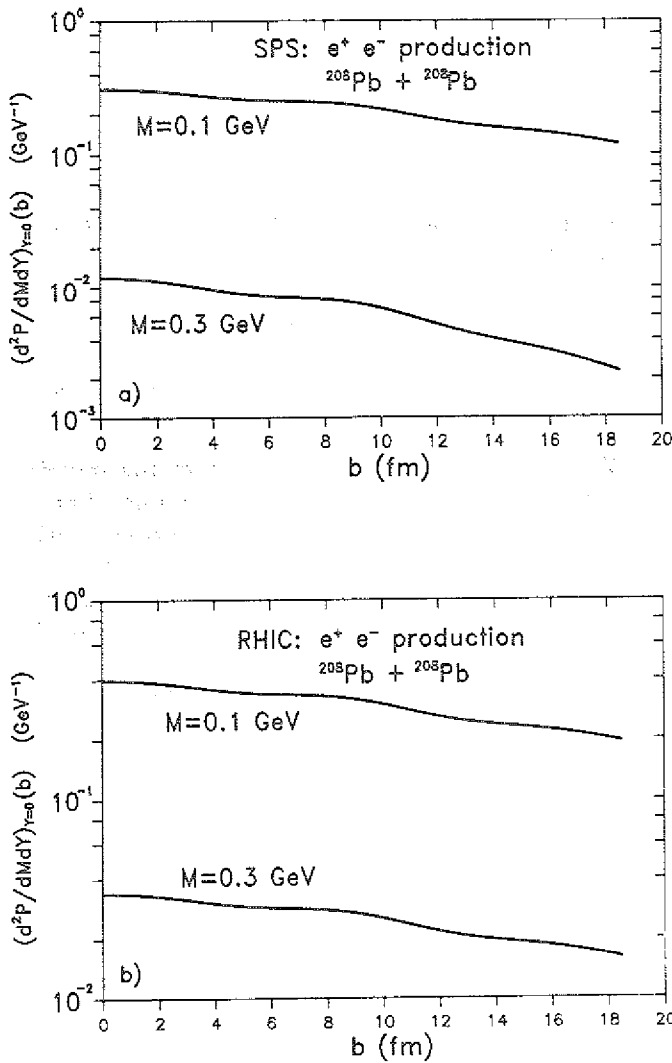


Fig. 1. Differential cross sections for coherent e^+e^- production in ^{208}Pb – ^{208}Pb collisions at rapidity $Y=0$ as a function of the invariant mass M at SPS, RHIC and LHC. a) disruptive collisions: $b < 2R$, b) non-disruptive collisions: $b > 2R$.

parison to the disruptive ones. This is a result of the unrestricted impact parameter range for those collisions. For the same reason the saturation property for disruptive collisions is not observed for non-disruptive ones: For very low values of the invariant mass the cross sections of RHIC and LHC do tend to the same values but for invariant masses above 0.3 GeV they are well separated. For invariant masses above 1 GeV the LHC cross section is more than one order of magnitude bigger than the RHIC cross section which itself is several orders of magnitude bigger than the SPS one.



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Fig. 2. Coherent e^+e^- production probability for $Y=0$ plotted as a function of the impact parameter b of the colliding $^{208}\text{Pb}-^{208}\text{Pb}$ nuclei at SPS (Fig. 2a) and RHIC (Fig.2b).

In Fig.2 the coherent e^+e^- production probability is plotted as a function of the impact parameter b of the colliding nuclei for $^{208}\text{Pb}-^{208}\text{Pb}$ collisions at SPS and RHIC. For $Y=0$ we plotted the distributions for two characteristic values of the invariant mass M .

For SPS (see Fig.2a) as well as for RHIC (see Fig.2b) the probability distributions smoothly decrease from a maximum value at $b = 0$. The distributions vary slowly and there are no significant characteristics at $b = 2R$. The absolute values of the probability distribution for RHIC are enhanced and the decrease for large b is not as strong as compared to SPS.

In the framework of the semi classical approach we are able to obtain a reliable guide through the process of coherent $\gamma-\gamma$ lepton pair production in relativistic heavy ion collisions [4] which explains the main characteristics of the process.

The production of electron positron pairs in $\pi^- p$ collisions at pion momenta of 16 GeV/c and 17 GeV/c has been measured almost 10 years ago [1,2]. Several mechanisms which have been proposed failed to explain the data (see for example [3,4]). Bottcher et al. [5] proposed the coherent $\gamma\gamma$ mechanism resulting in a good over all agreement with the data. These calculations seem to indicate a strong dependence of the cross section on the anomalous magnetic moment of the proton (i.e. the "spin flip" components of the process).

In the equivalent photon method the cross sections for coherent electron positron production can be calculated analytically since this is possible for the photon spectra of the pion and the proton [6]. In this approach we are able to show the significance of recoil contributions for the pion, which are neglected in Ref. [5]. Even the influence of the anomalous magnetic moment of the proton can be studied in a very transparent way.

Fig. 1 shows the equivalent photon spectrum of the pion in the $\pi^- p$ c.m. system for pion lab-momenta of 16 GeV/c as a function of the equivalent photon energy. The dashed line corresponds to the straight line approximation also used in Ref. [5], while the full line takes into account energy loss up to ω/E .

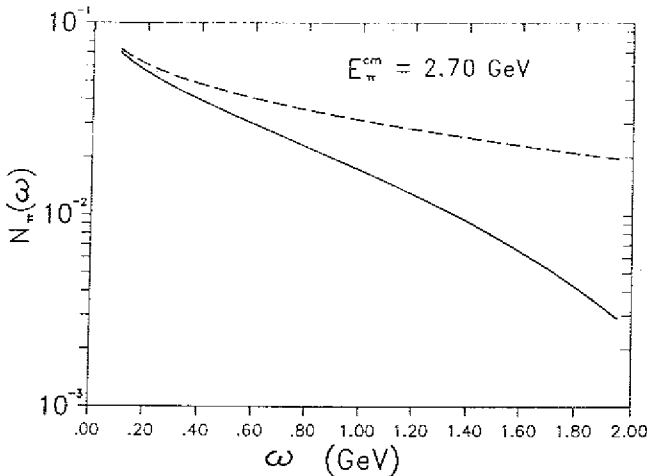


Fig. 1. Equivalent photon spectrum of the pion in the $\pi^- p$ c.m. system for 2.7 GeV as a function of the equivalent photon energy ω . The dashed line corresponds to the straight line approximation, the full one takes into account energy loss up to ω/E .

Even for low photon energies the difference is significant, for energies above 1.5 GeV the photon

spectrum including energy loss starts to decrease very fast. This leads to differences to the straight line approximation of nearly one order of magnitude for energies up to 2 GeV.

Fig. 2 shows the double differential cross section for electron positron production at Feynman $x_F = 0$ in the $\pi^- p$ c.m. system as a function of the invariant mass W . In the calculation of the dashed line the magnetic moment of the proton is neglected, while it is taken into account for the full line. For invariant masses below 1.5 GeV the discrepancy between the two calculations is not so strong. Not until after invariant masses of 2 GeV are reached, the cross section for vanishing proton magnetic moment starts to decrease much stronger than the one including the anomalous magnetic moment.

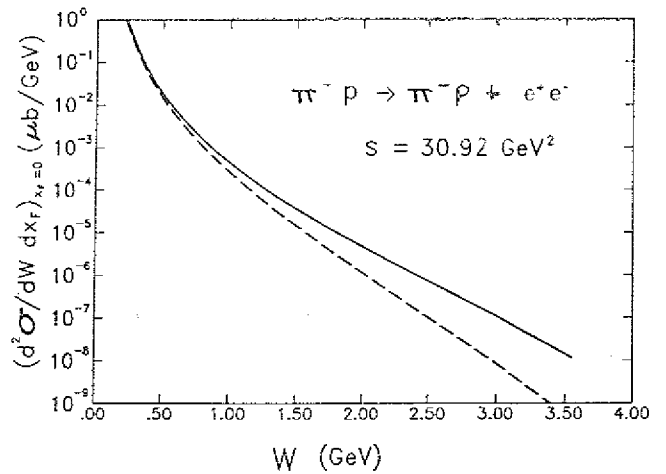


Fig. 2. Double differential cross section for electron positron production at Feynman $x_F = 0$ in the $\pi^- p$ c.m. system as a function of the invariant mass W for vanishing proton magnetic moment (dashed line) for the anomalous physical moment (full line).

Integrating over the invariant mass W in the same kinematical region used by Bottcher et al. ($0.2 \text{ GeV} < W < 1.2 \text{ GeV}$) leads to the differential cross section $d\sigma/dx_F$ as a function of the dimensionless Feynman x_F ; this is shown in Fig. 3. Again the dashed line corresponds to vanishing proton magnetic moment, the full line to the anomalous one. Over the whole range of possible values of x_F the difference between both calculations is found to

be not significant. The decrease of the cross section above $x_F = 0.7$ originates from the limited range of kinematically allowed values of x_F in consideration of pion recoil.

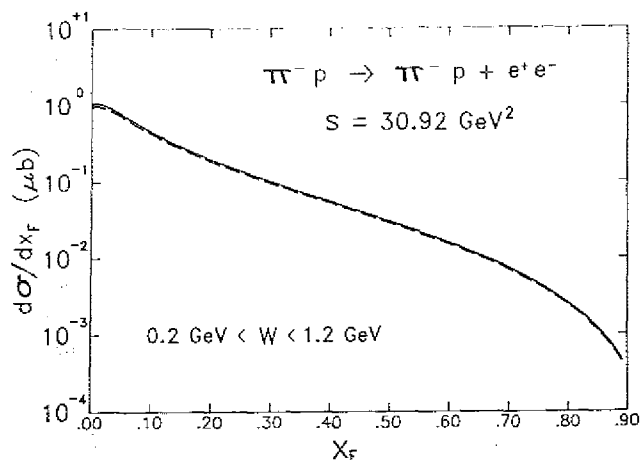


Fig. 3. Differential cross section for electron positron production as a function of Feynman x_F in the kinematical region $0.2 \text{ GeV} < W < 1.2 \text{ GeV}$ with the same label as in Fig. 2.

In Ref. [6] π^-p collisions are discussed in further detail within the equivalent photon method. This method provides a clear and transparent understanding of the relevant electromagnetic processes in π^-p collisions.

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The spin—longitudinal/spin—transverse response of ^{40}Ca to 500 MeV protons

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The study of the spin— isospin response of nuclei has been an area of intense interest, both, theoretically and experimentally. The total nuclear spin—response can be split in a spin—longitudinal and a spin—transverse component, closely related to the exchange of π and ρ mesons, respectively. Inelastic hadron—nucleus scattering is an appropriate tool for such studies. Energy and momentum can be transferred independently on the nucleus. The scattering process is sensitive to collective excitations of the nucleus. Both, spin—longitudinal as well as spin—transverse transitions are excited, whereas electron and pion scattering are sensitive only to spin—transverse excitations of the nucleus.

The availability of polarized proton beams made it possible to analyze the spin response in (p,p') scattering experiments. The inclusive cross sections as well as the spin—transfer coefficients are measured. This reaction is dominated by $\Delta S=0$ transitions, due to the strong scalar— isoscalar component of the projectile—target interaction. The spin—flip probability $S_{nn'}$, determined by the normal—component spin—transfer coefficient $D_{nn'}$, ($S_{nn'} = 1/2(1 - D_{nn'})$), can single out $\Delta S = 1$ transitions. The further determination of $D_{11'}$ and $D_{ss'}$ (longitudinal and sideways spin—transfer coefficients) allow the investigation of the longitudinal and transverse spin—flip probabilities S_L and S_T , respectively, given by:

$$S_L = \frac{1}{4}[1 - D_{nn'} + (D_{ss'} - D_{11'})]$$

$$S_T = \frac{1}{4}[1 - D_{nn'} - (D_{ss'} - D_{11'})]$$

Their sum gives again the total spin—flip probability $S_{nn'}$. So spin—longitudinal and spin—transverse excitations in the nucleus can be studied.

Such experiments with a complete set of spin observables have been performed recently at LAMPF on ^{40}Ca at 497 MeV incident energy, a momentum transfer of 0.56 fm^{-1} , corresponding to a scattering angle of 5.5° , and excitation energies ω between 30 and 40 MeV [1]. At this momentum transfer the pion—field is not expected to be strongly enhanced, and ordinary nuclear collectivity

should still be dominant. The excitation range corresponds to the region where the total nuclear spin response has found to be enhanced relative to the free NN scattering [2].

To analyze the data microscopic RPA—DWIA calculations have been performed for this reaction [3]. The single particle basis for the RPA is generated by a Woods—Saxon potential [4], and the residual interaction is based on the π and ρ exchange. The Franey—Love t -matrix [5] has been used as projectile—target interaction, and the optical potential has been taken from [6]. All multipoles up to $J=5$ have been summed up, which gives good convergence.

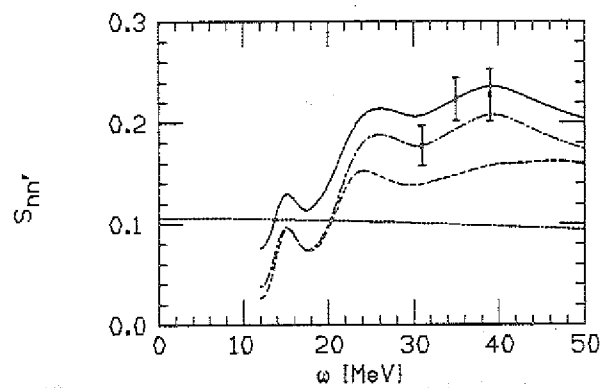


Fig. 1. Spin—flip probability $S_{nn'}$ for the reaction $^{40}\text{Ca} (p,p')$ at $T_{\text{lab}} = 497 \text{ MeV}$ and a scattering angle of 5.5° as function of excitation energy ω . The full curve represents the full RPA—DWIA calculation. For the dashed—dotted curve only central distortions have been used, and the dashed curve denotes the Plane Wave calculation. The dotted curve stands for free NN scattering (Fermi gas).

Fig. 1 shows the spin—flip probability $S_{nn'}$. The full curve denotes the full RPA—DWIA calculation, including central as well as spin—orbit distortions. The dotted curve represents free NN scattering (Fermi gas) for comparison. The calculation is in good agreement with the data. At lower excitation energies ($\omega \leq 20 \text{ MeV}$) $S_{nn'}$ is lower than the free values, due to the exhaustion of $\Delta S = 0$ strength in the giant resonance region. The spin—dipole resonance at about $\omega = 25 \text{ MeV}$ and the spin—quadrupole resonance at about 40 MeV lead

then to an enhancement of S_{nn} , at higher ω above the giant resonance region [3,7]. A comparison with the Plane Wave calculation (dashed curve) shows that S_{nn} is very sensitive to distortions of the projectile wave. Similar results have been obtained for different incident energies [3,8]. The dashed-dotted curve includes central distortions only. So both, central as well as spin-orbit distortions contribute to this effect.

For a better understanding of this, Fig. 2 shows a decomposition of S_{nn} into the spin-longitudinal

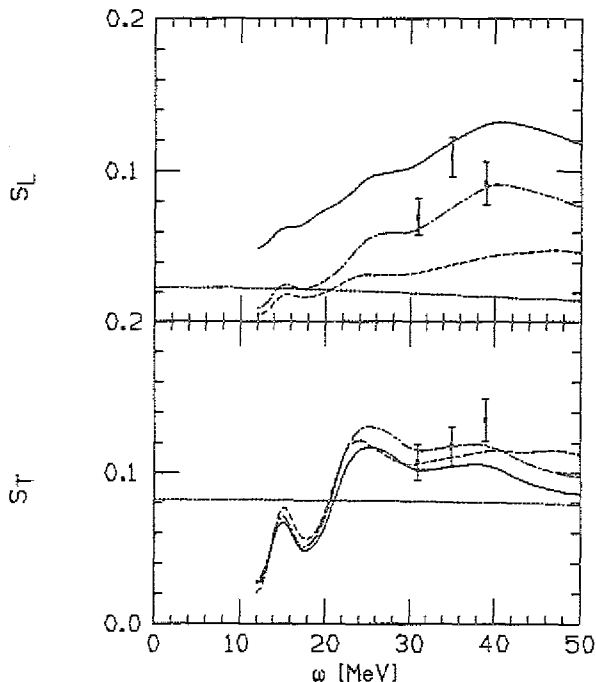


Fig. 2. Same as for Fig. 1 for S_L and S_T .

and spin-transverse contributions S_L and S_T , respectively. For S_L the effect of distortions is strikingly strong. At 50 MeV excitation energy they lead to an enhancement by more than a factor of 2 compared to the Plane Wave calculation (dashed curves). The effect on S_T is more moderate but goes in the opposite direction. So the ratio S_L/S_T is increased drastically by the distortions (Fig. 3). This makes it — for the moment — practically impossible to extract any information about an enhanced pion field in the nucleus by this kind of measurements.

The different sensitivity in the longitudinal and transverse channels is, as far as central distortions are concerned, related to the different ranges of the π and ρ exchange and the absorption by the imaginary part of the central optical potential.

Previous calculations for $\bar{p}$ -nucleus scattering have resulted in similar results [3]. For that reaction the

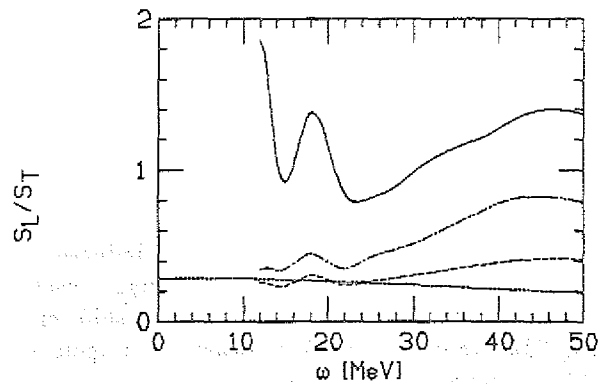


Fig. 3. Same as for Fig. 1 for S_L/S_T .

enhancement of S_L/S_T is much larger because of the stronger absorption of the projectile wave due to the annihilation.

The effect of spin-orbit distortions is surprisingly high. They account for approximately half of the total effect and lead to an enhancement of S_L while S_T is decreased. This goes back to the fact that for the spin-dipole as well as for the spin-quadrupole (both are dominant for S_L and S_T) the longitudinal and the transverse formfactor, i.e. the projectile-target interaction folded with the transition density, peak on different sides of the nuclear surface, also because of the different range of π and ρ exchange.

A more detailed analysis of distortion effects is in progress.

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The Δ -excitations in nuclei induced by intermediate energy charge exchange reactions reveal a systematic downward energy shift of the $\Delta(1232)$ -resonance peak position for targets with $A \geq 12$ compared to the peak position for a proton target [1,2]. The non-trivial part of this shift is caused by a nuclear medium correlation effect on the spin-longitudinal response function [3,4]. No shift occurs in the spin-transverse $(\vec{S} \times \vec{q} \cdot \vec{T})$ channel, since the interaction is here weakly repulsive.

The inclusive charge exchange cross sections contain information on both the spin-longitudinal (LO) and spin-transverse (TR) nuclear responses. These two responses can be experimentally separated by measuring the spin observables.

Therefore, the $(\vec{d},2p)$ reaction with a 2 GeV tensor polarized deuteron beam [5] has been used to analyze the spin-structure of the Δ -excitation and the spin-response of the Δ -propagation in nuclei [6]. The analysis of the $p(\vec{d},2p)\Delta^0$ -reaction enabled us to fix $t_{NN,N\Delta}$ as

$$t_{NN,N\Delta} = t'_{N\Delta} J_{\pi N\Delta} \left[\frac{\Lambda'^2 - m_\pi^2}{\Lambda_\pi'^2 - t} \right]^2$$

$$[V_{LO} \sigma_1 \cdot \hat{q} \vec{S}_2^\dagger \cdot \hat{q} + V_{TR} (\vec{\sigma}_1 \times \hat{q}) \cdot (\vec{S}_2^\dagger \times \hat{q})]$$

$$\vec{\tau}_1 \cdot \vec{T}_2 \quad (1)$$

with $V_{LO} = V_{TR}$. The other quantities in this equation are defined as $J_{\pi N\Delta} = 4\pi\hbar^2 f_{\pi NN} f_{\pi N\Delta} / m_\pi^2 \approx 800 \text{ MeV fm}^3$, $t'_{N\Delta} = 0.33$ and $\Lambda'_\pi = 650 \text{ MeV}$.

As measurements of the $p(\vec{p},\vec{n})\Delta^{++}$ -reactions have been performed years ago and measurements of the $^{12}\text{C}(\vec{p},\vec{n})$ -reaction are currently being planned, we calculated the spin-transfer coefficients C_{XX} and C_{ZZ} of these reactions to confirm our analysis of the driving interaction $t_{NN,N\Delta}$ and to get more information on the spin-response.

The spin-transfer coefficients of the $(\vec{p},\vec{n})$ experiments are defined as follows:

$$C_{XX} = \frac{-|V_{LO}|^2}{|V_{LO}|^2 + 2|V_{TR}|^2} \quad (2)$$

$$C_{ZZ} = \frac{|V_{LO}|^2 - 2|V_{TR}|^2}{|V_{LO}|^2 + 2|V_{TR}|^2} \quad (3)$$

In figs. 1 and 2 the calculated spin transfer coefficients C_{XX} and C_{ZZ} of the $p(\vec{p},\vec{n})\Delta^{++}$

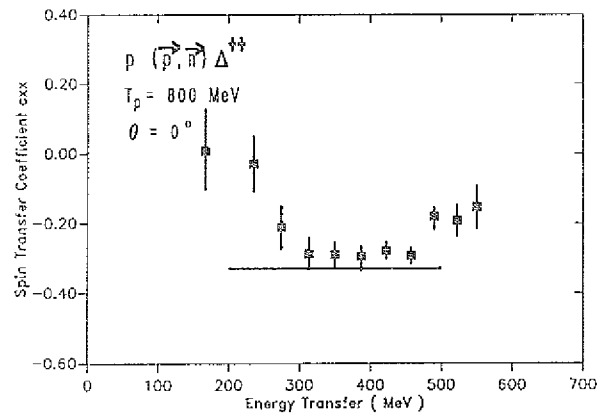


Fig. 1. Spin transfer coefficient C_{XX} for the $p(\vec{p},\vec{n})\Delta^{++}$ -reaction at $E = 800 \text{ MeV}$ and $\theta = 0^\circ$. Data are from ref. 7, and the solid line represents our calculation with $t_{NN,N\Delta}$ from eq. (1).

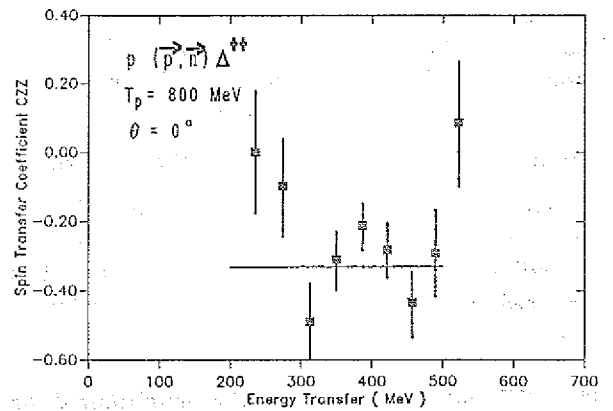


Fig. 2. As fig. 1, but spin transfer coefficient C_{ZZ} .

reaction at a proton incident energy of 800 MeV are compared with the zero degree data of [7]. These data can roughly be described by our calculations, where the $t_{NN,N\Delta}$ from eq. (1) was used, and thus this is another confirmation for the spin-structure of the $t_{NN,N\Delta}$ -amplitude fixed in [6] by analysis of the $p(\vec{d},2p)\Delta^0$ -reaction.

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Comparison of different models for the ΔN^{-1} -residual interaction
in the Δ -propagation in nuclei

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Different models have been used for the residual interaction in Δ -hole calculations. All these models are based on the π -exchange model, because the pion is the dominant feature in this excitation energy range due to its pole-structure. In momentum space the π -exchange potential is given by:

$$V_{\pi}(\omega, \vec{q}) = 4\pi\hbar c \frac{f_{\pi N \Delta}^2}{M_{\pi}^2} \left[\frac{\Lambda_{\pi}^2 - M_{\pi}^2}{\Lambda_{\pi}^2 - t} \right]^2 \frac{1}{t - M_{\pi}^2} (\vec{S}_2^+ \cdot \vec{q})(\vec{S}_1 \cdot \vec{q})(\vec{T}_2^+ \cdot \vec{T}_1) \quad (1)$$

where $t = \omega^2 - \vec{q}^2$, $f_{\pi N \Delta}$ is the coupling constant, and Λ_{π} is the cutoff parameter of the π -formfactor.

Moniz et al. [1,2] have used a π -exchange potential with a cutoff parameter $\Lambda_{\pi} = 300$ MeV in their calculations of the pion- and photoabsorption. The π -exchange potential (1) is purely spin-longitudinal (LO). In the preceding report we have seen that a photon excites the Δ spin-transversal (TR). Therefore the electromagnetic excited Δ is not affected by this residual interaction.

Udagawa et al. [3] used a π -exchange potential as residual interaction with a cutoff parameter $\Lambda_{\pi} = 780$ MeV. Additionally they had to introduce a Landau-Migdal term, which describes the short range correlations, with a strength of $g'_{\Delta\Delta} = 1/3$ in units of $J_{\pi\Delta\Delta} = 4\pi\hbar c f_{\pi N \Delta}^2 / m_{\pi}^2 \approx 1600$ MeV fm³:

$$V_{\delta}(\omega, \vec{q}) = 4\pi\hbar c \frac{f_{\pi N \Delta}^2}{M_{\pi}^2} \left[\frac{\Lambda_{\pi}^2 - M_{\pi}^2}{\Lambda_{\pi}^2 - t} \right]^2 g'_{\Delta\Delta} (\vec{S}_2^+ \cdot \vec{S}_1)(\vec{T}_2^+ \cdot \vec{T}_1) \quad (2)$$

This additional term is necessary, because the pure π -exchange potential with $\Lambda_{\pi} = 780$ MeV is too attractive in the (LO) channel and shifts the resonance to too low energies.

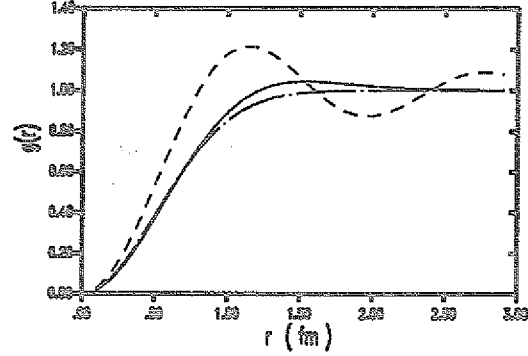


Fig. 1. NN-correlation functions in the range from 0 to 3 fm. The dashed curve is given by eq. (4). The dash-dotted and solid curves show other parametrizations of the correlation function.

Another way to introduce the short range correlations in the π -exchange model is to multiply the π -exchange potential in coordinate space with a correlation function $g(r)$ [4]:

$$\hat{V}_{\pi}(\omega, r) = V_{\pi}(\omega, r) \cdot g(r) \quad (3)$$

Fig. 1 shows three typical correlation functions, which have been used in studies on the NN-interaction. The dashed curve is given by

$$g(r) = 1 - j_0(q_c r) \quad (4)$$

with q_c being the Compton wavelength of the ω -meson and j_0 being the spherical Bessel function of the first kind. The other two curves are other parametrizations of the correlation function.

The effect on the residual interaction by adding either the Landau-Migdal term or by multiplying the π -exchange potential with a correlation function is shown in Fig. 2. The potentials in Fig. 2 are calculated in the (LO) channel for a energy transfer of $\omega = 300$ MeV. The pure π -exchange potential ($\Lambda_{\pi} = 780$ MeV) is drawn as the short dashed curve and the Landau-Migdal term is plotted as the dash-dotted curve. The long dashed curve represents the sum of the two former potentials. The solid curve shows the "correlated" π -exchange potential $\hat{V}_{\pi}$. Obviously the Landau-Migdal term and the correlation function

have a very similar effect on the π -exchange potential: both make the potential more repulsive and reduce the attraction in the (LO) channel.

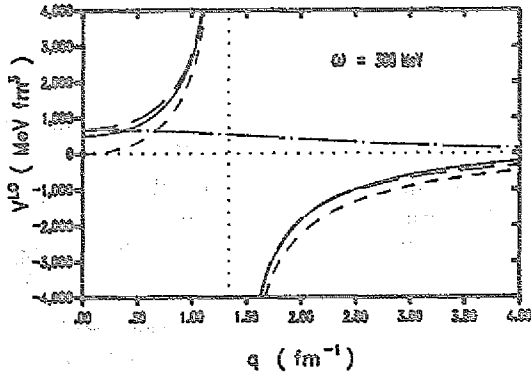


Fig. 2: Different models for the residual interaction in the (LO) channel for $\omega = 300$ MeV. Short dashed curve: π -exchange potential (eq. (1)); dash-dotted: Landau-Migdal term (eq. (2)); long dashed: sum of short dashed and dash-dotted curve; solid: "correlated" π -exchange potential (eq. (3)).

The Landau-Migdal term (2) is a central-force. Any central-force can be split into a (LO) and a (TR) part, using the following relation:

$$(\vec{S}_2^+ \cdot \vec{S}_1) = (\vec{S}_2^+ \cdot \hat{q})(\vec{S}_1 \cdot \hat{q}) + (\vec{S}_2^+ \times \hat{q}) \cdot (\vec{S}_1 \times \hat{q}) \quad (5)$$

where $\hat{q} = \vec{q} / |\vec{q}|$. Adding the Landau-Migdal term in the residual interaction generates therefore a repulsive potential in the (TR) channel. This potential is shown as solid curve in Fig. 3. Multiplication of the π -exchange potential in coordinate space by a correlation function destroys the pure (LO) spin-structure of the π -exchange and generates a contribution to the (TR) channel, which is plotted as dashed curve in Fig. 3.

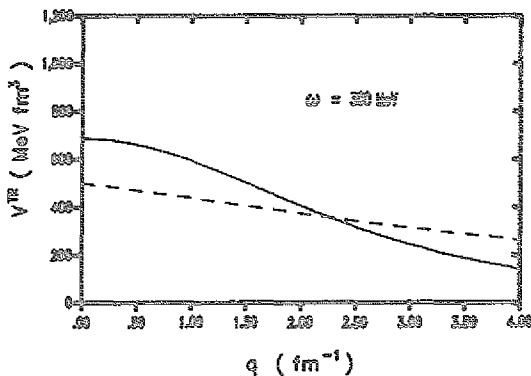


Fig. 3: Different models for the residual interaction in the (TR) channel for $\omega = 300$ MeV. Dashed curve: Landau-Migdal term (eq. (2)); solid: "correlated" π -exchange potential (eq. (3)).

In Fig. 4 we compare calculations of the total photo-absorption cross section on ^{12}C for the different models of the residual interaction. The short dashed curve shows the result without inclusion of the residual interaction. The result obtained by using a π -exchange potential alone ($\Lambda_\pi = 780$ MeV) is given as long-dashed curve and the result for a ΔN - N -interaction consisting of π -exchange potential plus Landau-Migdal term is plotted as dash-dotted curve. The solid curve in Fig. 4 shows the result for $V_{\Delta N, \Delta N} = \hat{V}_\pi$.

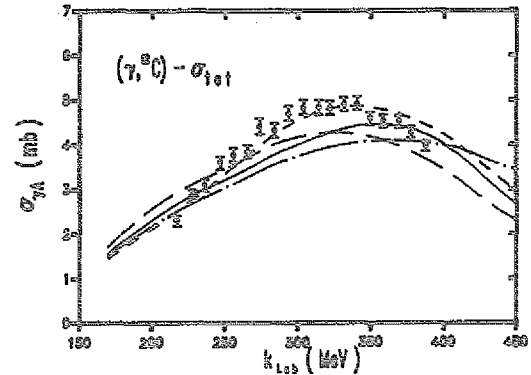


Fig. 4: The total photo-absorption cross section for ^{12}C . The data are taken from ref. [5]. Short dashed curve: calculation without inclusion of the residual interaction; long dashed: result obtained by using a π -exchange potential alone; dash-dotted: calculation with inclusion of π -exchange potential plus Landau-Migdal term as residual interaction; solid: result for the "correlated" π -exchange potential.

The results show that the residual interaction which includes the Landau-Migdal term or the correlation function is repulsive, whereas the Δ -resonance is only little shifted by a pure π -exchange potential. The reason for this is, as we have seen, that the Landau-Migdal term and the correlation function generate a repulsive potential in the (TR) channel. But we have to include the Landau-Migdal term or the correlation function in order to reproduce the data for charge-exchange reactions and pionabsorption, if we use $\Lambda_\pi = 780$ MeV [3,4] (see also the following report).

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Pionabsorption on ^{12}C in the Δ -resonance region

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In this report we present our results for the total pionabsorption cross section. The model we use has been already described for photo-absorption in report [1].

The excitation operator has been determined by analysing pion-photoproduction data on the free nucleon. The excitation operator for the absorption of a π^-, π^0, π^+ ($\alpha = -1, 0, +1$) with momentum $\vec{q}$ on a nucleon has been found to be [2,3]

$$F_{\pi N \Delta}^+ = \frac{g_{\pi N \Delta}}{M_{\pi}} v(q) (\vec{q} \cdot \vec{S}^+) T_{\alpha}^+ \quad (1)$$

where $v(q)$ is the π -formfactor

$$v(q) = \frac{1}{1+q^2/\beta^2} \quad (2)$$

with the cutoff parameter $\beta = 300$ MeV. $g_{\pi N \Delta}$ is the pion-nucleon coupling constant and $\vec{S}^+$ and T_{α}^+ are the spin- and isospin- ($1/2 \rightarrow 3/2$)-transition operators, respectively.

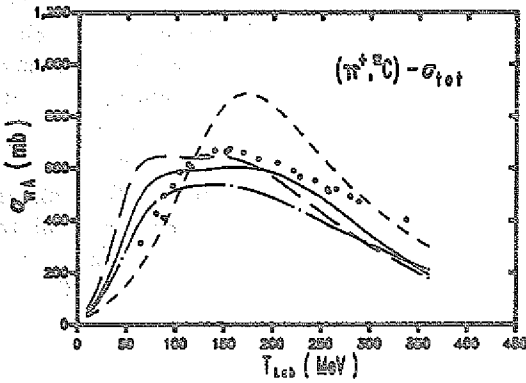


Fig. 1: The total π^+ -absorption cross section for ^{12}C . The data are taken from ref. [4]. Short dashed curve: calculation without inclusion of the residual interaction; long dashed: result obtained by using a π -exchange potential alone; dash-dotted: calculation with inclusion of π -exchange potential plus Landau-Migdal term as residual interaction; solid: result for the "correlated" π -exchange potential $\hat{V}_{\pi}$.

In Fig. 1 we compare calculations of the total pionabsorption cross section for different models of the residual interaction. The short dashed curve shows the result without inclusion of the residual interaction. The result obtained by using a π -exchange potential alone ($\Lambda_{\pi} = 780$ MeV) is given as long dashed curve and the result for a ΔN -residual interaction consisting of π -exchange potential plus Landau-Migdal term is plotted as dash-dotted curve. The solid curve in Fig. 1 shows the result for $V_{\Delta N, \Delta N} = \hat{V}_{\pi}$ (see previous report).

It is obvious, that the pure π -exchange potential with the cutoff parameter $\Lambda_{\pi} = 780$ MeV is too attractive. The Landau-Migdal term and the correlation function make the residual interaction less attractive. Then the Δ -peak position can be reproduced.

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In charge-exchange reactions at intermediate energies the Δ -peak position in nuclei is shifted downwards in energy compared to the Δ -excitation of the proton. For electromagnetic probes the experimental data do not show such a shift. Udagawa et al. [1] have analysed the charge-exchange data using the isobar-hole-model. We present in this report an approach which applies this model to the electromagnetic excitation of the Δ in nuclei.

In contrast to charge-exchange reactions the electromagnetic excitation process of the Δ is well known from studies of the pion-photoproduction and Compton scattering on the free nucleon [2,3]. There one finds that the magnetic dipole $M_{1+}(3/2)$ is the dominant excitation mechanism. We use therefore the $M_{1+}(3/2)$ -excitation operator to describe the absorption of a photon with momentum $\vec{k}$ and polarisation-vector $\vec{\epsilon}_{\vec{k}\lambda}$:

$$\tilde{F}_{\gamma N \Delta}^+ = \frac{\tilde{g}_{\gamma N \Delta}(E)}{M_{\Delta}} \vec{\epsilon}_{\vec{k}\lambda} \cdot \vec{k} \times \vec{S}^+ T_3^+ \quad (1)$$

with the effective coupling

$$\tilde{g}_{\gamma N \Delta}(E) = g_{\gamma N \Delta} e^{i\Phi(E)} \quad (2)$$

Here $\vec{S}^+$ and T_3^+ are the spin- and isospin- ($1/2 \rightarrow 3/2$) - transition operators, respectively. The coupling constant $g_{\gamma N \Delta}$ and the phase $\Phi(E)$ are fixed by comparison with the data for pion-photoproduction on the free nucleon.

We describe the excitation of the Δ in the nucleus in the impulse approximation. The total photo-absorption cross section is then given by (optical theorem):

$$\sigma_{\gamma A} = \frac{1}{2} \sum_{\lambda} \frac{1}{k} \frac{\sqrt{M_T}}{\sqrt{W}} \text{Im}[-\langle \rho | G | \rho \rangle] \quad (3)$$

where k is the photon momentum and W the total energy in the γ -nucleus-cms; M_T is the target nucleus mass. $|\rho\rangle$ is the initial doorway state

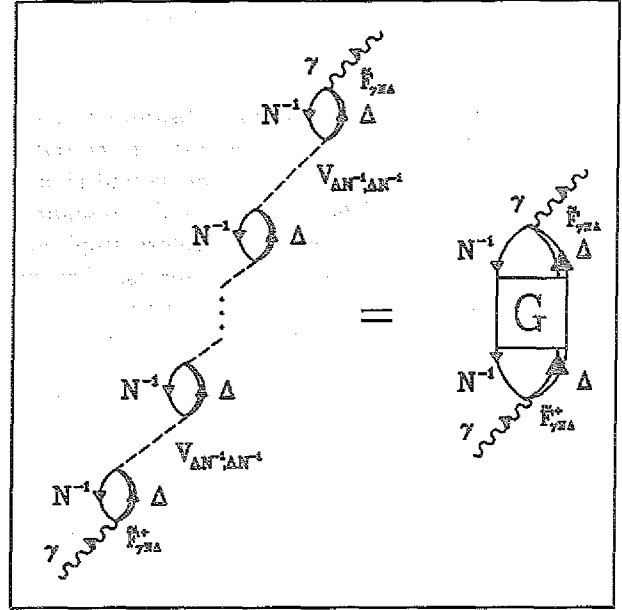


Fig. 1. Graphical representation of the Δ -hole propagation.

excited by the photonabsorption and G is the Green's function that describes the propagation of the Δ in the nuclear medium. This is schematically shown in Fig. 1. $|\rho\rangle$ has the explicit form:

$$|\rho\rangle = \tilde{F}_{\gamma N \Delta}^+ |\Phi_A\rangle \quad (4)$$

with Φ_A being the initial target ground state wave function.

For the Green's function G , we make the following ansatz [1,4]:

$$G = \frac{1}{E + i/2 \Gamma_{\Delta} - H_B - T_{\Delta} - U_{\Delta} - V_{N \Delta, N \Delta}} \quad (5)$$

In (5) E is the total energy of the Δ , Γ_{Δ} is the free decay width, T_{Δ} is the kinetic energy operator; U_{Δ} is the Δ -nucleus optical potential and H_B is the Hamiltonian of the hole-nucleus-system. $V_{N \Delta, N \Delta}$ is the ΔN^{-1} -residual interaction.

The Δ -nucleus optical potential U_{Δ} is taken to be a complex Woods-Saxon potential, consisting of the Δ -nucleus single particle potential (depth $V_{\Delta} = -65$ MeV) and a complex spreading potential (depth $W_{\text{Spread}} = +30$ MeV - i 40 MeV).

The ΔN^{-1} -residual interaction $V_{N\Delta, N\Delta}$ is a sum of a π -exchange potential and a phenomenological short range interaction term.

Fig. 2 shows the total photo-absorption cross section for ^{12}C . The dash-dotted curve is the result of the calculation without spreading potential and residual interaction. The position of the resonance is reproduced very well, but the absolute amplitude is overestimated. The reason for this is, that we have neglected all higher order absorptive processes. Including these processes by taking the spreading potential into account reduces the cross section and shifts the peak position by about 10 to 20 MeV to higher excitation energies (dashed curve). The calculation is still in good agreement with the data. The solid curve in Fig. 2 shows the result obtained by using the full propagator of eq. (5). The residual interaction shifts the resonance about 20 to 30 MeV to higher energies.

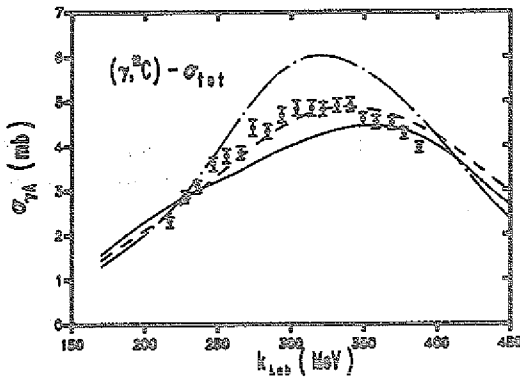


Fig. 2: The total photo-absorption cross section for ^{12}C . The data are taken from ref. [5]. Dash-dotted curve: total cross section calculated without inclusion of the spreading potential and residual interaction; dashed: total cross section calculated without inclusion of the residual interaction; solid: calculation using the full propagator of eq. (5).

The contributions of the natural parity multipoles $J^{\pi} = 1^{-}, 2^{+}$ to the total photo-absorption cross section are shown in Fig. 3. In the calculations including the residual interaction (solid and dash-dotted curves) the Δ -peak position is shifted about more than 50 MeV to higher energies compared to the calculations without residual interaction (short and long dashed curves). The multipoles $J^{\pi} = 1^{-}, 2^{+}$ are obviously responsible for the shift already observed in Fig. 2.

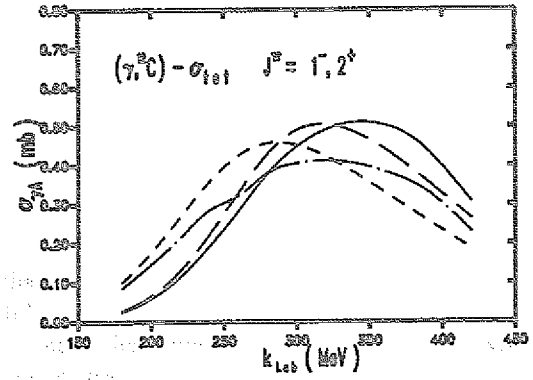


Fig. 3: Contributions of the natural parity multipoles $J^{\pi} = 1^{-}, 2^{+}$ to the total cross section. Short dashed and dash-dotted curves: contribution of the multipole $J^{\pi} = 1^{-}$ to the total cross section calculated without inclusion of the residual interaction and including the residual interaction, respectively; long dashed and solid: the same for the multipole $J^{\pi} = 2^{+}$.

This result is in contradiction to former calculations by Moniz et al. [3] and our original expectation. Because the excitation operator (1) has a spin-transversal structure ($\vec{k} \times \vec{S}^+$), whereas the π -exchange potential has a spin-longitudinal structure ($\vec{q} \cdot \vec{S}^+$), the spin-transversal excited Δ should not be affected by the spin-longitudinal residual interaction. In order to understand this result we have to study different models for the residual interaction in detail. We present such a study in the next report.

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A possible medium modification of the nucleon formfactor is one of the most intriguing but not yet fully understood phenomena [1]. No consensus has been attained as to whether the observed peculiarities in the cross section, especially the suppression of its longitudinal against transverse component, are explainable entirely in terms of conventional nucleonic degrees of freedom or if the quarks manifest their presence as well.

The MIT [2] and NIKHEF-K [3] groups report about the transverse/longitudinal relative enhancement detected in ^{12}C . The observed enhancement turns out to be stronger for the s -shell than for the p -shell. It appears impossible to explain the s -shell data within the distorted-wave impulse-approximation (DWIA) using the free nucleon formfactor [2]. Also, the meson exchange currents do not provide a significant correction for the kinematics considered [4].

Here we explore the problem of two-body rescattering effects which have been found to be important for inclusive processes, even at comparatively high momenta. In exclusive processes the coupling between protons and neutrons causes a depletion of the cross section in the $(e,e'p)$ channel [5].

This effect may also depend on the density. Intuitively one would expect that the probability for rescattering is enhanced with increasing density. One should, however, remember that the effective interaction in the relevant $(\tau\tau')$ channel becomes weaker at higher density [6]. Which of the two effects dominates?

It turned out that the rescattering effects are stronger for the $s_{1/2}$ than for the $p_{3/2}$ shell [7].

The effect is extremely important for the longitudinal cross section. The transverse cross section

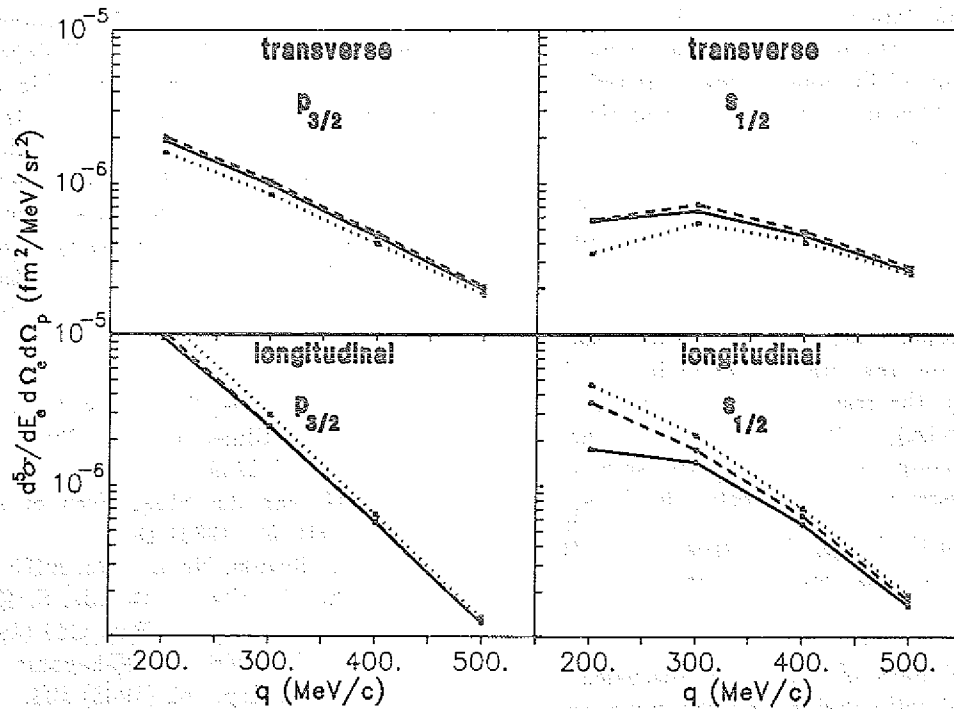


Fig.1. The separated $p_{3/2}$ (l.h.s.) and $s_{1/2}$ (r.h.s) contributions to the transverse (upper part) and longitudinal (lower part) components of the $(e,e'p)$ cross sections (in $[\text{fm}^2 \text{ MeV}/\text{sr}^2]$) in ^{12}C for parallel kinematics. The dotted lines display the mean field result, dashed lines the RPA result with the proton-neutron coupling discarded and the solid lines the full RPA result.

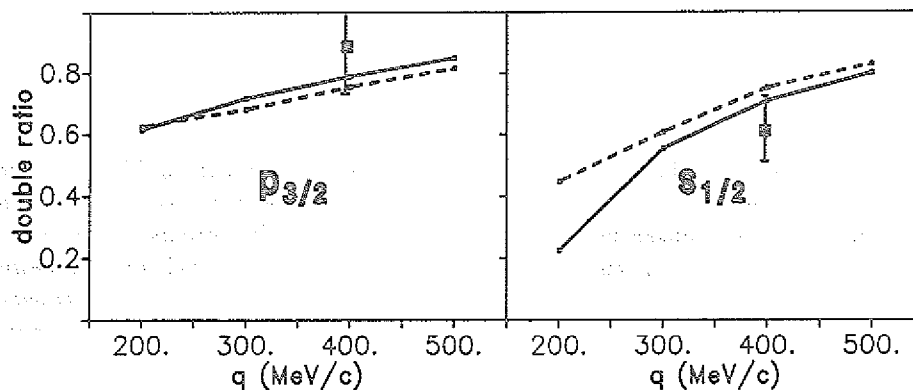


Fig.2. The RPA longitudinal/transverse ratio divided by the same ratio calculated on the mean field level. The dashed line displays the same quantity with the proton-neutron coupling in RPA ignored. The experimental data are those of ref. [2].

remains much less affected by such effects because in this channel the electron couples to the neutron already on the mean field level.

For parallel kinematics the coincidence cross section is a sum of the longitudinal and transverse components only. Their q dependence corresponding to three different calculations is illustrated in Fig. 1. The contributions for the p and s shells are separated. The bigger slope of the longitudinal part reflects the fact that at higher q -values the cross section is dominated by the spin-flip (transverse) processes. Since at the higher transferred momenta one penetrates the deeper regions of the nucleus, such a slope is smaller for the s than for the p shell both in the longitudinal and transverse cross sections. For the s shell it even bends down in the low q -region which reflects inaccessibility of this shell. More important for the present discussion is the difference between the results corresponding to the mean field (DWIA), RPA and RPA without proton-neutron coupling. As is clearly seen the rescattering processes are more effective in the $s_{1/2}$ shell. Moreover, in the longitudinal cross section the charge-exchange effects are dominant in this respect and lead to its reduction.

The experimental data of ref.[2] are discussed in terms of the L/T ratio divided by the same ratio for the distorted-wave impulse approximation. In such a double ratio the off-shell $p(e,e')p$ cross sections cancel out and the result is identical to the same ratio of the spectral functions. The analogous ratios extracted from our results are displayed in Fig. 2. At the momentum transfer which coincides with the measured values ($q \sim 400$ MeV/c) this corresponds to 0.79 for the p -shell and 0.71 for the s -shell. This is to be compared to experimental [2] $0.89 \pm 0.09 \pm 0.12$ (the same ratio extracted from

the data of ref.[3] is 0.67 ± 0.22) and $0.61 \pm 0.08 \pm 0.07$, respectively. In this kinematic region the difference of 10% in the calculation for different shells is caused entirely by the proton-neutron coupling. Without this coupling both shells give almost the same result (about 0.75).

In conclusion, the longitudinal/transverse suppression and in particular the shell dependence of this effect does not necessarily provide evidence for a medium modification of the nucleon formfactor. The observed deviations from the mean field picture largely arise from rescattering processes. Clearly, more data is needed, also for other kinematics, in order to create a more systematic picture. It would also be extremely important to obtain experimental data on the $(e,e'n)$ cross sections.

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Coupling between $(e,e'p)$ and $(e,e'n)$ channels in the quasi-elastic domain

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In recent years the $(e,e'p)$ reaction has been successfully exploited in order to understand the nuclear structure. In the quasi-free domain momentum distributions of valence nucleons have been

extracted from the data with great precision. The alternative $(e,e'n)$ channel has never been studied experimentally in the quasi-elastic region. First detailed studies of this channel have recently been

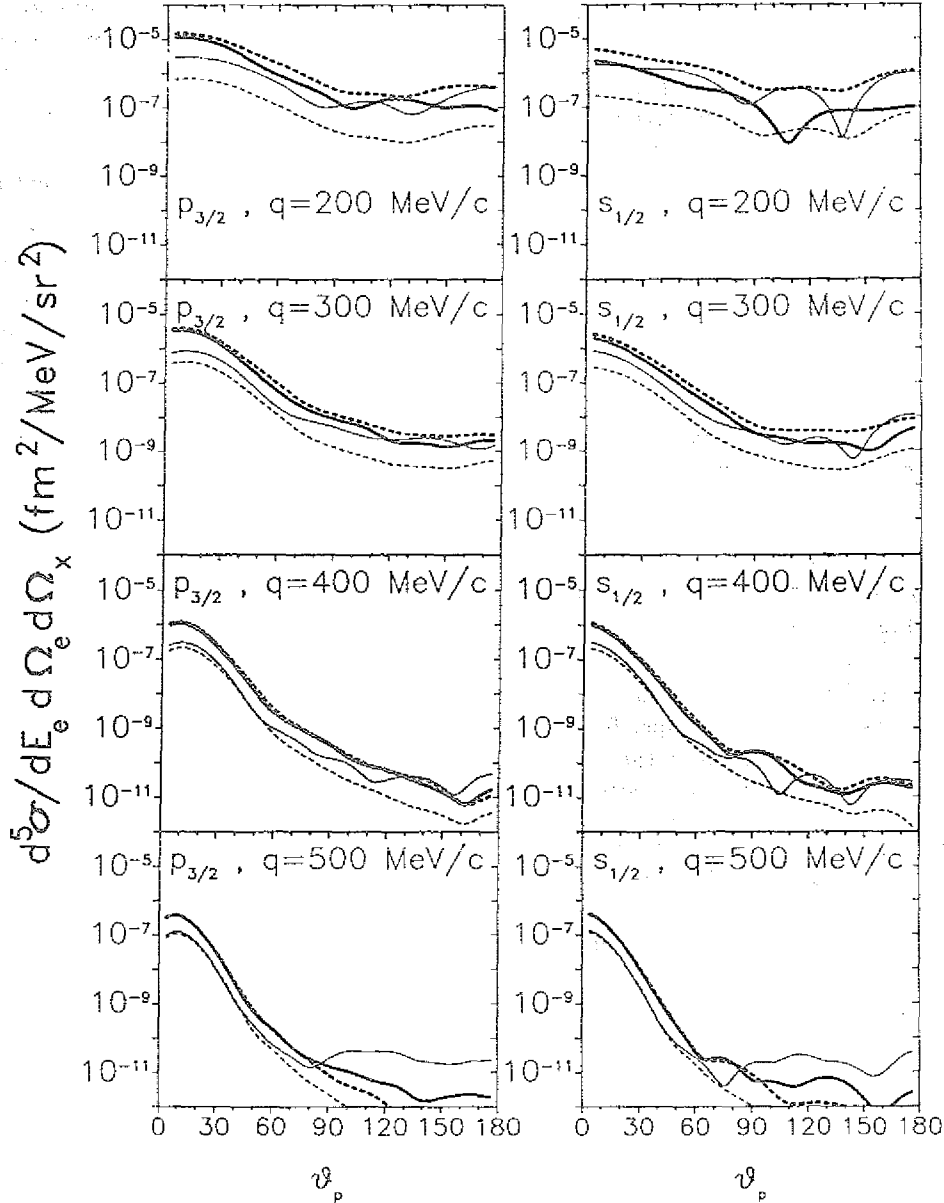


Fig. 1. The separated $p_{3/2}$ (l.h.s.) and $s_{1/2}$ (r.h.s.) contributions to the $(e,e'p)$ and $(e,e'n)$ cross sections (in $[fm^2/MeV/sr^2]$) in the quasielastic peak of ^{12}C plotted as functions of the nucleon angle measured with respect to the direction of the momentum transfer q . The thicker lines correspond to protons and the thinner ones to neutrons. In each case the dashed lines describe the mean field result and the solid lines include also the rescattering effects.

proposed by the NIKHEF-K group [1]. In the present note we quantitatively explore this problem.

The general coincidence cross section in the one-photon exchange approximation can be expressed as

$$\frac{d^6\sigma}{dE'_e d\Omega'_e dE_p d\Omega_p} = \sigma_M^p E_p (V_L R_L + V_T R_T + V_{LT} R_{LT} \cos\phi + V_{TT} R_{TT} \cos 2\phi), \quad (1)$$

where σ_M is the Mott cross section, the energies and solid angles of the final electron and proton are denoted as E'_e , E_p and Ω'_e , Ω_p respectively, and ϕ is the angle between the electron-scattering plane and the one determined by the momentum transfer q and the proton momentum p . The nuclear structure functions R involve matrix elements of the corresponding current operators between the initial and final nuclear states and the coefficients V are determined by the electron variables.

The most natural frame for studying $(e,e'x)$ processes is provided by the coupled channel method (CCM) in the basis generated by the nuclear mean field. In the present analysis we apply the recently developed method [2] to simultaneously account for the mean-field distortions and two-body rescattering processes. This method allows using finite range interactions with an explicit inclusion of exchange terms and tensor interactions. The residual interaction is specified in terms of the G -matrix [3] derived from the Bonn meson exchange model [4] in nuclear matter at densities correlated with the average densities in different shells. The mean-field part of the Hamiltonian is represented by a local Woods-Saxon potential which reproduces the ground state. The free nucleon formfactor in its standard dipole parametrization is used.

Fig. 1 shows contributions to the $(e,e'p)$ and $(e,e'n)$ coincidence cross sections coming from $p_{3/2}$ and $s_{1/2}$ shells in ^{12}C at four different q values and for the energy transfer corresponding to the position of the quasielastic peak. These cross sections are plotted versus the angle between the momentum p of the knocked-out nucleon and the transferred momentum q . Thus the small angles correspond to the parallel kinematics where the interference terms (V_{LT} and V_{TT} in Eq. 1) disappear. As is seen, the rescattering processes (the difference between CCM and pure mean field results) dramatically enhance the number of emitted neutrons at the expense of reduced proton cross section. Also, the proton angular distribution is modified due to the coupling between the $(e,e'p)$ and $(e,e'n)$ channels. Moreover, such effects are stronger for the $s_{1/2}$ than for the $p_{3/2}$ shell. The extra emission of neutrons caused by the rescattering turns out to be more probable for the s -shell, i.e. at higher densities. Consequently, measuring only protons one misses part of the cross section as compared to the mean field picture. The effect is dramatic for the longitudinal cross section. The transverse cross section remains much less affected by such effects because in this channel the electron couples to the neutron already on the mean field level. Of course, increasing q one gets closer to the mean field picture because the residual interaction becomes weaker.

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Δ -Excitations in Nuclei by the $(\vec{p}, \vec{n})$ -reaction

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In the first report we confirmed the choice of the driving interaction $t_{NN,N\Delta}$ for the description of

Δ -excitations in nuclei by analysis of the $p(\vec{p}, \vec{n})$ Δ^{++} interaction.

Using this interaction $t_{NN,N\Delta}$, charge exchange reactions in nuclei in the Δ -resonance region can be analyzed using the Distorted Waves Impulse Approximation for the nuclear reaction part and the Δ -hole model for the nuclear structure part. A detailed description of the method and of our results can be found in [1]. The analysis of the spin-response of the Δ -propagation in nuclei by the $^{12}\text{C}(\vec{d}, 2p)$ -reaction did, however, not prove the collective medium effect on the spin-longitudinal response function [2], as the spin response was obviously severely distorted by other reaction channels.

The spin-transfer coefficients of the $(\vec{p}, \vec{n})$ experiments are defined as follows:

If we denote the momentum transfer with q , the driving interaction $t_{NN,N\Delta}$ can be decomposed as

$$t_{NN,N\Delta} = t'_{N\Delta} \pi_{N\Delta} \left[\frac{\Lambda'^2 - m^2}{\Lambda'^2 - t} \right]^2$$

$$[V_{LO} \sigma_1 \cdot \hat{q} \vec{S}_2^\dagger \cdot \hat{q} + V_{TR} (\vec{\sigma}_1 \times \hat{q}) \cdot (\vec{S}_2^\dagger \times \hat{q})] \vec{T}_1 \cdot \vec{T}_2^\dagger. \quad (1)$$

Defining the spin-longitudinal response function as

$$R_L(\omega, q) = \sum_n \left| \langle n | \sum_{j=1}^A \exp[i\vec{q} \cdot \vec{r}_j] \vec{S}_j \cdot \hat{q} \vec{T}_j | 0 \rangle \right|^2 \delta(\hbar\omega - (E_n - E_0)) \quad (2)$$

and the spin-transverse response function as

$$R_T(\omega, q) = \sum_n \left| \langle n | \frac{1}{\sqrt{2}} \sum_{j=1}^A \exp[i\vec{q} \cdot \vec{r}_j] \vec{S}_j \cdot \hat{q} \vec{T}_j | 0 \rangle \right|^2 \delta(\hbar\omega - (E_n - E_0)) \quad (3)$$

where the sum over n runs over a complete set of ΔN^{-1} states with eigenenergies E_n , the spin transfer coefficients can be defined as follows:

$$C_{XX} = \frac{-|V_{LO}|^2 |R_L|^2}{|V_{LO}|^2 |R_L|^2 + 2|V_{TR}|^2 |R_T|^2} \quad (4)$$

$$C_{ZZ} = \frac{|V_{LO}|^2 |R_L|^2 - 2|V_{TR}|^2 |R_T|^2}{|V_{LO}|^2 |R_L|^2 + 2|V_{TR}|^2 |R_T|^2} \quad (5)$$

The angular dependence of the $^{12}\text{C}(p, n)$ cross sections in the Δ -resonance region is shown in figs. 1 and 2. The calculated cross sections were multiplied with an overall normalization factor of

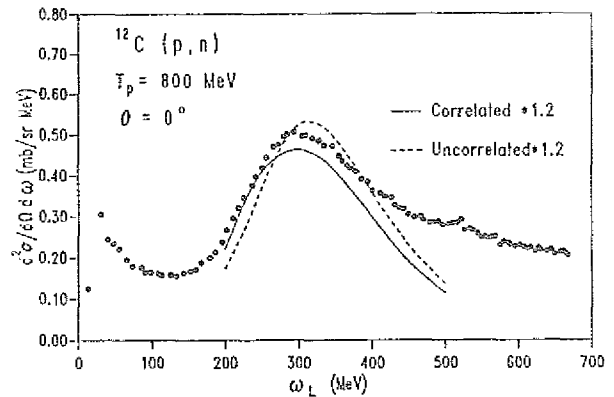


Fig. 1. Cross section of the $^{12}\text{C}(p, n)$ reaction at $E = 800$ MeV for $\theta = 0^\circ$. The solid line is the calculated cross section including Δ -hole correlations, whereas the broken line does not include them.

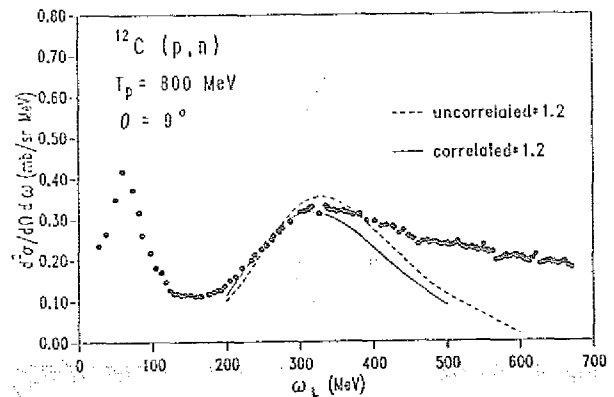


Fig. 2. As fig. 1, but for $\theta = 90^\circ$.

$N=1.2$. The full curves, which include the ΔN^{-1} correlations, are shifted by about 30 MeV due to the collective effect on the spin-longitudinal response function. It can be seen that both for a scattering angle of zero degrees and for one of nine degrees the spectra can be reproduced.

The calculations of the spin transfer coefficients are shown in figs. 3 and 4. In fig. 3, where the observable C_{XX} is shown, we can see that the Δ -hole correlations lead to a reduction of this observable relative to the uncorrelated (dashed) curve on the left side of the Δ -resonance. This reduction reflects the shift of the spin-longitudinal response function relative to the spin-transverse one. It is less apparent for higher scattering angles compared to the zero degree calculations. A corresponding result can be drawn out of the calculations of the spin transfer coefficients C_{ZZ} , which is shown in fig. 4, where we observe that the values for C_{ZZ} are shifted upwards at the left side of the Δ -resonance, and also for C_{ZZ} , this shift becomes less significant for higher scattering angles.

similar as in the $(\vec{d}, 2p)$ -reaction, contributions different from the one step Δ -excitation distort the spin observables. It seems likely though that these should have, if at all, less effect than in the $(\vec{d}, 2p)$ -reaction on the left side of the Δ -resonance, as the (p, n) reaction seems to yield a "cleaner" signal in this region.

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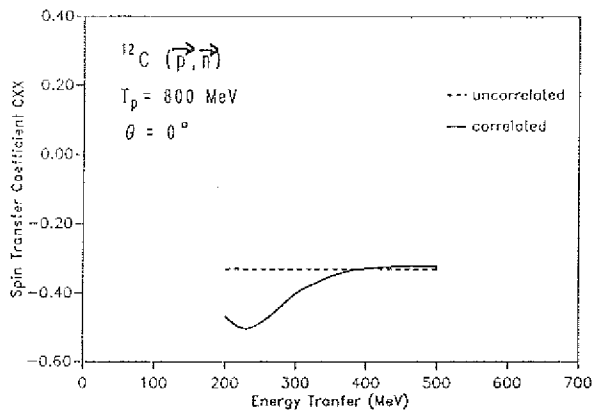


Fig. 3. As fig. 1, but spin transfer coefficient C_{XX} .

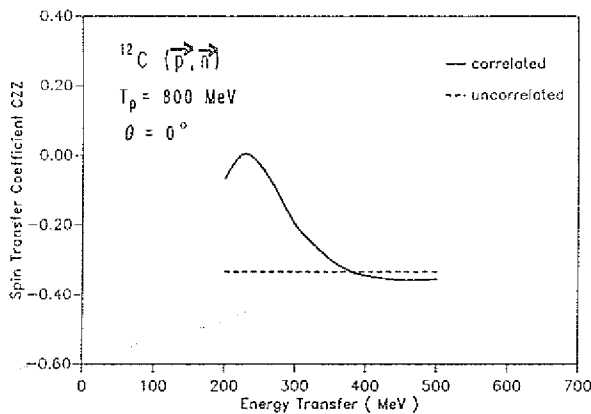


Fig. 4. As fig. 1, but spin transfer coefficient C_{ZZ} .

Therefore, at least for zero degree scattering, a large effect is expected which is caused by the Δ -hole correlations. It cannot, however, be ruled out that,

Exclusive cross sections in (h,t) reactions to the Δ -excitations in nuclei

T. Udagawa*, P. Oltmanns, F. Osterfeld

Charge exchange reactions on nuclei in the Δ -resonance region reveal a systematic downward shift which is caused by a coherent medium effect on the spin-longitudinal response function of the Δ -resonance as compared to the peak position of the proton target. This is caused by a coherent medium effect on the spin-longitudinal response function [1]. In order to obtain more evidence about this effect, data of exclusive cross sections of Δ -decay have been measured [2,3], among these are cross sections of pion plus proton and two protons which come from the decay processes $\Delta \rightarrow \pi + p$ and $\Delta + N \rightarrow N + N$, respectively, in coincidence with the outgoing triton.

The aim of the present work is to make a simple, but quantitative estimate of the coincidence cross sections based on the formalism that we used in our previous calculations of the single spectra in Ref. 1, and also to discuss some physical implications on the coincidence cross sections. Particularly, we shall point out the importance of the exclusive coherent pion production cross section which carries information on the Δh -correlations.

Let us first briefly discuss our coincidence cross section formulae, which can be derived from the inclusive (singles) cross section formula used in our previous work in Ref. 1. It reads as (we denote the $(^3\text{He}, t)$ reaction as $A + a \rightarrow B + \Delta + b$, where a (b) is the projectile (ejectile), A is the target and B is the residual hole nucleus)

$$\frac{d^2\sigma}{dE_b d\Omega_b} = \frac{E_a E_b E_A E_{B+\Delta}}{(2\pi\hbar^2 c^2 W)^2} \frac{k_b}{k_a} \Sigma_{ljm}$$

$$\begin{aligned} & [\langle \psi_{ljm} | \frac{\Gamma_{\Delta}}{2} | \psi_{ljm} \rangle + \langle \psi_{ljm} | -W_{\Delta} | \psi_{ljm} \rangle \\ & + \langle \psi_{ljm} | -\text{Im}[V_{N\Delta, N\Delta} | \psi_{ljm} \rangle] , \end{aligned} \quad (1)$$

where E_i is the total energy of particle i ($i=a, b, A$, and $B+\Delta$), $k_a(k_b)$ is the wave number of a (b), and W is the total energy of all particles in the center of momentum system. $|\psi_{ljm}\rangle = G|\rho_{ljm}\rangle$ is the Δh wave function of the system with definite orbital angular momentum l , total angular momentum j and the projection m . ($|\rho_{ljm}\rangle$ is the corresponding doorway state.) G is the Green's function and describes the propagation of the $B+\Delta$ system. The detailed expression for ρ and G are given in

Ref. 1. The three terms in the above equation describe contributions from decay of the Δ into $\pi + p$, the spreading of the ΔN^{-1} state into $2p2h$ states (conversion of the Δ into the nucleon), and the coherent pion production leaving the target in its ground state. The $\pi + p$ and $2p$ production cross sections are contained in the first and second terms, respectively. These terms may also contain contributions from the decay processes followed by absorption of one or both of the decay products. Therefore, in order to obtain the real $\pi + p$ and $2p$ production cross sections, we have to decompose the first and second terms in Eq. (4) into the emission and absorption components. This decomposition may be carried out if one assumes Feshbach type formulae for Γ_{Δ} and $W_{\Delta}(r)$ with an optical model intermediate Green's function, and then following the prescription developed for deriving the breakup-fusion formula [4].

The essential idea of the method is to replace Γ and W_{Δ} in eq. (1) by microscopic operators that express explicitly the decay modes. The expression for W_{Δ} is the well known Feshbach formula, while Γ may be given in a very similar form as

$$\Gamma_{\Delta} = -2\text{Im}[V_{DP} \frac{1}{E - H_{PP} + i\epsilon} V_{PD}] \quad (2)$$

where V_{PD} (V_{DP}) describes the interaction of the $\Delta \rightarrow \pi + N$ ($\pi + N \rightarrow \Delta$) process, while H_{PP} is the one pion plus residual nucleus Hamiltonian, which may be given as

$$H_{PP} = T_{\pi} U_{\pi} + T_N U_N + H_h \quad (3)$$

In the above equation, T_i and U_i ($i=\pi$ and N) are the kinetic energy operator and the potential energy of the decay product i . The potential U_i describes the nuclear medium effects on the particle i , which we take to be complex. The imaginary part of U then takes into account the absorption of the particle i while it escapes from the nucleus. The decomposition of Γ_{Δ} yields

$$\frac{1}{2} \hat{\Gamma}_{\Delta}^{\dagger} = \pi V_{DP} \omega_P^{(-)} \omega_{\pi}^{(-)} \delta(\omega - T_{\pi} - T_P) \omega_P^{(-)\dagger} \omega_{\pi}^{(-)\dagger} V_{PD}^{\dagger} \quad (4)$$

where the wave operators $\omega_P^{(-)}$ and $\omega_{\pi}^{(-)}$ in Γ_{Δ} describe the emission of both p and π . A similar decomposition can be done for W_{Δ} .

In the present work, we shall not follow such a prescription, but rather make the following simple ansatz that the resulting emission parts of Γ_{Δ} and

W_{Δ} are simple local operators $\Gamma_{\Delta}^{\dagger}(r)$ and $W_{\Delta}^{\dagger}(r)$, respectively, having nonvanishing values only outside a certain cutoff radius R_i ($i = \pi p$ and $2p$). We further assume that in the outer region $\Gamma_{\Delta}^{\dagger}(r)$ and $W_{\Delta}^{\dagger}(r)$ are equal to the free Γ_{Δ} -value and the original spreading potential $W_P(r)$, respectively. We thus have

$$\Gamma_{\Delta}^{\dagger}(r) = \Gamma_{\Delta} \Theta(r - R_{\pi p}) \quad (5)$$

and

$$W_{\Delta}^{\dagger}(r) = W_{\Delta} \Theta(r - R_{2p}) \quad (6)$$

We may then treat R_i ($i = \pi p$ or $2p$) as a phenomenological parameter that can be determined from fit to the experimental cross section.

Fig. 1 shows the fit obtained. The full lines are theoretical predictions assuming $R_{\pi p} = 3.6$ fm and $R_{2p} = 2.1$ fm, while the experimental data are shown by the filled circles. It is remarkable that the theoretical predictions fit not only the magnitude but also the shape of the spectra fairly well. Particularly, the observed difference between the peaks in the $\pi+p$ and $2p$ coincidence spectra is reproduced well in the calculations. In Fig. 1, we presented by broken lines $1/8$ and $1/2$ of the total $\pi+p$ and $2p$ cross sections, respectively, both obtained by setting $R_i = 0$, i.e., without introducing any cutoff. The fact that the magnitude of the two cross sections (full and broken lines) agree for each other indicates that the reduction factors due to the absorption is 8 for the $\pi+p$ and 2 for the $2p$ case. The large reduction factor (equivalently the large cutoff radius value of 3.6 fm) for the $\pi+p$ case is due to the large absorption of π . Note that for the $2p$ emission case, both protons may carry approximately 150 MeV kinetic energy. At that energy, the absorption becomes small. The rather small absorp-

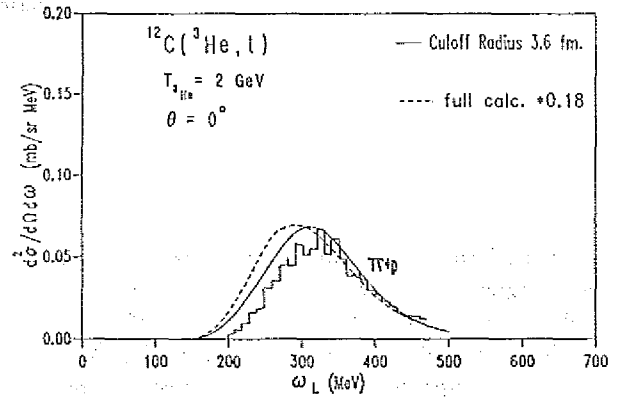


Fig. 1. πp coincidence cross section of the $^{12}\text{C}(^3\text{He}, t)$ reaction at $E = 2$ GeV for $\theta = 0^\circ$. The solid line is the calculated cross section with $R_{\pi p} = 3.6$ fm while the broken is that with $R_{\pi p} = 0$ and $N_{\pi p} = 0.18$.

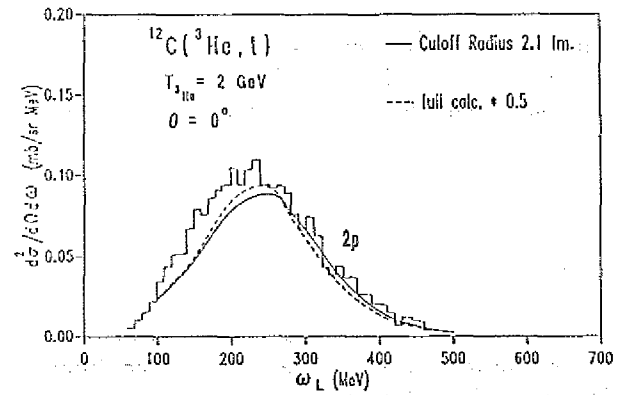


Fig. 2. $2p$ coincidence cross sections of the same reactions as considered in fig. 1. The solid line is the calculated cross section with $R_{2p} = 2.1$ fm while the broken line is that with $R_{2p} = 0$ and $N_{2p} = 0.5$.

tion factor of 2.1 or equivalently the small cutoff radius of $R_{2p} = 2.1$ fm may reflect this.

It is important to note that the peak position of the $\pi+p$ cross section obtained with $R_{\pi p} = 3.6$ fm (full line) is located at a higher energy as compared with the peak of the cross section obtained without introducing any cutoff (broken line). This result may be understood as a result from the fact that the downward shift due to the Δh correction is more important for the Δ created in the nuclear interior rather than the surface. In these calculations, the limited acceptance of the Diogene detector has not been taken into account. Calculations of Gaarde et al. [5] indicate, however, that only the magnitude of the cross sections, but not the shape of the cross sections is different from those of an

Rescattering contribution to proton—proton bremsstrahlung in a relativistic approach

V. Herrmann and K. Nakayama*

Interest in the nucleon—nucleon bremsstrahlung (NN γ) reaction as a means of testing the off—energy—shell behavior of the NN interaction has been rejuvenated considerably in the past few years. In fact, it has been demonstrated for the first time by the TRIUMF group [1,2] that NN γ reactions, or more specifically, the analyzing power in the proton—proton bremsstrahlung (pp γ) process at an incident energy near the pion production threshold shows a clear signature of off—shell effects. However, a rather disturbing disagreement has been observed between recent pp γ section data [3] at an incident energy of $T_{lab} = 280$ MeV and calculations [1,4] using a modern potential model. In order to achieve reasonable agreement, the data of Ref. 3 have been scaled down by a factor of 2/3; however, no obvious reason for an error in the absolute normalization has been found. A large part of the disagreement between the calculated and the non—normalized data may be attributed to the relativistic correction (RSC) which reduces the calculated cross section by about 20% — 30% with respect to the non—relativistic results. If the normalization of the TRIUMF's data is correct, the potential model may be unable to describe NN γ processes. The calculation accompanying the experiment reported in Ref. 3, however, does not include the rescattering (RES) contribution which, in the non—relativistic approach, enhances the cross section by about 25%. Given the comparable sizes of the RSC and RES contributions, it is, therefore, crucial that both corrections are included simultaneously for a proper assessment of the current status of agreement between the theory and experiment.

We investigate here the role of the RES contribution in the pp γ reaction within a relativistic approach and compare with the recent data by the TRIUMF group [3]. This is the first NN γ calculation which includes, simultaneously, both the RSC and relativistic (one—body) RES corrections.

The one—body RES bremsstrahlung amplitude is given by

$$R = \langle \vec{\epsilon}, \vec{k}; \phi_f | (T^-)^{\dagger} G_f V_{em} G_i T^+ | 0; \phi_i \rangle, \quad (1)$$

where ϕ denotes the two—nucleon unperturbed wave function, G is the two—nucleon Greens

function consistent with the approximation used for solving the relativistic T —matrix equation, and $T^{\pm}$ denotes the NN T —matrix. The subscript $i(f)$ refers to the initial (final) two—nucleon state, while the superscript $+$ or $-$ indicates the boundary condition associated with the outgoing (+) or incoming (—) waves. $\vec{k}$ and $\vec{\epsilon}$ stand for the momentum and polarization of the emitted photon, respectively.

V_{em} in eq.(1) denotes the relativistic electromagnetic transition potential to be used with two—component wave functions. It is obtained from the relativistic nuclear Hamiltonian via the minimal substitution, $\vec{p} \rightarrow \vec{p} - e\vec{A}$, and by considering the electromagnetic coupling to nucleon only to first order. We write V_{em} as the sum of four terms

$$V_{em} = V_{conv} + V_{mag} + V_{rsc} + V_{rem}. \quad (2)$$

In eq.(2), the dominant RSC is given by V_{rsc} . The last term in eq.(2), V_{rem} , contains the remaining relativistic corrections. It has a completely different operator structure than the first three terms and its contribution is negligible compared to those from the other terms.

In Fig.1 we illustrate the effect of the one—body RES term in the coplanar geometry pp γ cross section as a function of symmetric proton scattering angles $\theta_1 = \theta_2$ at an incident energy of $T_{lab} = 280$ MeV and at a fixed photon emission angle of $\theta = 175^\circ$. As can be seen, in the relativistic approach, the RES contribution (lower solid curve) enhances the cross section by $\sim 20\%$ with respect to the relativistic result without the RES term (lower dashed curve) for proton angles $\theta_1 = \theta_2 > 10^\circ$. However, for proton angles $\theta_1 = \theta_2 < 10^\circ$, the RES contribution reduces the cross section; moreover, its effect is more pronounced than at larger proton scattering angles, especially for $\theta_1 = \theta_2 \sim 5^\circ$ where the cross section is reduced by $\sim 2\mu b/(sr^2 \cdot rad)$ (or $\sim 28\%$). This is due, mainly, to the strong 1S_0 state contribution from the RSC term V_{rsc} (see eq.(1)). For very small proton scattering angles the emitted photon carries most of the available kinetic energy and, consequently, the energy in the c.m. of the two scattered protons is very small. The NN γ process, then, necessarily samples the 1S_0 state which is quite strong for such small energies. In

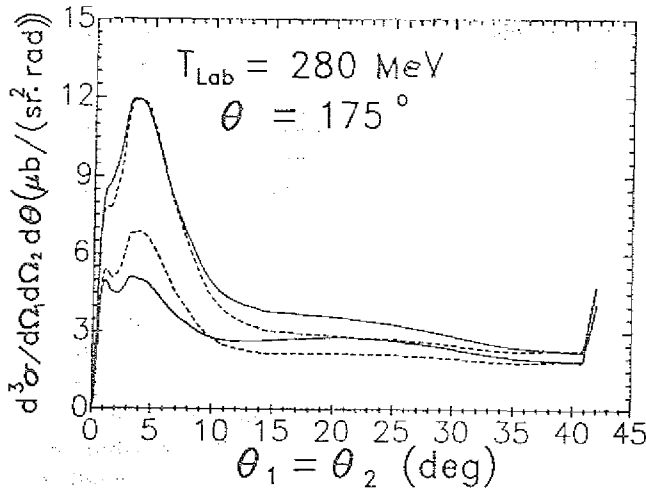


Fig. 1. Coplanar geometry $pp\gamma$ cross sections in the laboratory frame as a function of symmetric proton scattering angles $\theta_1 = \theta_2$ at an incident energy of $T_{\text{Lab}} = 280$ MeV and at a fixed photon emission angle of $\theta = 175^\circ$. The lower (upper) two curves correspond to the relativistic (non-relativistic) calculation with (solid lines) and without (dashed lines) the one-body RES contribution.

Fig. 1 we also display the results with (upper solid curve) and without (upper dashed curve) the RES contribution in the non-relativistic approach. Here, the larger enhancement of the cross section (due to the RES term) compared to the relativistic calculation in the region of $\theta_1 = \theta_2 > 10^\circ$ is due to the fact that in the relativistic case there is destructive interference between V_{mag} and V_{rsc} , while in the non-relativistic case the RSC term V_{rsc} is simply absent. For proton scattering angles $\theta_1 = \theta_2 < 10^\circ$, the RES contribution in the non-relativistic case is relatively small, in contrast to the relativistic case. From Fig.1 it is clear that this is due to the absence of V_{rsc} in the non-relativistic approach. We emphasize that relativistic corrections, other than that due to V_{rsc} , are negligible for the kinematical conditions considered here. It is interesting to note that, for proton scattering angles between $\sim 15^\circ$ and $\sim 35^\circ$ where most of the calculations have been made in the past, the relativistic calculation with the RES contribution yields practically the same results as the much simpler non-relativistic calculation without the RES term. At proton angles around

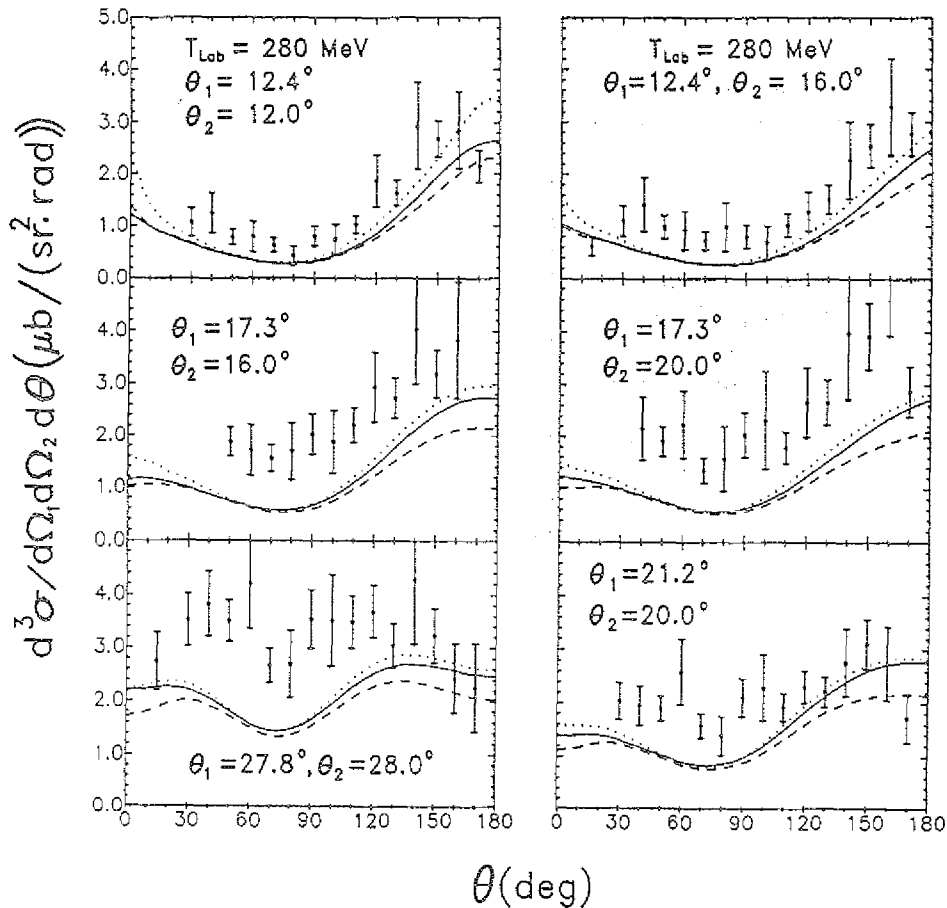


Fig. 2. Coplanar geometry $pp\gamma$ cross sections in the laboratory frame as a function of photon emission angle θ at an incident energy of $T_{\text{Lab}} = 280$ MeV for various asymmetric proton scattering angles θ_1 and θ_2 . The solid and dashed curves correspond to the relativistic calculation with and without the one-body RES contribution, respectively, while the dotted curves correspond to the non-relativistic calculation without the RES term. The data are from Ref.3 but do not contain the arbitrary normalization factor of 2/3.

$\theta_1 = \theta_2 = 5^\circ$, however, inclusion of the RES term further enhances the relativistic effects. An understanding of the effects of different terms on $NN\gamma$ reactions at these small angles, in particular the effect of the RES, is especially relevant for experiments currently being performed at IUCF [5].

Fig.2 shows a comparison of the present cross section calculations and the data of Ref.3 in the coplanar geometry as a function of photon emission angle θ at an incident energy of $T_{lab} = 280$ MeV and for various asymmetric proton scattering angles. The solid and dashed curves correspond to the relativistic calculations with and without the RES contribution, respectively. The non-relativistic results without the RES term are also shown (dotted curves). As we expect from the discussion of Fig. 1, the RES term enhances the cross section for these proton scattering angles and for $\theta_1 \sim \theta_2$ reduces the discrepancy between the theory and experiment significantly. However, the enhancement is not large enough to remove the discrepancy completely. As the proton angles increase, the relativistic results with the RES contribution become very close to the non-relativistic results without the RES term as discussed before.

In summary, we have investigated the role of the one-body RES contribution to the $pp\gamma$ reaction within the relativistic approach. A significant difference in its role relative to that in the non-relativistic approach arises for kinematical conditions where the photon energy is close to its endpoint energy. Under these conditions the 1S_0 state contribution from V_{RSC} becomes very important; we note that V_{RSC} is absent in the non-relativistic approach. For these kinematical conditions the inclusion of the RES term in the relativistic approach further amplifies the relativistic effects on the $pp\gamma$ cross section. It would be very interesting to have cross section data under these kinematical conditions where potential model

calculations within the relativistic approach can be better tested. In this connection, $pp\gamma$ experiments being performed at IUCF [5] for very small proton scattering angles are of special interest. For the coplanar geometry experiment of the TRIUMF group [3] at $T_{lab} = 280$ MeV and proton scattering angles larger than 10° , the inclusion of the RES term enhances the cross section. Although this enhancement significantly reduces the difference observed between the data and theoretical calculations which only include the external current in the relativistic approach, the RES correction is not large enough to completely resolve the discrepancy. The RES contribution to the analyzing power (not shown) is not negligible; however, the data are not sufficiently accurate to disentangle such an effect. Moreover, the analyzing power is less affected by relativistic corrections than the corresponding cross section since this spin observable is not very sensitive to the 1S_0 state where the dominant (and large) RSC occurs in the cross section.

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Spin observables in pp bremsstrahlung and off-shell effects of the NN interaction

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The nucleon-nucleon (NN) bremsstrahlung reaction has been studied intensively, especially in the sixties and seventies, in order to learn about the off-shell behavior of the NN interaction. However, most of those studies have been restricted to investigating NN γ cross sections and almost no attention has been given to the investigation of spin polarization observables. Due to the weakness of the electromagnetic interaction combined with the experimental techniques which were available, accurate measurements of spin polarization observables in the NN bremsstrahlung process were extremely difficult to perform in the past. More recently, the TRIUMF group [1-4] has measured the analyzing power for the proton-proton (pp) bremsstrahlung process and has demonstrated that it is very sensitive to the off-shell behavior of the NN interaction. In particular, it has been shown that the analyzing power probes primarily the off-shell behavior of the NN tensor force [5]. Motivated by these findings, we investigate in the present work the sensitivity of the spin-correlation coefficients to off-shell effects in the pp bremsstrahlung reaction. These spin observables have not been considered previously in this type of reaction and, therefore, it is of interest to see to what extent they provide a signature of similar or complimentary off-shell effects, especially in view of the sensitivity to off-shell effects shown by the analyzing power.

The spin-correlation coefficients C_{ij} which provide a measure of the sensitivity of the bremsstrahlung reaction to the initial spin orientations of the incoming particle 1 and the target particle 2 are defined as

$$C_{ij} = \frac{\text{Tr}(\tilde{M} \vec{\sigma}_1 \cdot \hat{n}_i \vec{\sigma}_2 \cdot \hat{n}_j \tilde{M}^\dagger)}{\text{Tr}(\tilde{M} \tilde{M}^\dagger)} \quad (1)$$

where $\vec{\sigma}_1$ and $\vec{\sigma}_2$ stand for the Pauli spin matrices; $\hat{n}_i$ and $\hat{n}_j$ denote unit vectors along the x-, y- and z-axes. In the following we will restrict our considerations to the cases where $\hat{n}_i = \hat{n}_j$, i.e., C_{ii} . The spin-correlation coefficient C_{ii} is given by

$$C_{ii} = \frac{[\sigma_i(++) + \sigma_i(--)] - [\sigma_i(+-) + \sigma_i(-+)]}{[\sigma_i(++) + \sigma_i(--)] + [\sigma_i(+-) + \sigma_i(-+)]} \quad (2)$$

with σ_i denoting the cross section with the spins of the beam and target nucleons along the direction $\hat{n}_i$. The arguments of σ_i indicate whether the spins of the beam (first argument) and target (second argument) nucleons are parallel (+) or antiparallel (-) to $\hat{n}_i$.

In the present work, the pp bremsstrahlung transition amplitude $\tilde{M}$ in eq.(1) is calculated in momentum space within the framework of the relativistic one-boson-exchange potential model developed by the Bonn group [6]. The present calculation treats the coupling of the photon to the nucleon relativistically to all orders in photon momentum and thereby includes the relativistic spin correction which is known to reduce the pp bremsstrahlung cross section compared to the non-relativistic calculation [1,7]. The details of the formalism will be given elsewhere. In the following, the off-shell result means that the off energy-shell T-matrix elements are used in the calculation as prescribed by the bremsstrahlung reaction, while the on-shell result means that the T-matrix elements, and only the T-matrix elements, are forced to be on the energy-shell according to eq.(2) of Ref. 5 which is equivalent to setting the half-off-shell extension function [2] to unity, its on-shell value.

Fig. 1 illustrates off-shell effects in the pp γ spin-correlation coefficients C_{xx} , C_{yy} and C_{zz} in the coplanar geometry at an incident energy of $T_{\text{lab}} = 280\text{MeV}$ for symmetric proton scattering angles of $\theta_1 = \theta_2 = 10^\circ$ and 20° . For comparison, in the upper part of Fig. 1 we also display the corresponding analyzing power results. As one can see, for forward proton scattering angles of $\theta_1 = \theta_2 = 10^\circ$ (column a), C_{xx} exhibits a degree of sensitivity to off-shell effects which is larger than that observed in the corresponding analyzing power. For the particular kinematical condition of Fig. 1a, the spin-correlation coefficients C_{yy} and C_{zz} are less sensitive to off-shell effects than C_{xx} . At larger proton scattering angles of $\theta_1 = \theta_2 = 20^\circ$ (column b), the spin-correlation coefficients show a similar or even stronger sensitivity to off-shell effects than that observed at smaller scattering angles. In contrast, the analyzing power becomes much less sensitive to off-shell effects at these larger proton scattering angles. We note that the insensitivity of C_{yy} to the off-shell effects shown in

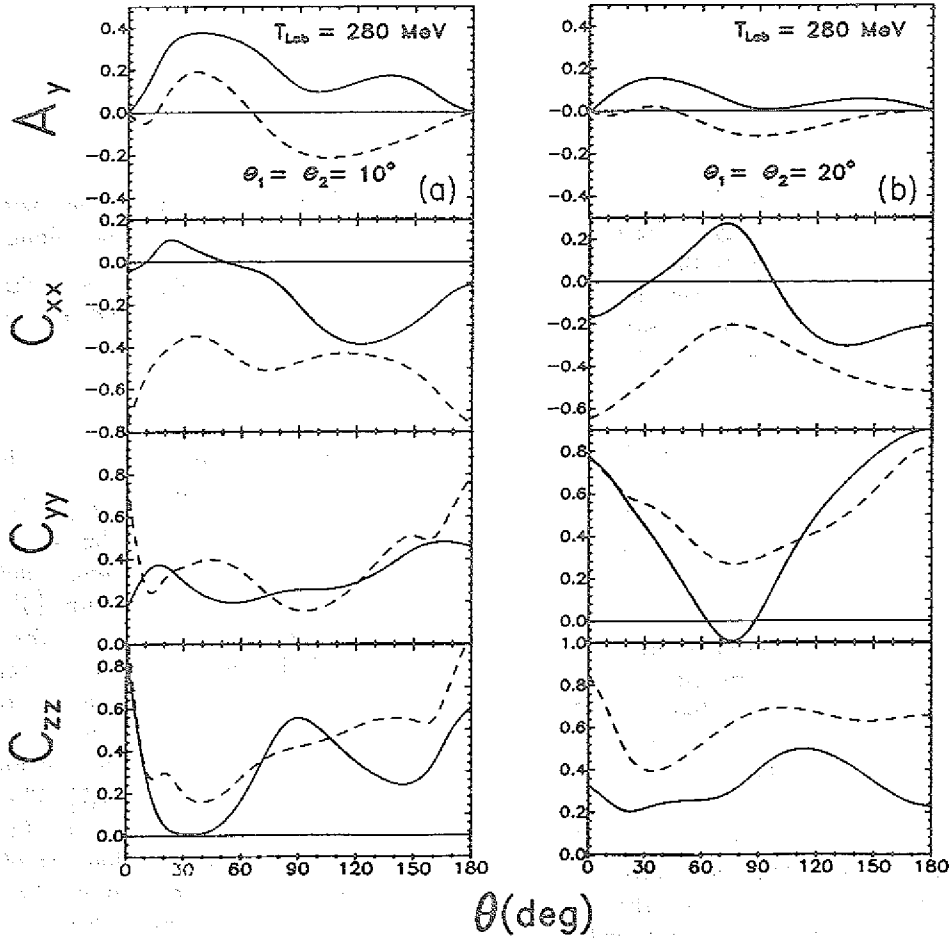


Fig. 1. Coplanar geometry $pp\gamma$ spin polarization observables in the laboratory frame as a function of photon emission angle θ . From the top to bottom are the observables analyzing power A_y , the spin correlation coefficients C_{xx} , C_{yy} and C_{zz} . Column a shows the results for symmetric proton scattering angles of $\theta_1 = \theta_2 = 10^\circ$, while column b shows the results for $\theta_1 = \theta_2 = 20^\circ$. The incident energy is $T_{lab} = 280\text{MeV}$. The solid and dashed curves denote the off- and on-shell results, respectively.

Fig. 1 at $\theta_1 = \theta_2 = 10^\circ$ is peculiar to the particular T-matrix interaction used in the present work. For this T-matrix the overall insensitivity at 10° results from destructive interference between appreciable off-shell effects from various partial-wave states. Indeed, using an alternative Bonn T-matrix, we find that C_{yy} is sensitive to off-shell effects when all of the important partial-wave states are included.

In summary, we have tested the sensitivity of the spin-correlation coefficients in the $pp\gamma$ reaction to the off-shell behavior of the NN interaction. At proton scattering angles of $\theta_1 = \theta_2 = 10^\circ$, we have found that C_{xx} is more sensitive to off-shell effects than the corresponding analyzing power. At larger proton scattering angles where measurements should be less difficult, this sensitivity is similar to or even enhanced (in the case of C_{yy} and C_{zz}) with respect to that observed at $\theta_1 = \theta_2 = 10^\circ$. In contrast, the analyzing power becomes much less sensitive to off-shell effects as the proton scattering

angles increase. Overall, the spin-correlation coefficients sample predominantly the off-shell behavior of the NN tensor force as does the analyzing power. Although, the measurements of spin-correlation coefficients are more difficult than those of the analyzing power, they would appear to offer an alternative means of investigating off-shell effects of the NN interaction, especially at large proton scattering angles where, experimentally, problems with the background should be smaller than at small proton scattering angles. Moreover, in the coplanar geometry, results at large proton scattering angles are also less subject to theoretical uncertainties than are at small proton scattering angles, because, the energy carried by the emitted photon decreases with increasing proton scattering angles; consequently, effects such as the Coulomb correction and the two-body current contribution (which cannot be incorporated fully into the calculation at the present) are expected to be much less important.

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On the relativistic spin correction in proton—proton bremsstrahlung

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Relativistic effects in proton—proton (pp) bremsstrahlung reactions have been investigated by a number of authors in the past and are known to reduce the pp bremsstrahlung cross section with respect to non—relativistic results. The dominant relativistic correction arises from the so called relativistic spin correction (RSC). However, to our knowledge, only the work by Liou and Sobel [1] examines the physics behind this correction to pp bremsstrahlung process. They analyze in terms of kinematical factors involved but no complete explanation for the behavior and magnitude of the RSC has been given.

In the present work, we investigate relativistic effects, and in particular the RSC, in an effort to better understand its role in pp bremsstrahlung reactions. We demonstrate that, for energetic photons, the effect is largely dynamical in origin and arises predominantly from the strong energy dependence of the 1S_0 state and that this can be understood in a very simple way.

The relativistic electromagnetic transition potential (to be used with two component wave functions) is obtained from the relativistic nuclear Hamiltonian via the minimal substitution, $\vec{p} \rightarrow \vec{p} - e\vec{A}$, and by considering the electromagnetic coupling to nucleon only to first order. We write the coupling as the sum of four terms

$$V_{em} = V_{con} + V_{mag} + V_{RSC} + V_{rem}, \quad (1)$$

where, in the NN center—of—mass (c.m.) frame,

$$V_{con} = -\sqrt{\frac{2\pi}{k}} \frac{e}{2m} \vec{\epsilon} \cdot (\vec{p} + \vec{p}') \times \left\{ \tilde{\epsilon}_1^{\pm} \delta(\vec{p} - \vec{p}' - \frac{\vec{k}}{2}) - \tilde{\epsilon}_2^{\pm} \delta(\vec{p} - \vec{p}' + \frac{\vec{k}}{2}) \right\} \quad (2)$$

$$V_{mag} = -i\sqrt{\frac{2\pi}{k}} \frac{e}{2m} \times \left\{ \tilde{\mu}_1^{\pm} \vec{\epsilon} \cdot (\vec{k} \wedge \vec{\sigma}_1) \delta(\vec{p} - \vec{p}' - \frac{\vec{k}}{2}) + \tilde{\mu}_2^{\pm} \vec{\epsilon} \cdot (\vec{k} \wedge \vec{\sigma}_2) \delta(\vec{p} - \vec{p}' + \frac{\vec{k}}{2}) \right\} \quad (3)$$

and

$$V_{RSC} = i\sqrt{\frac{2\pi}{k}} \frac{e}{2m} \times \left\{ \tilde{\nu}_1^{\pm} \vec{\epsilon} \cdot (\vec{p}' \wedge \vec{\sigma}_1) \delta(\vec{p} - \vec{p}' - \frac{\vec{k}}{2}) - \tilde{\nu}_2^{\pm} \vec{\epsilon} \cdot (\vec{p}' \wedge \vec{\sigma}_2) \delta(\vec{p} - \vec{p}' + \frac{\vec{k}}{2}) \right\}. \quad (4)$$

In the above equations, $\vec{p}$ and $\vec{p}'$ denote the relative momenta of the two interacting nucleons before and after the emission of a photon with momentum $\vec{k}$ and polarization $\vec{\epsilon}$. The nucleon mass and electric charge are denoted by m and e , respectively, and $\vec{\sigma}_i$ stands for the Pauli spin matrix for nucleon i ($= 1, 2$). The factors $\tilde{\epsilon}_i^{\pm}$, $\tilde{\mu}_i^{\pm}$ and $\tilde{\nu}_i^{\pm}$ are functions of nucleon and photon momenta. The leading terms in $\tilde{\epsilon}_i^{\pm}$ and $\tilde{\mu}_i^{\pm}$ are given by $\tilde{\epsilon}_i^{\pm} = 1$ and $\tilde{\mu}_i^{\pm} = \mu_i$ which give rise to the convection and magnetization current operators of the conventional non—relativistic approach. Here, μ_i denotes the anomalous nuclear magnetic moment in units of nuclear magnetons (for protons $\mu_1 = \mu_2 = \mu = 2.793$). If we consider terms through p/m in the factors $\tilde{\epsilon}_i^{\pm}$, $\tilde{\mu}_i^{\pm}$ and $\tilde{\nu}_i^{\pm}$, one recovers the usual expression for the interaction Hamiltonian obtained via the Foldy—Wouthuysen reduction when terms through p/m are kept.

In eq.(1), the dominant RSC is denoted by V_{RSC} . It results from taking the matrix element of the relativistic four component electromagnetic transition operator between the upper and lower components of the Dirac spinors. The remaining relativistic correction, V_{rem} in eq.(1), has a different operator structure than the first three terms and its contribution is negligible compared to those from the other terms.

Before discussing the major results of the present work, we note that a simple comparison of eq.(3) and eq.(4), together with the fact that $\tilde{\mu}_1^{\pm} = \tilde{\mu}_2^{\pm} = \mu$ and $\tilde{\nu}_1^{\pm} = \tilde{\nu}_2^{\pm} = (\mu - 1/2)k/m$ to leading orders, reveals that it is not possible to analyze the effect of the RSC in terms of an overall reduction factor of the cross section such as that discussed in an earlier work [1]. We also note that while V_{mag} in eq.(3) contains an overall factor of $k = |\vec{k}|$, V_{RSC}

in eq.(4) contains an overall factor of $p' = |\vec{p}'|$. Therefore, together with the leading term in $\tilde{\nu}_1^-$ and $\tilde{\nu}_2^+$, V_{RSC} is proportional to the product $p'k/m$; V_{mag} is proportional to k . Due to energy-momentum conservation, when k increases (decreases) p' decreases (increases); therefore, the RSC varies slowly as far as the product $p'k/m$ is concerned.

As has been shown elsewhere, for energetic photons produced in the pp bremsstrahlung process, the dominant current is the magnetization current, V_{mag} . Moreover, among the various ways in which a photon may be created in a NN bremsstrahlung reaction, the dominant one is that in which the photon is emitted before the strong interaction takes place. The spin matrix element of V_{mag} is, then, proportional to

$$\begin{aligned} \langle S' | V_{\text{mag}} | S \rangle &\propto (1 - \delta_{S',0} \delta_{S,0}) \\ &\times [(-)^{S_1} \tilde{\mu}_1 f(\vec{p}', \vec{p} - \vec{k}/2) + (-)^{S_2} \tilde{\mu}_2 f(\vec{p}', \vec{p} + \vec{k}/2)], \end{aligned} \quad (5)$$

where $S(S')$ denotes the NN total spin in the initial (final) state and f is a function common to $\tilde{\mu}_1^-$ and $\tilde{\mu}_2^+$. For the range of photon momentum $\vec{k}$ involved in most of the existing NN bremsstrahlung experiments, both $f(\vec{p}', \vec{p} - \vec{k}/2)$ and $f(\vec{p}', \vec{p} + \vec{k}/2)$ have the same sign. Then, since to leading order $\tilde{\mu}_1^- = \tilde{\mu}_2^+$ for identical particles, the above matrix element will be suppressed to a large extent if a spin-singlet state is present in either the initial or final state. In particular, the strong 1S_0 state only contributes significantly to the $pp\gamma$ cross section for those kinematical conditions where the cancellation in eq.(5) is minimized.

It is now easy to understand the effect of the RSC. Given the form of the spin matrix element of V_{mag} (eq.(5) and comparing eq.(3) and eq.(4), it is easy to see that the spin matrix element of V_{RSC} is proportional to

$$\begin{aligned} \langle S' | V_{\text{RSC}} | S \rangle &\propto -(1 - \delta_{S',0} \delta_{S,0}) \\ &\times [(-)^{S_1} \tilde{\nu}_1 g(\vec{p}', \vec{p} - \vec{k}/2) - (-)^{S_2} \tilde{\nu}_2 g(\vec{p}', \vec{p} + \vec{k}/2)], \end{aligned} \quad (6)$$

where g is a function common to $\tilde{\nu}_1^-$ and $\tilde{\nu}_2^+$. To leading order, $\tilde{\nu}_1^- = \tilde{\nu}_2^+ = (\mu - 1/2)k/m$ as mentioned before. Unlike the magnetization current contribution, the two terms in brackets in the above matrix element differ by a minus sign. These two terms correspond to contributions from the interacting nucleons 1 and 2, respectively. The origin of this relative sign is due to the fact that V_{RSC} is proportional to terms involving the cross product of the Pauli spin matrix with the sum of

the nucleon momenta in the initial and final state, i.e., $\vec{p}_1 + \vec{p}_1'$ in the first term and $\vec{p}_2 + \vec{p}_2'$ in the second term. In the NN c.m. frame the two interacting nucleons have momenta which differ by a minus sign. This does not occur in the case of the magnetization current contribution, where the photon momentum appears in place of the nucleon momentum in both terms (1 and 2). Due to this relative minus sign, the contribution from singlet states in general, and the strong 1S_0 state in particular, will not be suppressed as it is in the magnetization current term. Therefore, when the two nucleons interact after losing energy by emitting energetic photons, the 1S_0 state contribution becomes significant. This occurs because the final state c.m. energy is small in this case which favors scattering in the 1S_0 state where the T-matrix element is very large. We emphasize that the contribution from V_{RSC} itself is much smaller than that from V_{mag} but what makes the RSC term, V_{RSC} , so important is the interference between V_{RSC} and the dominant magnetic current term V_{mag} .

The features just discussed above are illustrated in Fig. 1. First, we note that in the relativistic calculation which includes the RSC the cross sections for forward and backward photon emission angles are reduced by $\sim 25\%$ at $T_{\text{lab}} = 280$ MeV and by $\sim 50\%$ at $T_{\text{lab}} = 100$ MeV with respect to the corresponding non-relativistic results. The reason for relatively larger RSC effect at $T_{\text{lab}} = 100$ MeV than at $T_{\text{lab}} = 280$ MeV is due to the relatively larger influence of the 1S_0 state at lower incident energy. The dashed curves are the results when the 1S_0 state contribution from V_{RSC} (and only from V_{RSC}) is switched off. As discussed above, the RSC is dominated by scattering in the 1S_0 state; this is particularly clear at $T_{\text{lab}} = 100$ MeV where the RSC is almost entirely due to the strong 1S_0 state contribution.

In a coplanar geometry experiment, the RSC becomes more pronounced as the proton scattering angles decrease. This is illustrated in Fig. 2 where the $pp\gamma$ cross section is shown as a function of symmetric proton scattering angles, $\theta_1 = \theta_2$, and for a fixed photon emission angle. We see that around $\theta_1 = \theta_2 = 5^\circ$ the RSC (solid curve) reduces the cross section by more than a factor of 2 compared to the non-relativistic result. This is particularly relevant for experiments currently being performed at IUCF [2]. As we have shown in Fig. 1b, for sufficiently small forward proton scattering angles, the RSC is large even at $T_{\text{lab}} = 100$ MeV. In Fig. 2, a nearly constant RSC effect for $\theta_1 = \theta_2 > 10^\circ$ where the photon energy is rapidly decreasing is due to the fact that in this region of proton scattering angles the T-matrix in the 1S_0 state is much smaller and less energy dependent than at smaller proton angles. In addition, the decrease in the photon energy is to a

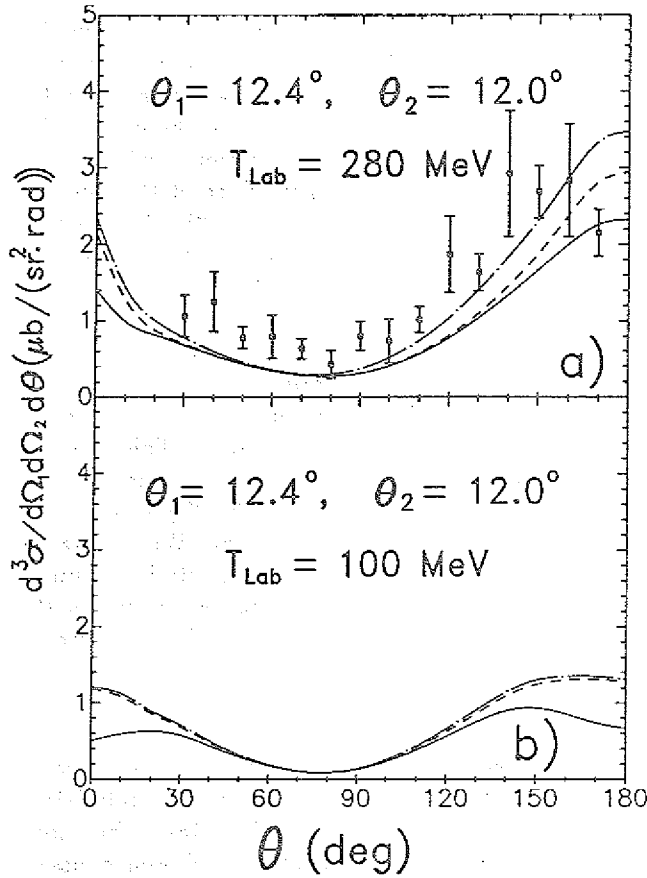


Fig. 1. $pp\gamma$ cross sections in the relativistic (solid lines) and non-relativistic (dash-dotted lines) approaches. Relativistic results when the contribution from the 1S_0 state in V_{RSC} (see eq.(1)) is switched off are represented by dashed lines. The one-boson-exchange potential of Ref. 3 has been used for generating the NN T-matrix.

(a) coplanar geometry cross section in the laboratory frame as a function of photon emission angle at a proton incident energy of $T_{lab} = 280$ MeV and for proton scattering angles of $\theta_1 = 12.4^\circ$ and $\theta_2 = 12.0^\circ$. The data are from Ref. 4 and do not contain the arbitrary normalization factor of $2/3$.

(b) same as (a) at $T_{lab} = 100$ MeV.

large extend compensated by the increase in the nucleon relative momentum p' as discussed before.

In summary, we have shown in a very simple way that the RSC in pp bremsstrahlung reactions is large when the emitted photon has an energy in the region near the maximum value allowed kinematically, even when the incident energy is relatively low. This is due, primarily, to the 1S_0 state contribution which is large and strongly energy dependent compared to other states.

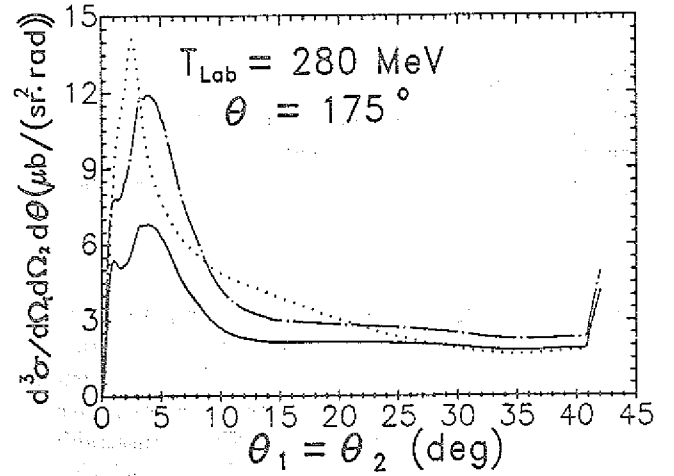


Fig. 2. $pp\gamma$ coplanar geometry cross section in the laboratory frame as a function of symmetric proton scattering angle $\theta_1 = \theta_2$ at a proton incident energy of $T_{lab} = 280$ MeV and for a fixed photon emission angle of $\theta = 175^\circ$. The relativistic (solid lines) and non-relativistic (dash-dotted lines) approaches are shown. Relativistic on-shell results according to Ref. 5 are represented by dotted lines. The one-boson-exchange potential of Ref. 3 has been used for generating the NN T-matrix elements.

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A comparison of nucleon—nucleon potential models in proton—proton bremsstrahlung

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Recent investigations of the proton—proton (pp) bremsstrahlung reaction have demonstrated that off-shell effects of the nucleon—nucleon (NN) interaction can be clearly seen in this type of reaction under certain kinematical conditions. Once the sensitivity of the $pp\gamma$ process to the off-shell behavior of the NN interaction is established for different observables, the next interesting and natural question to ask is to what extent the NN bremsstrahlung reaction can distinguish different off-shell behaviors of modern realistic NN interactions. This question has been addressed by a number of authors in the past and, more recently, by the TRIUMF group [1,2] and is also the issue of the present work. We calculate the pp bremsstrahlung process within a relativistic approach so that the important relativistic corrections are taken into account. The calculation includes the single-scattering contribution as well as the rescattering contribution within the relativistic approach.

In the present work we consider three different interactions, two of them are based on the one-boson-exchange (OBE) model for the NN force developed by the Bonn group [3,4]; the third one is based on the PARIS potential [5]. The two Bonn potentials considered are: OBEPQ and OBEPT potentials [3,4]; the latter is an energy-dependent one.

The different off-shell behavior exhibited by different NN interactions arises either from the differences in the structure of the corresponding NN potentials or from the different approximations made in obtaining a T-matrix interaction from a particular NN potential.

The NN bremsstrahlung reaction requires NN T-matrix elements (in the NN c.m. frame) of the form

$$T(|\vec{p}' \pm \frac{\vec{k}}{2}|, p; E_p) \text{ and } T(|\vec{p} \pm \frac{\vec{k}}{2}|, p'; E_{p'}) \quad (1)$$

corresponding to the process in which the photon is emitted after and before the strong interaction takes place. In eq.(1), $\vec{k}$ denotes the momentum of the emitted photon. E_q , with $\vec{q}$ denoting either $\vec{p}$ or $\vec{p}'$ is the starting energy. "Half off the energy-shell" or simply "off-shell" means that $|\vec{q}' \pm \vec{k}/2| \neq |\vec{q}|$.

In NN elastic scattering, where no photon is present, $|\vec{q}'| = |\vec{q}|$ (on-shell) by energy-momentum conservation.

Since the on-shell limit of the NN T-matrix interaction from eq.(1) is not unique, we have to specify which on-shell limit we are considering when discussing off-shell effects. The on-shell limit we are interested in here is

$$\begin{aligned} T(|\vec{p}' \pm \frac{\vec{k}}{2}|, p; E_p) &\rightarrow T(p, p; E_p) \\ T(|\vec{p} \pm \frac{\vec{k}}{2}|, p'; E_{p'}) &\rightarrow T(p', p'; E_{p'}) \end{aligned} \quad (2)$$

We refer to these on-shell limits as the OESA. In this OESA limit the T-matrix elements, and only the T-matrix elements, are forced to be on-shell according to eq.(2).

Although, realistic NN interactions are said to be equivalent on the energy-shell (phase-shift equivalent) they are not completely equivalent because of the uncertainties arising in fitting the phase-shifts. Since we want to test the sensitivity of the pp bremsstrahlung reaction to off-shell differences of these interactions we have to check whether they yield the same on-shell results. Fig.1a shows a comparison of the OESA (eq.(2)) on-shell results for the coplanar geometry $pp\gamma$ cross section for NN interactions based on the OBEPQ (solid), OBEPT (dashed) and PARIS (dash-dotted) potentials. The incident energy is $T_{lab} = 280$ MeV and the proton scattering angles are $\theta_1 = 12.4^\circ$ and $\theta_2 = 12.0^\circ$. The data are from Ref. 6. We see that as far as the pp bremsstrahlung process in the coplanar geometry is concerned, the OBEPQ and PARIS based T-matrices yield practically the same on-shell results. We observe that the OBEPQ and PARIS interactions have the same NN propagators. The difference observed between the on-shell results based on the OBEPT interaction and those based on the other two interactions is due only to the difference in the corresponding NN propagators G in the extra intermediate states which are present in the NN bremsstrahlung process compared to the NN elastic scattering process because of the emission of the photon. Indeed, if we force G in those intermediate states (and only in those intermediate states) to be the same for all these three interactions, the resulting $pp\gamma$ cross sections

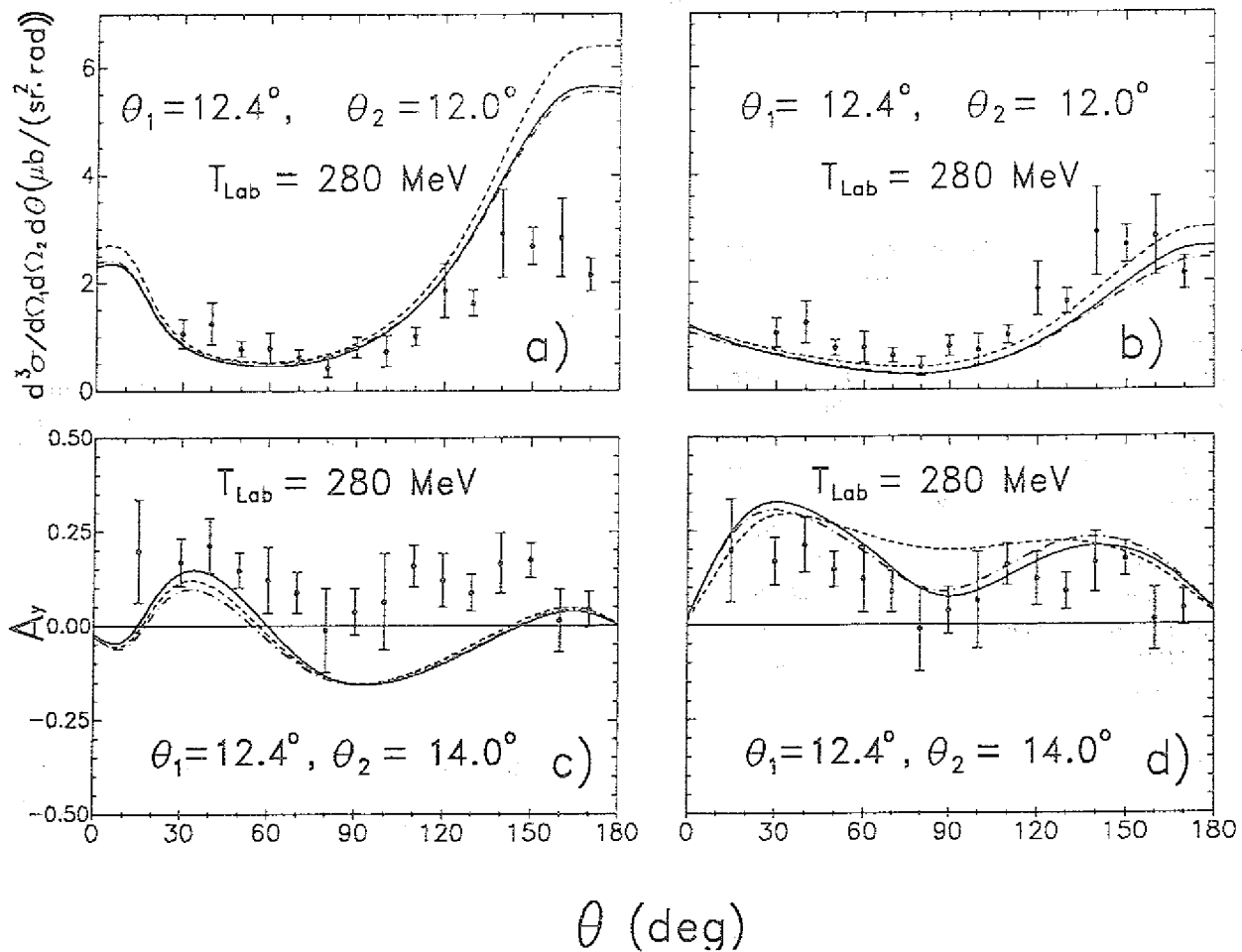


Fig. 1. Coplanar geometry $pp\gamma$ cross section and analyzing power in the laboratory frame as a function of photon emission angle θ at an incident energy of $T_{\text{Lab}} = 280$ MeV. The solid, dashed and dash-dotted curves are the results corresponding to the OBEPQ, OBEPT [3] and PARIS [5] potentials, respectively. Parts (a) and (b) correspond to the OESA (eq.(2)) and exact (off-shell) cross sections for $\theta_1 = 12.4^\circ$ and $\theta_2 = 12.0^\circ$, respectively. The cross section data have not been renormalized as was done in Ref. 6. Parts (c) and (d) correspond to the OESA and exact (off-shell) results, respectively, for the analyzing power for $\theta_1 = 12.4^\circ$ and $\theta_2 = 14.0^\circ$. The data [6] are multiplied by a factor of -1 to agree with our convention.

in the OESA limit are practically the same for all these interactions. Therefore, as far as the pp bremsstrahlung process is concerned, all the three interactions have the same on-shell values.

In Fig. 1b the corresponding off-shell results as prescribed by the NN bremsstrahlung reaction are shown. Here, we do not see any significant differences among these results, indicating that the $pp\gamma$ cross section in this geometry is not sensitive enough to distinguish among the differences in the off-shell behavior of these three interactions. Figs. 1c and 1d show the on- and off-shell results, respectively, for the analyzing power in a similar kinematical condition. A comparison of these two figures shows clearly that this quantity is sufficiently sensitive to reveal the off-shell difference of the OBEPT-based interaction from the other two interactions manifested in the results around $\theta = 90^\circ$. Here, we stress that, although the OBEPT version predicts a result which is in

disagreement with the data, a caution must be taken to conclude from this that its off-shell behavior is wrong. This is because the theory of the NN bremsstrahlung reaction and, in particular, the present potential model calculation, is not yet complete.

Very recently, we have investigated the sensitivity of spin-correlation coefficients in pp bremsstrahlung to off-shell effects of the NN interaction [7] and have found that these coefficients are very sensitive to off-shell effects, even at proton scattering angles where the analyzing power is weakly sensitive to these effects. In Fig. 2 we show the on- and off-shell results for the spin-correlation coefficients C_{xx} , C_{yy} and C_{zz} , as defined in Ref. 7, at an incident energy of $T_{\text{Lab}} = 280$ MeV and for symmetric proton scattering angles of $\theta_1 = \theta_2 = 10^\circ$. In the upper part of the figure we also display the corresponding analyzing power for comparison. We see that all three

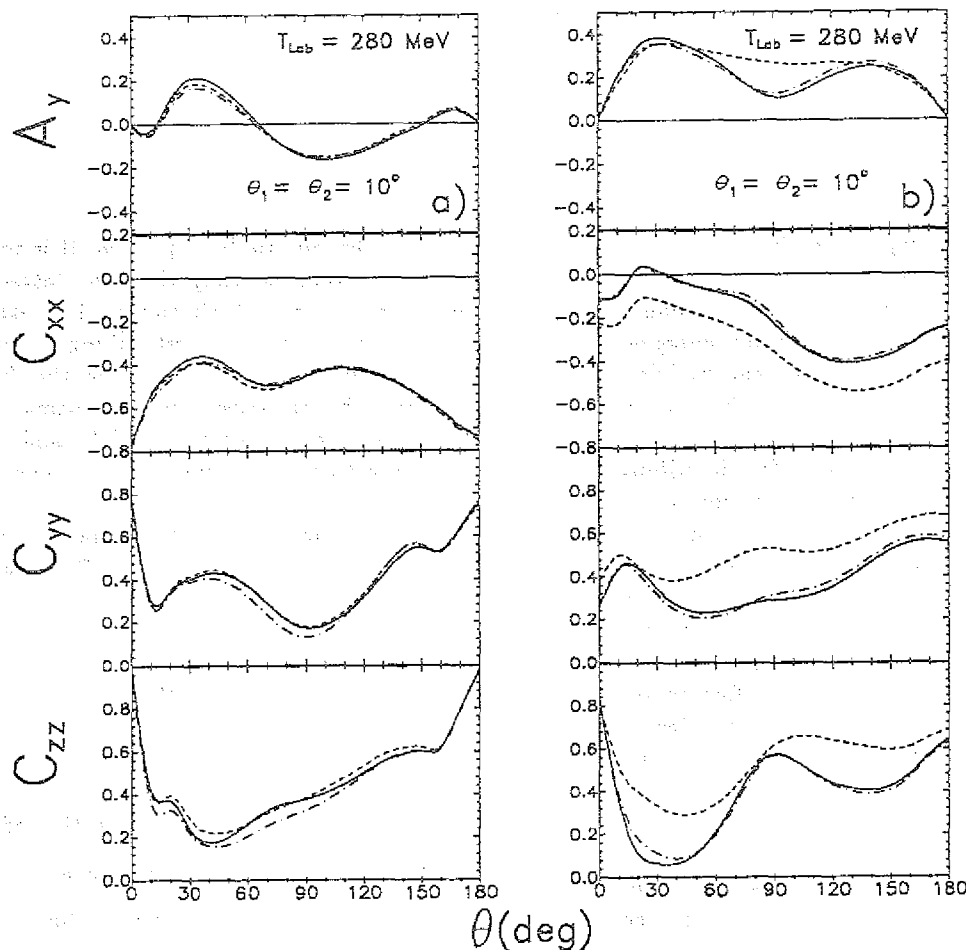


Fig. 2. Coplanar geometry $pp\gamma$ spin-correlation coefficients and analyzing power in the laboratory frame as a function of photon emission angle θ at an incident energy of $T_{\text{Lab}} = 280$ MeV and at symmetric proton scattering angles of $\theta_1 = \theta_2 = 10^\circ$. From the top to bottom are the analyzing power A_y and the spin-correlation coefficients C_{xx} , C_{yy} and C_{zz} . Part (a) corresponds to the on-shell results, while part (b) corresponds to the exact (off-shell) results. For details see caption of Fig.1.

interactions considered yield practically the same on-shell results (Fig. 2a). In Fig. 2b where the off-shell results are shown, we see a clear difference between results based on the OBEPT interaction and those based on the OBEPQ and PARIS interactions. The latter two interactions exhibit basically the same off-shell results.

We have tested the sensitivity of the $pp\gamma$ reaction to the different off-shell behaviors of three NN interactions based on modern potentials. We have found that cross section measurements are unable to distinguish between these off-shell differences. In order to discriminate between different interactions, one would need much larger off-shell differences than those generated by modern interactions. The analyzing power, however, is sufficiently sensitive to distinguish between the off-shell differences in the tensor component of different NN interactions. It is shown that measurements of spin-correlation coefficients offer a potential means of distinguishing between

different interactions based on their off-shell differences.

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Hard photon production in proton-deuteron reactions at intermediate energies

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There has been a resurgence of interest in proton-neutron (pn) bremsstrahlung reactions since the first observation of energetic photons produced in heavy-ion collisions at intermediate energies was reported a few years ago. Theoretical models using transport equations to describe the heavy-ion dynamics require the elementary pn bremsstrahlung process as the basic ingredient for describing the photon production in these heavy-ion collisions. However, due to experimental difficulties, pn bremsstrahlung data are far too sparse to provide the necessary information required in heavy-ion calculations. Consequently, these calculations have to rely on theoretical predictions of the elementary pn bremsstrahlung process. In an effort to provide reliable input for calculations of heavy-ion photo-production, the pn bremsstrahlung process based on meson-exchange models has been investigated theoretically [1,2].

Quite recently, pd bremsstrahlung cross sections have been measured at a proton incident energy of $T_{\text{lab}} = 200$ MeV by the Grenoble group [3] and at $T_{\text{lab}} = 145$ and 195 MeV by the Michigan State group [4] under kinematical conditions appropriate for understanding of the pn bremsstrahlung process as it enters the calculation of energetic photon production in heavy-ion collisions. Although the pd bremsstrahlung process is different from the

pure pn bremsstrahlung process, it is much easier to measure experimentally than the latter process and offers an important alternative for testing theoretical models of pn bremsstrahlung reactions. In the present work we calculate, for the first time, pd bremsstrahlung cross sections using modern meson-exchange potential models and compare the results with the data mentioned above.

The formalism for calculating pd bremsstrahlung is basically the same as given in Refs.1,5. We express the photon emission probability per unit solid angle and unit photon energy as

$$\frac{d^2P_{pd\gamma}}{d\omega d\Omega} = \frac{1}{W_{NN}} \frac{d^2W_{pd\gamma}}{d\omega d\Omega}, \quad (1)$$

where $d^2W_{pd\gamma}/d\omega d\Omega$ denotes the pd bremsstrahlung differential transition rate and W_{NN} is the total nucleon-nucleon (NN) collisional rate irrespective of the bremsstrahlung process.

We calculate the pd bremsstrahlung transition rate by simply folding the elementary NN bremsstrahlung transition rate with the momentum distribution function of the nucleons in the target deuteron

$$\begin{aligned} \frac{d^2W_{pd\gamma}}{d\omega d\Omega} &= \frac{1}{\epsilon_1} \int \frac{d^3p_2}{(2\pi)^3 \epsilon_2} n(\vec{p}_2) 2\pi\omega \int \frac{d^3p'_1}{(2\pi)^3 \epsilon'_1} \frac{d^3p'_2}{(2\pi)^3 \epsilon'_2} Q \\ &\times \left\{ \frac{1}{4} \sum_{SM_S} \sum_{\vec{\epsilon}} \left| \sqrt{\epsilon'_1 \epsilon'_2 \omega} < \vec{\epsilon}, \vec{k}; \vec{p}_1', \vec{p}_2', S' M_{S'} | V_{em} | 0; \vec{p}_1, \vec{p}_2, S M_S > \sqrt{\epsilon_1 \epsilon_2} \right|^2 \right\} \\ &\times \delta^3(\vec{p}_1 + \vec{p}_2 - \vec{p}_1' - \vec{p}_2' - \vec{k}) \delta(\epsilon_1 + \epsilon_2 - \epsilon'_1 - \epsilon'_2 - \omega). \end{aligned} \quad (2)$$

In the above equation, the primed (unprimed) quantities refer to the final (initial) two interacting nucleons' energies ϵ'_1, ϵ'_2 (ϵ_1, ϵ_2), momenta $\vec{p}_1', \vec{p}_2'$ ($\vec{p}_1, \vec{p}_2$), total spin S' (S), and spin projection $M_{S'}'$ (M_S). The matrix element in curly brackets denotes the photon emission amplitude for a photon

with polarization $\vec{\epsilon}$ and momentum $\vec{k}$; ω stands for the photon energy. The momentum distribution function of the nucleons inside the target deuteron is denoted by n and Q is the Pauli blocking operator which prevents nucleons in the final state from having the same momenta.

Analogously, the total NN collisional rate is given by

$$W_{NN} = \frac{1}{\varepsilon_1} \int \frac{d^3 p_2}{(2\pi)^3 \varepsilon_2} n(\vec{p}_2) \int \frac{d^3 p'_1}{(2\pi)^3 \varepsilon'_1} \frac{d^3 p'_2}{(2\pi)^3 \varepsilon'_2} Q \quad (3)$$

$$\times \left\{ \frac{1}{4} \sum_{\substack{SM_S \\ S'M_{S'}}} |\sqrt{\varepsilon'_1 \varepsilon'_2} \langle \vec{p}_1', \vec{p}_2', S'M_{S'} | T | \vec{p}_1, \vec{p}_2, SM_S \rangle \sqrt{\varepsilon_1 \varepsilon_2}|^2 \right\}$$

$$\times (2\pi)^4 \delta^3(\vec{p}_1 + \vec{p}_2 - \vec{p}_1' - \vec{p}_2') \delta(\varepsilon_1 + \varepsilon_2 - \varepsilon'_1 - \varepsilon'_2),$$

where T denotes the NN T matrix associated with the bare NN potential used.

The single-particle energies appearing in eq.(2) and eq.(3) are assumed to have the relativistic form

$\varepsilon(p) = \sqrt{p^2 + m^2}$, where m stands for the nucleon mass.

The elementary NN bremsstrahlung amplitude has been calculated [1,2] within a meson exchange potential model where the strongly interacting particles are treated to all orders in perturbation theory while the coupling to the photon is considered only in first order via the minimal substitution in the non-relativistic nuclear Hamiltonian. The model includes the convection, magnetization, as well as important exchange current contributions. Since the proton-proton (pp) bremsstrahlung process is known to be negligible compared to the pn process [2], we consider only the latter reaction.

The photon cross section is obtained by multiplying the photon emission probability rate given by eq.(1) by the total pd cross section, σ_{pd} ,

$$\frac{d^2 \sigma}{d\omega d\Omega} = \sigma_{pd} \frac{d^2 P_{pd \gamma}}{d\omega d\Omega}. \quad (4)$$

Fig. 1 shows the results for pd bremsstrahlung cross sections as a function of photon energy at a proton incident energy of $T_{lab} = 195$ MeV and for various photon emission angles. For the pd total cross section we use the experimental value of $\sigma_{pd} = 61 \pm 4$ mb at $T_{lab} = 208$ MeV [6]. The data are from the Michigan State group [4]. In order to access the level of agreement of the calculated and measured shapes of the data, the theoretical results have been multiplied by a normalization factor of $N=1.67$. This would correspond to using the total pd cross section of $\sigma_{pd} = 100$ mb instead of its experimental value of $\sigma_{pd} = 61 \pm 4$ mb. We see that, overall, the shape of both the energy and angular distributions are in good agreement with the data, although, absolute cross sections are underestimated by $\sim 40\%$. Similar results are obtained for the cross section data obtained by the Grenoble group [3]. In view of the pd

bremsstrahlung data from Refs.3,4 being much more accurate than the existing pn bremsstrahlung data, the underprediction of $\sim 40\%$ observed in Fig.1 which is consistent with the results obtained for the elementary pn bremsstrahlung reaction [1,2] strongly suggests that some significant contribution to the elementary pn bremsstrahlung process is still not taken into account. It is, however, not easy to think of a process(es) which gives such a significant contribution, at least in the NN center-of-mass energy domain involved in the present calculation. The only obvious term which has been left out of the present calculation is the relativistic spin correction. This correction is known to reduce the cross section considerably in pp bremsstrahlung [7,8]; if the inclusion of this term should also reduce the pn bremsstrahlung cross section, this would further increase the discrepancy with the existing data.

In summary, we have investigated the production of hard photons in intermediate energy pd reactions. As far as the shape of both the energy and angular distributions are concerned, the present calculation is, overall, in good agreement with those observed in recent data at $T_{lab} = 195$ MeV and 200 MeV. Absolute cross sections, however, are underpredicted by $\sim 40\%$ which is consistent with the results obtained for the pn bremsstrahlung reaction and it suggests that further work should be done for a better understanding of the elementary pn bremsstrahlung process. This underprediction of the pn bremsstrahlung cross section should be considered when applying our pn bremsstrahlung calculation to the description of more complicated processes. At $T_{lab} = 145$ MeV incident energy, the present model underestimates the cross section more in the low photon energy region than in high energy region (not shown). This kind of discrepancy is not easy to understand since, for pd bremsstrahlung, we do not expect secondary collisions to play a significant role. Such a discrepancy may be due to the inadequacy of the present model. It would be interesting to have pd bremsstrahlung data at lower proton incident energies in order to see whether the discrepancy increases or decreases at lower energies. This could provide helpful insight as to the origin of the disagreement.

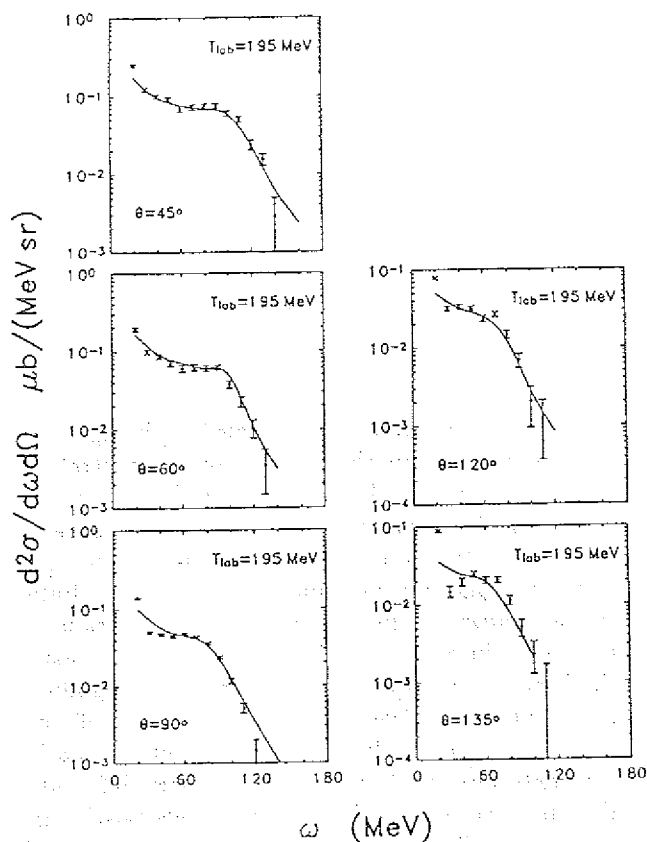


Fig. 1. Proton-deuteron bremsstrahlung cross section in the laboratory frame as a function of photon energy ω at a proton incident energy of $T_{lab} = 195$ MeV and for various photon emission angles θ . The solid lines correspond to the present prediction multiplied by an arbitrary normalization factor of $N = 1.67$. The data are from Ref.4.

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Consistent two-phase model of the nucleon

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In the following we assume that the nucleon consists of a bare nucleon (later referred to as the core) consisting of u and d quarks and a component which contains a virtual meson and an associated recoil baryon (later referred to as the cloud). As a natural consequence of this fact the quark distribution functions have the following structure:

$$q_f(x) = q_f^{\text{val}}(x) + q_f^{\text{sea}}(x), \quad (1)$$

$$q_f^{\text{val}}(x) = q_{f_{\text{core}}}^{\text{val}}(x) + q_{f_{\text{cloud}}}^{\text{val}}(x), \quad (2)$$

$$q_f^{\text{sea}}(x) = q_{f_{\text{core}}}^{\text{sea}}(x) + q_{f_{\text{cloud}}}^{\text{sea}}(x), \quad (3)$$

where q_f stands for $u, d, s, \bar{u}, \bar{d}$ and $\bar{s}$, respectively.

In the present approach we take into account explicitly the fact that the core constitutes only a fraction of the nucleon. This requires an additional renormalization, which is crucial in order to assure the number sum rules for valence quarks. The core contributions to the valence-quark distribution functions can be expressed as follows

$$q_{f_{\text{core}}}^{\text{val}}(x) = Z q_{f_{\text{bare}}}^{\text{val}}(x), \quad (4)$$

where Z is an overall renormalization factor. It can be calculated as

$$Z = \frac{1}{1 + \sum_i \int_0^1 f_{M_i}^i(y) dy}, \quad (5)$$

and satisfies $0 \leq Z \leq 1$.

The functions $f_{M_i}^i(y)$ are momentum distributions of the meson M_i in the nucleon. They depend on the coupling constants and form factor. Explicit expressions have already been presented elsewhere [1]. In the present approach, similarly as in ref. [1], in addition to the Sullivan processes involving mesons, we include also those connected to the recoil baryons. This requires, in principle, a calculation of momentum distribution functions for baryons [1], which involve a less known form

factor, because the baryon, not the meson, is off-shell. In this investigation, however, instead of introducing a few new parameters we follow the approach proposed in ref. [2]. That is, we assume that the amplitude $f_B(y)$ to find a baryon B with momentum fraction y accompanied by a meson M is connected to the corresponding amplitude $f_M(y)$ to find a meson M accompanied by the baryon B via the following relation

$$f_B(y) = f_M(1-y). \quad (6)$$

Since, for simplicity, in practical calculation we assume the same cut-off parameter in the form factors for all meson vertices, the above approach allows us to study dependencies on only one parameter instead of investigation a many-dimensional space. The way we treat vector mesons has already been described in the previous contribution [3].

We assume that the distribution functions of the valence quarks of the bare nucleon satisfy the following number sum rules

$$\int_0^1 u_{\text{bare}}^{\text{val}}(x) dx = 2(1), \quad (7)$$

$$\int_0^1 d_{\text{bare}}^{\text{val}}(x) dx = 1(2) \quad (8)$$

for proton or neutron, respectively.

We define the cloud contribution to the valence quark distribution functions as

$$q_{f_{\text{cloud}}}^{\text{val}}(x) = q_f^{\text{Sull}}(x) - q_{f_{\text{core}}}^{\text{Sull}}(x), \quad (9)$$

The quantities on the right-hand sides of the above equations are contributions to the quark distribution functions calculated from the generalized Sullivan processes. Formally they can be written as

$$q_f^{\text{Sull}}(x) = \sum_i q_{f_i}^{\text{Sull},i}(x), \quad (10)$$

where the index i runs over all possible virtual decompositions of the nucleon into meson and

baryon. The contribution of quarks from a meson or the associated baryon can be written as a convolution of the meson (baryon) structure function and its momentum distribution $f_M^i(y)$ in the nucleon [4]

$$xqf_{Sull,i}^{sea}(x,Q^2) = Z \int_x^1 dy f_M^i(y)(x/y)q_{f,M}^i(x/y,Q^2), \quad (11)$$

where $q_{f,M}^i(x/y,Q^2)$ is the quark distribution function in the meson M_i . A similar equation holds for the recoil baryons. The only difference in comparison to ref. [4] is the presence of the factor Z in front of the right hand side of Eq.(11) which takes into account the proper renormalization of the valence quarks.

As a natural consequence of Eq.(9) sea-quark distribution functions may be calculated as follows

$$u_{cloud}^{sea}(x) = \bar{u}_{cloud}^{sea}(x) = \bar{u}^{Sull}(x) \quad (12)$$

$$d_{cloud}^{sea}(x) = \bar{d}_{cloud}^{sea}(x) = \bar{d}^{Sull}(x) \quad (13)$$

$$s_{cloud}^{sea}(x) = s^{Sull}(x) \quad (14)$$

$$\bar{s}_{cloud}^{sea}(x) = \bar{s}^{Sull}(x). \quad (15)$$

According to the definitions (2) and Eqs.(7)–(10) the numbers of valence quarks are given by

$$\begin{aligned} \int_0^1 u^{val}(x)dx &= 2Z + N^{Sull}(u-\bar{u}), \\ \int_0^1 d^{val}(x)dx &= Z + N^{Sull}(d-\bar{d}), \end{aligned} \quad (16)$$

where

$$\begin{aligned} N^{Sull}(u-\bar{u}) &= \int_0^1 (u^{Sull}(x) - \bar{u}^{Sull}(x))dx, \\ N^{Sull}(d-\bar{d}) &= \int_0^1 (d^{Sull}(x) - \bar{d}^{Sull}(x))dx. \end{aligned} \quad (17)$$

It may be shown [5] that

$$\begin{aligned} N^{Sull}(u-\bar{u}) &= 2(1-Z), \\ N^{Sull}(d-\bar{d}) &= (1-Z), \end{aligned} \quad (18)$$

therefore the valence quark distribution functions defined by Eq.(2) also satisfy the required number sum rules

$$\int_0^1 u^{val}(x)dx = 2(1), \quad (19)$$

$$\int_0^1 d^{val}(x)dx = 1(2). \quad (20)$$

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On the role of the vertex form factor for deep-inelastic lepton-nucleon scattering

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The main ingredients of our model [1] are vertex coupling constants, parton distribution functions for mesons and baryons and vertex form factors.

The coupling constants for the nucleon-meson-nucleon vertices have been taken from the Bonn nucleon-nucleon potential [2], while those for the nucleon-meson-hyperon vertices from [3], where they were adjusted to the low-energy nucleon-hyperon scattering data.

The quark distribution functions for mesons and baryons necessary for calculating contributions from the generalized Sullivan processes [1] were obtained in the following way. In the case of the nucleon we use a recently proposed parametrization of Owens [4]. The parametrizations for other baryons can be obtained by the requirement of SU(3) flavor symmetry. In the case of π meson we have adopted the parametrization of the NA3 Collaboration [5]. The structure functions for other pseudo-scalar mesons may be obtained via flavor SU(3) rotations. The approximations for vector mesons have already been discussed in [6].

For the sake of simplicity we use for all vertices the same form factors. In the present paper we parametrize the form factor in the standard dipole form:

$$F(t) = \left(\frac{\Lambda^2 - m_i^2}{\Lambda^2 - t} \right)^2, \quad (1)$$

where m_i is a mass of the meson M_i .

In distinction to ref. [7], where individual cut-offs in different sectors have been varied to get agreement with experimental data on parton distribution functions, in the present paper there is only one cut-off parameter which we shall try to constrain by the experimental data.

We find the data recently obtained by the CCFR Collaboration at Fermilab [8] very useful in limiting the range of the cut-off parameter. In Figure 1 we show the dependence on the values of the cut-off parameter of the following quantities

$$\kappa = \frac{2 \langle x \rangle_s}{\langle x \rangle_u + \langle x \rangle_d}, \quad (2)$$

$$\eta_s = \frac{2 \langle x \rangle_s}{\langle x \rangle_u + \langle x \rangle_d}, \quad (3)$$

$$R_Q = \frac{\langle x \rangle_u + \langle x \rangle_d + \langle x \rangle_s}{\langle x \rangle_u + \langle x \rangle_d + \langle x \rangle_s}. \quad (4)$$

The $\langle x \rangle_{q_f}$ symbols here are defined as

$$\langle x \rangle_{q_f} = \int_0^1 x q_f(x) dx. \quad (5)$$

The horizontal lines in the figure show the range allowed by the experiment. The indicated error bars have been obtained by adding the statistical and systematic error bars from ref. [8] quadratically.

In the present paper we consider two extreme limits as far as the sea-quark distribution functions are considered. In the first limit we assume that the sea-quark distribution of the nucleon can be entirely attributed to the meson cloud [7]. In the second limit we include in addition a contribution from the core. For this purpose we use a very recent parametrization of the sea-quark distributions [4]. The two solid lines in Figure 1 refer to the two mentioned approximations. In our opinion the real situation should be in between those two extreme cases. Since the distribution of valence quarks in the core is theoretically rather unknown, we assume the experimental distributions [4]. For this reason our practical approach is not fully self-consistent. A consistent approach would require a readjustment to experimental data of the valence quark distributions in the core, including the mesonic effects discussed in the present paper. In the present note we will only estimate the possible inconsistency. For this purpose we replace our distribution function for valence quarks by the parametrization of experimental data [4]. The so-corrected results are presented in Figure 1 by the dashed lines. The result shows that the inconsistency for large cut-off parameters is considerable and should be ultimately corrected. The correction in reality might be slightly smaller, since the momentum carried by valence quark in the Owens parametrization [4] is larger than in other parametrizations of experimental data. If we assumed that the sea distribution in the nucleon can be entirely attributed to the meson cloud we have to accept quite large cut-off parameters ($\Lambda > 1600$ MeV). If the other scenario were true, i.e. there were a large

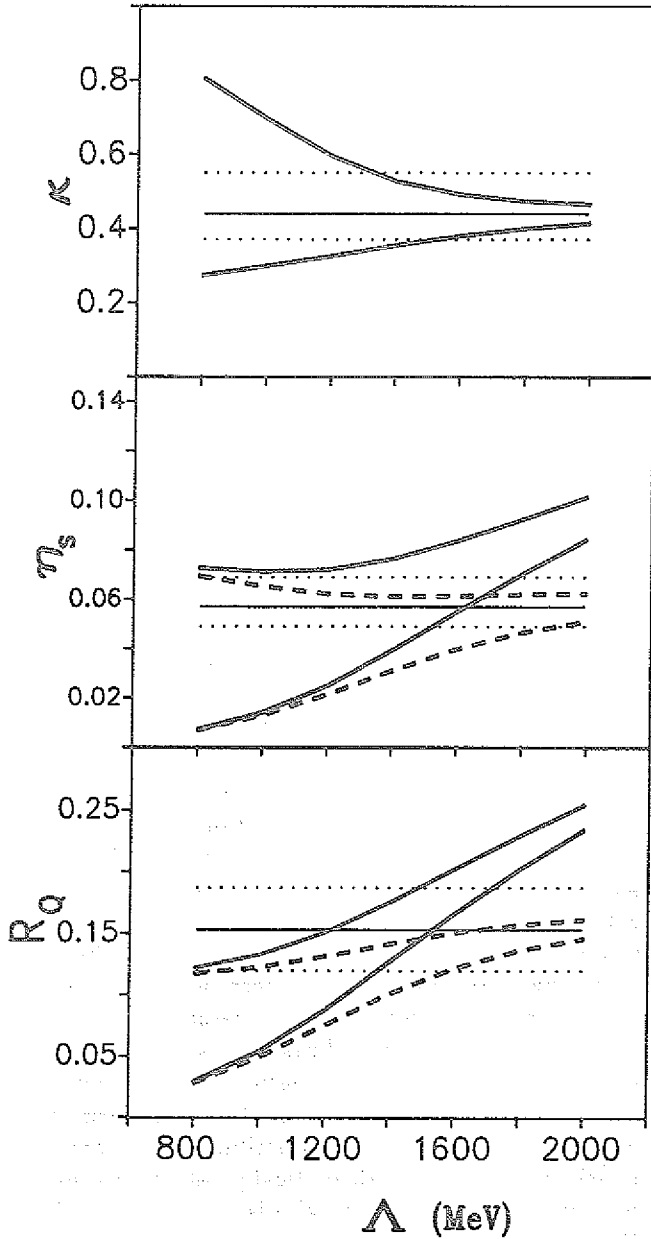


Fig.1. Dependence of K , η_s and R_Q on the cut-off parameter of the form factor compared with experimental result obtained by the CCFR Collaboration [8]. The horizontal lines show the experimental result (thin solid) and the range (thin dotted) allowed by the experiment. The solid lines show the result of our calculation in the two extreme cases, without sea quarks in the core (the lower curve) and with the sea in the core included (the upper curve). The dashed lines denote the calculations with the "model" valence quark distribution functions replaced by the parametrization of experimental data.

amount of perturbative sea in the core, we could use cut-off parameters in a fairly broad range.

It may be concluded from the above analysis that it is possible to get a good description of experimental

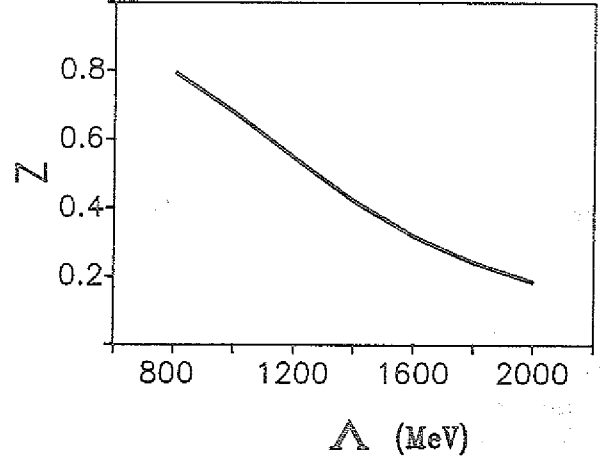


Fig.2. The renormalization coefficient Z as a function of the cut-off parameter of the form factor.

data, at least the global effect, with fairly hard form factors as compared to earlier studies. The large cut-off parameter would have deeper consequences for the nucleon structure. In the Cloudy Bag Model the cut-off parameter Λ is related to the bag radius. The large cut-off parameter would translate to small bag radius. On the other hand, it influences strongly the renormalization factor Z [1], as may be seen from Figure 2. The renormalization factor may be interpreted as a probability to find the core inside the nucleon. Large cut-off parameters imply a rather small probability. This would mean that the bulk properties of the nucleon may be determined by the meson cloud, instead of, as conventionally accepted, by the bare nucleon. This is very interesting in the light of the so-called nucleon spin crisis [9] and requires quantitative analysis. We point out that, within our model, we are able to resolve the known discrepancy between the cut-off parameter used in the nucleon-nucleon potential and that used for the deep-inelastic scattering.

In Table 1 we present contributions of all included Sullivan processes for a given meson-baryon pair for different values of cut-off parameter. As one would expect the most important contributions come from the core, the $\pi + N$ and $\pi + \Delta$ processes. The other processes are almost negligible for small cut-off parameters. They become, however, important for large cut-off parameters. This is especially spectacular for vector mesons. Of course the main reason is the bigger masses of vector mesons (see formula (1)). The low energy phenomenology requires rather large cut-off parameters. If this scenario is true, then different mesons play a crucial role in the description of the nucleon structure and have to be included in the

analysis of the EMC effect and/or in the explanation of the spin structure of the nucleon.

process	800	1200	1600	2000
0.core	79.53	54.54	31.77	18.39
1. $\pi + N$	12.26	19.43	18.64	15.21
2. $\rho + N$		1.59	6.00	9.22
3. $\omega + N$		1.36	4.05	5.07
4. $\sigma + N$	0.45	3.32	4.63	4.31
5. $\eta + N$	0.01	0.11	0.19	0.21
6. $\pi + \Delta$	7.60	16.06	19.90	20.64
7. $\rho + \Delta$		2.32	10.85	20.25
8. $K + \Lambda$	0.12	0.88	1.50	1.67
9. $K^* + \Lambda$		0.09	0.74	1.48
10. $K + \Sigma$	0.01	0.09	0.15	0.17
11. $K^* + \Sigma$		0.06	0.53	1.06
12. $K + Y^*$	0.01	0.06	0.13	0.18
13. $K^* + Y^*$		0.09	0.91	2.14

Table 1. The contributions of different processes in deep-inelastic scattering (in %), for different cut-off parameters (in MeV) of the form factor.

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Consequences of SU(2) flavor symmetry breaking in the nucleon sea for the Drell-Yan processes

A. Szczurek* and J. Speth

The SU(2)_F symmetry breaking in the nucleon sea may also be tested in the Drell-Yan processes. Recently Ellis and Stirling [1] have suggested a new Drell-Yan experiment which could shed new light on the symmetry breaking. They have proposed to measure the dilepton production in the proton-proton and proton-neutron collisions (pN → l⁺l⁻X). The corresponding cross section, distribution in rapidity (y) and the mass of the produced lepton pair (M),

$$\sigma^{pp} \equiv s \frac{d^2\sigma^{pp}}{d\sqrt{\tau}dy}, \quad \sigma^{pn} \equiv s \frac{d^2\sigma^{pn}}{d\sqrt{\tau}dy} \quad (1)$$

is sensitive to the distribution of sea antiquarks

$$\sigma^{pn} = \frac{8\pi\alpha^2}{9\sqrt{\tau}} \sum_f e_f^2 (q_f(x_1, M) \bar{q}_f(x_2, M) + \bar{q}_f(x_2, M) q_f(x_1, M)) \quad (2)$$

where $\tau = M^2/s$, $x_1 = \sqrt{\tau} \exp(y)$, $x_2 = \sqrt{\tau} \exp(-y)$. At $y = 0$ we get

$$\sigma^{pp}|_{y=0} = 2 \text{ const} \left(\frac{4}{9} (u_{val-sea}^{\bar{u}} + u_{sea-sea}^{\bar{u}}) + \frac{1}{9} (d_{val-sea}^{\bar{d}} + d_{sea-sea}^{\bar{d}}) + \frac{1}{9} (s_{sea-sea}^{\bar{s}}) \right) \quad (3)$$

and assuming isospin invariance

$$\sigma^{pn}|_{y=0} = \text{const} \left(\frac{4}{9} (u_{val-sea}^{\bar{u}} + u_{sea-sea}^{\bar{u}} + d_{val-sea}^{\bar{u}} + u_{sea-sea}^{\bar{d}}) + \frac{1}{9} (d_{val-sea}^{\bar{d}} + d_{sea-sea}^{\bar{d}} + u_{val-sea}^{\bar{d}} + u_{sea-sea}^{\bar{s}}) + \frac{1}{9} (s_{sea-sea}^{\bar{s}} + s_{sea-sea}^{\bar{s}}) \right) \quad (4)$$

where all distributions are evaluated at $x = \sqrt{\tau}$ and $Q^2 = M^2$.

Ellis and Stirling have found that the so-called p-n cross section asymmetry defined as

$$A_{DY} = \frac{\sigma_{pp} - \sigma_{pn}}{\sigma_{pp} + \sigma_{pn}}|_{y=0} \quad (5)$$

is very sensitive to the $\bar{u} - \bar{d}$ asymmetry and can

therefore provide a definite measure of isospin breaking in the sea. How it looks in our model [2] is illustrated in Figure 1. For comparison we enclose predictions of three different sets of parton distributions. The results for different parton distributions respecting SU(2) flavor symmetry of the sea are rather similar, but are quite different from the prediction of our model.

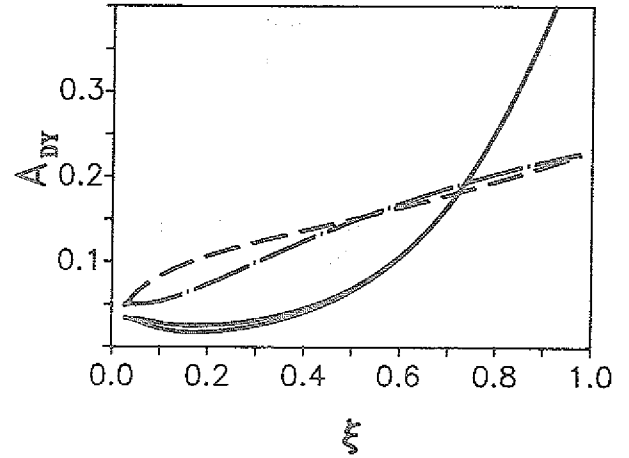


Fig.1. The Drell-Yan rapidity asymmetry A_{DY} as

a function of $\sqrt{\tau}$. The solid line represents predictions of our model. For comparison we present predictions of three sets of parton distributions: dashed curve [3], and dash-dotted curve [4].

Summarizing, we believe that the measurement of the asymmetry with a proton beam on hydrogen and deuterium targets simultaneously (in order to eliminate systematic errors) could give information on the SU(2)_F symmetry breaking in the nucleon's sea, independent of that obtained from the Gottfried sum rule.

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Gottfried sum rule violation

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Recently the NMC group at CERN obtained a new experimental result [1] for the Gottfried Sum Rule

$$S_G(x=0.004-0.8) = 0.227 \pm 0.007(\text{stat.}) \pm 0.014(\text{syst.}) \quad (1)$$

This implies that the sea of the nucleon may break flavor SU(2) symmetry, which is built into almost all current standard parametrizations of parton distributions. Different explanations have been offered.

Since the Gottfried sum rule contains a factor dx/x , it is strongly weighted by values of the quark distribution function at small x . The global effect of the NMC data can be described by the distribution proposed by Martin et al. [2]. In their approach, the behaviour of the valence quark distributions at small x is determined by the data, retaining SU(2)_F symmetry in the sea. As a consequence, more of the sum-rule strength is to be found at very small x than in the conventional picture. Although it predicts a value of 0.23 for the Gottfried Sum Rule, the global fit to $F_2^{\mu p}(x) - F_2^{\mu n}(x)$ is not particularly impressive.

The idea that the sea might not be SU(2) symmetric was suggested in a paper of Field and Feynman [3], who proposed arguments based on the Pauli exclusion principle. According to them the production of $u\bar{u}$ pairs would be suppressed relative to $d\bar{d}$, since the proton contains only one valence down quark but two valence up quarks.

Within the quark-parton model, the deep-inelastic structure functions of the nucleon can be written in terms of the quark distribution functions for proton and neutron, respectively. Assuming isospin invariance one can relate the quark distribution functions of a proton and a neutron to each other. Then the Gottfried sum rule can be written as

$$S_G = 1/3 \int_0^1 dx (u(x) + \bar{u}(x) - d(x) - \bar{d}(x)). \quad (2)$$

In our approach [4] the Gottfried sum rule can be written in the form

$$S_G = Z/3 \int_0^1 dx (u_{\text{bare}}^{\text{val}}(x) - d_{\text{bare}}^{\text{val}}(x)) +$$

$$1/3 \int_0^1 dx (u^{\text{Sull}}(x) + \bar{u}^{\text{Sull}}(x)) -$$

$$1/3 \int_0^1 dx (d^{\text{Sull}}(x) + \bar{d}^{\text{Sull}}(x)). \quad (3)$$

Finally, making use of the number sum rules for valence quarks in the core, it can be written in the simple form

$$S_G = 1/3 (Z + \sum_i A_i) \quad (4)$$

with

$$A_i = \int_0^1 \{ (u_i^{\text{Sull}}(x) + \bar{u}_i^{\text{Sull}}(x)) - (d_i^{\text{Sull}}(x) + \bar{d}_i^{\text{Sull}}(x)) \} dx \quad (5)$$

Since the number sum rules are determined by the valence quarks only one can write S_G as

$$S_G = 1/3 (1 + 2 \int_0^1 dx (\bar{u}^{\text{sea}}(x) - \bar{d}^{\text{sea}}(x))) \quad (6)$$

Finally

$$S_G = 1/3 (1 + 2 \sum_i B_i), \quad (7)$$

where

$$B_i = \int_0^1 \{ \bar{u}_i^{\text{Sull}}(x) - \bar{d}_i^{\text{Sull}}(x) \} dx \quad (8)$$

Form (7) is fully equivalent to (4). The coefficients A_i and B_i are presented in Table 1 for all included Sullivan processes.

Since different authors used different representations for the GSR, there has been a little controversy over which Sullivan processes contribute to the Gottfried sum rule violation. According to Signal et al. [5] "there is no

contribution coming from the pion being struck, the effect is entirely due to the recoil baryon being struck by the incoming photon". On the other hand it is obvious that this is the pion which brings $\bar{u}-\bar{d}$ asymmetry into the nucleon sea. In our opinion the paradox is artificial and comes from a comparison of interpretations based either on Eq.(4) or on Eq.(7), equivalent representations of the Gottfried sum rule. In fact, for full consistency all processes have to be included, since they enter at least in the renormalization factor Z .

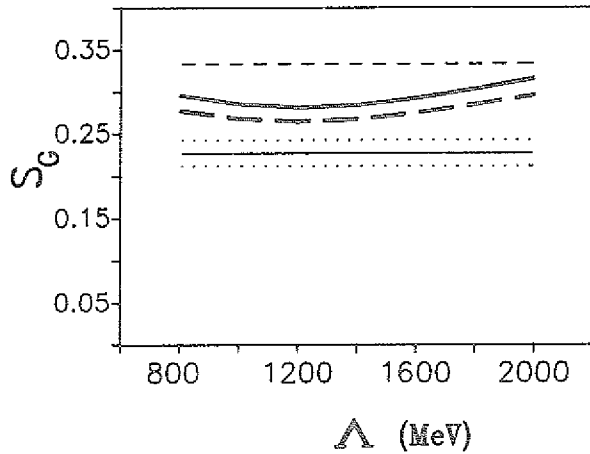


Fig.1. Dependence of the Gottfried sum rule on the momentum cut-off parameter Λ of the form factor. The solid line denotes the result for the integral from 0 to 1, the dashed line denotes the integral from 0.005 to 1. The result of the NMC Collaboration, with experimental uncertainties is shown by the horizontal thin solid and dotted lines. The dashed line, corresponding to the naive quark model, has been displayed for comparison.

In Figure 1 we present the dependence of the Gottfried sum rule on the momentum cut-off parameter Λ of the form factor. The horizontal lines show the range allowed by the NMC data. The results of our calculation are significantly smaller than the predictions of the naive quark model, larger, however, than the experimental value. One should remember, however, that the very small x region was not included in the experimental data. To demonstrate importance of small x we have calculated also the following integrals

$$S_G = \int_{0.005}^1 dx \frac{F_2^{\mu p}(x) - F_2^{\mu n}(x)}{x}. \quad (9)$$

The result is now much closer to the experimental value. This demonstrates the importance of small- x region. The effect is extremely sensitive to the shape of the valence-quark distribution functions in the core, at small x . The effect would be even

larger for distributions like [2]. The best agreement with experimental data occurs for the cut-off parameter of about 1200 MeV. Of course we cannot exclude the presence of some other mechanisms like the Pauli exclusion principle [3], for instance. We find, however, that any phenomenological treatment of such effects cannot be trusted, although it may allow a better description of the experimental data.

process	A_i^M	A_i^B	B_i^M	B_i^B
meson+baryon				
1. $\pi + N$	0	-1/3	-2/3	0
2. $\rho + N$	0	-1/3	-2/3	0
3. $\omega + N$	0	1	0	0
4. $\sigma + N$	0	1	0	0
5. $\eta + N$	0	1	0	0
6. $\pi + \Delta$	0	5/3	1/3	0
7. $\rho + \Delta$	0	5/3	1/3	0
8. $K + \Lambda$	1	0	0	0
9. $K^* + \Lambda$	1	0	0	0
10. $K + \Sigma$	-1/3	4/3	0	0
11. $K^* + \Sigma$	-1/3	4/3	0	0
12. $K + Y^*$	-1/3	4/3	0	0
13. $K^* + Y^*$	-1/3	4/3	0	0

Table 1. The coefficients A_i and B_i are parameters of two different representations of the Gottfried sum rule (4),(7). All numbers in the table have to be multiplied by the probabilities of a given meson-baryon pair inside the nucleon.

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Calculations of E1 resonances in ^{40}Ca , ^{48}Ca and ^{208}Pb
in the 1p1h+2p2h+continuum approximation

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From a microscopic point of view the widths of giant resonances originate from two processes: the coupling of simple RPA excitations with more complex configurations, first of all the 2p2h ones (so-called spreading width $\Gamma\downarrow$) and the decay of excited states into the single particle continuum (escape width $\Gamma\uparrow$). It is clearly necessary now to develop approximate methods to take into account both 2p2h configurations and the single particle continuum and to unify them in one approach. This is done in the present work.

The main idea of our method of taking into account 2p2h configurations is to use a small parameter, i.e., the creation amplitude of the collective phonon g , within the consistent Green function formalism to consider 1p1h $\otimes$ phonon configurations instead of 2p2h ones and to account for a small number of the most collective phonons. In refs. [1,2] a linearized version of the theory has been developed and an equation for the density matrix ρ has been obtained. The basic ideas of the method of taking into account 1p1h $\otimes$ phonon configurations have been shortly described in refs. [3,4].

The integral equation for the transition density ρ , which differs from the RPA-like equations only by the form of the generalized propagator A , has been solved in coordinate representation. This reduces the computational difficulties of the problem appreciably. To obtain the generalized propagator $A(\vec{r}, \vec{r}')$ we have applied the method of the combined $(\vec{r}, \lambda)$ representation, which is similar to that used in [5], to take into account the continuum within the RPA for nuclei with pairing:

$$A(\vec{r}, \vec{r}') = \tilde{A}_{\text{cont}}^{\text{RPA}}(\vec{r}, \vec{r}') + \sum_{1234} (A_{1234} - \tilde{A}_{12}^{\text{RPA}} \delta_{13} \delta_{24}) \tilde{\varphi}_1^*(\vec{r}) \tilde{\varphi}_2(\vec{r}) \tilde{\varphi}_3(\vec{r}') \tilde{\varphi}_4^*(\vec{r}') \quad (1)$$

Here $\tilde{A}_{\text{cont}}^{\text{RPA}}$ means the "refined" 1p1h propagator in which the single particle continuum is taken into account exactly [6]. The summation given in eq.(1) is performed over two shells above and below the

Fermi level. Thus in our method the single particle continuum is included completely only in the 1p1h part of the propagator A . In the 2p2h (1p1h $\otimes$ phonon, to be exact) part of A the continuum is accounted for only approximately.

Due to all this approximations the problem becomes solvable using a relatively simple numerical algorithm even for a nonseparable interaction between the nucleons.

All phonons and E1 resonances in the nuclei ^{40}Ca , ^{48}Ca and ^{208}Pb have been calculated using the same set of the well-known parameters of the Migdal's force:

$$f_{\text{ex}} = -3.74, f_{\text{in}} = -0.002, f'_{\text{ex}} = 2.30, \\ f'_{\text{in}} = 0.76, g = 0.05, g' = 0.96, \\ C_0 = 300 \text{ MeV} \cdot \text{fm}^3$$

In order to simulate the finite experimental resolution we have to introduce a smearing parameter Δ in the calculations. Accordingly, we have taken the values $\Delta = 0.4$ MeV for ^{40}Ca , ^{48}Ca and $\Delta = 0.25$ MeV for ^{208}Pb which are approximately equal to the experimental resolution. For further experimental details see refs. [3,4].

The main results of the calculations which are presented in Figs. 1-3 and Table 1 are as follows:

1. The RPA with continuum does not agree with experiment except, perhaps, for the centroid energy $\bar{E}$ in ^{40}Ca .

2. The inclusion of both the 1p1h $\otimes$ phonon configurations and the continuum shift the $\bar{E}$ values in the required direction, except for the case of ^{40}Ca . We have good

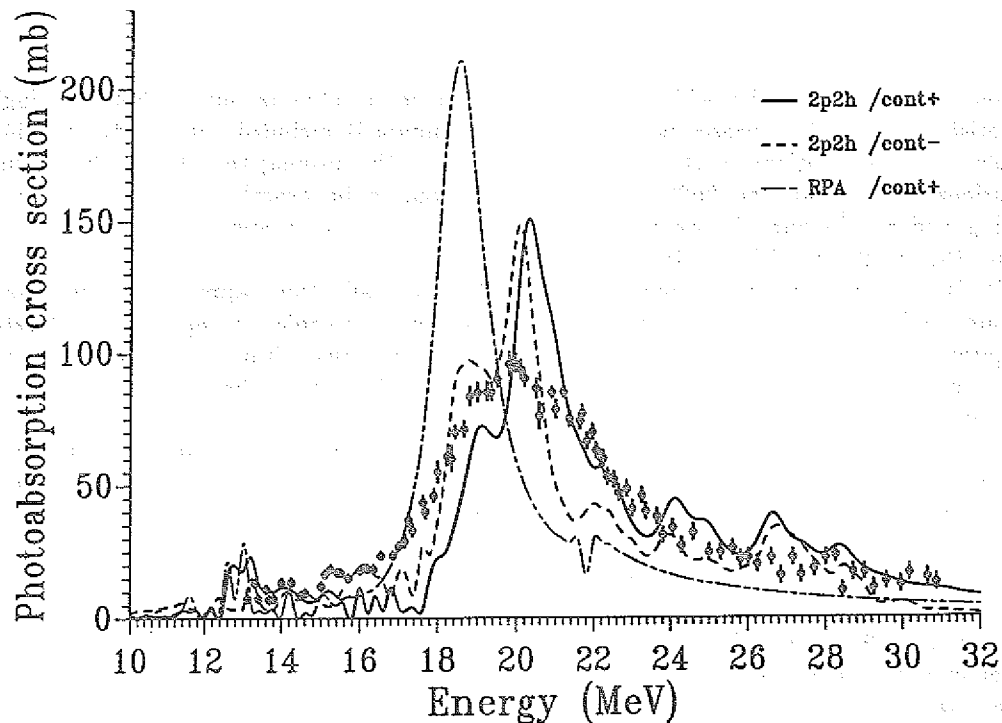


Fig. 1. The E1 photoabsorption cross section for ^{40}Ca obtained in the RPA with continuum (— · — · —), in the $1p1h + 1p1h \otimes \text{phonon}$ (— —) and $1p1h + 1p1h \otimes \text{phonon} + \text{continuum}$ (— · — · —) approximation. The experimental data are taken from ref. [7].

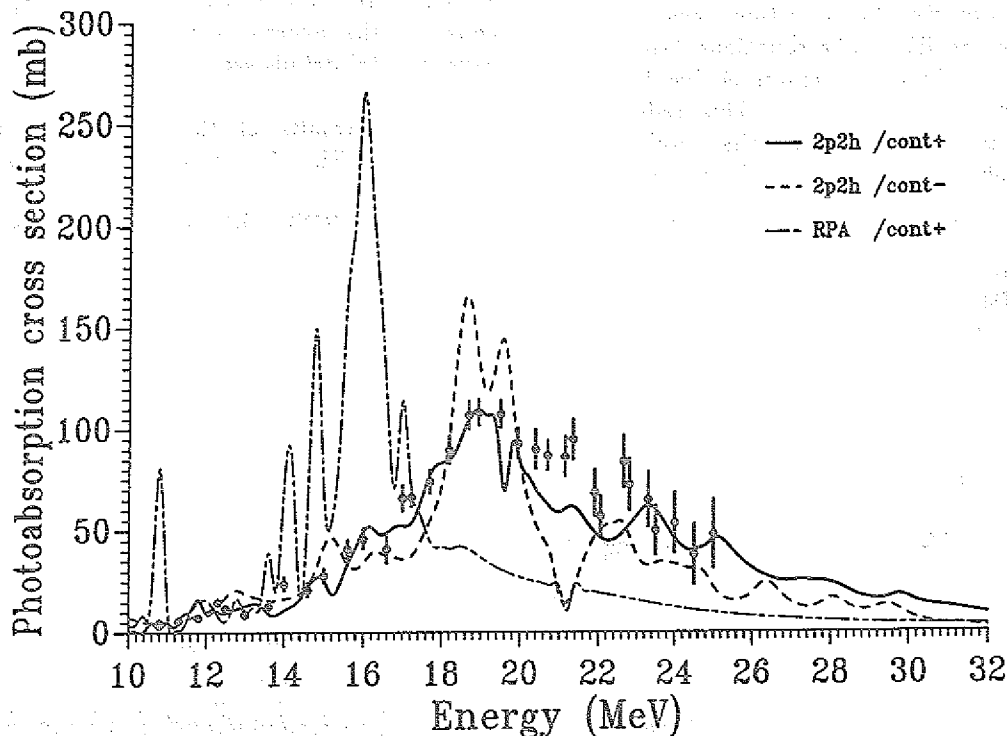


Fig. 2. Same as Fig. 1, but for ^{48}Ca . The experimental data are taken from ref. [8].

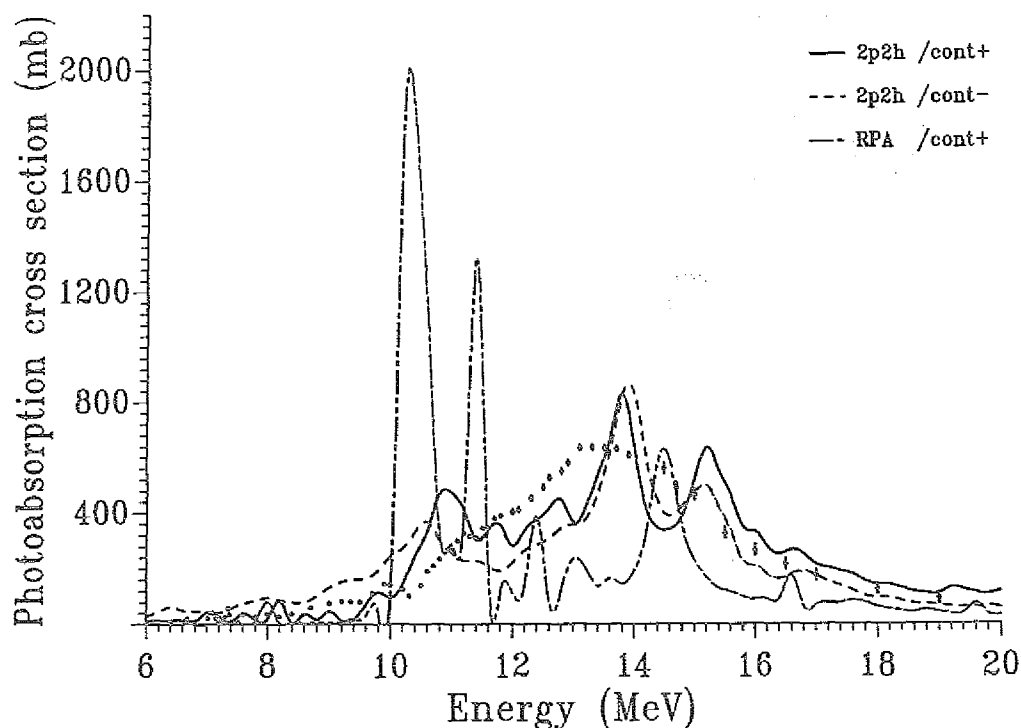


Fig. 3. Same as Fig. 1, but for ^{208}Pb . The experimental data are taken from ref. [9].

	^{40}Ca				^{48}Ca				^{208}Pb			
	σ_{max}	$\bar{E}$	Γ	σ_0	σ_{max}	$\bar{E}$	Γ	σ_0	σ_{max}	$\bar{E}$	Γ	σ_0
	mb	MeV	MeV	mb·MeV								
1p1h+cont.	210	19.3	1.9	586.3	150.0	16.0	3.30	652.5	945	12.2	2.00	2709
1p1h+2p2h	114	20.5	3.7	589.7	118.7	19.0	4.53	696.9	657	13.0	3.60	3110
1p1h+2p2h+cont.	111	21.2	4.3	651.8	101.3	19.6	6.27	756.5	680	13.8	3.87	3367
Exp.	95	20.0	5.0	637.7	102.7	19.6	7.10	836.6	640	13.42	4.05	3306.8
	$\sigma_{\text{cl}}=600$				$\sigma_{\text{cl}}=700$				$\sigma_{\text{cl}}=2984$			

Table 1. Integral characteristics of the E1 resonance.

agreement with experiment for ^{48}Ca and ^{208}Pb .

3. The most important result of our calculation is the satisfactory description of the E1 widths. The two new ingredients (as compared with the RPA) increase the Γ values. Now we underestimate the Γ_{exp} values only by 14% for ^{40}Ca , 11% for ^{48}Ca and 4% for ^{208}Pb .
4. We have obtained a reasonable description of the envelopes of the E1 resonances as one can see from the figures.

To our knowledge this kind of calculation has been done for the first time.

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Microscopic Calculation of the Giant Dipole Resonance in Hot Deformed Nuclei with Realistic Forces

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Giant E1 resonances in hot deformed nuclei are continuing to be studied rather intensively in both experimental and theoretical works. In microscopic calculations of these nuclei separable particle-hole (ph) effective interactions have been used as a rule and the parameters of the forces have been fitted to the position of the giant resonance rather often. For the time being there are no calculations with nonseparable forces for deformed nuclei going beyond the QRPA. It is known, in a consistent QRPA approach, that the pp and ph interactions should be taken into account together with the ph interaction, see e.g. [1]. So far there are no calculations for hot deformed nuclei where the dynamic effects (rather than static ones, e.g. pairing) of these forces have been studied especially together with the ph forces.

Thus the answers provided by the QRPA with realistic nonseparable ph and pp (hh) forces should be reliable for a number of interesting features. The parameters of these forces are known relatively well, e.g. for the Migdal forces see ref. [2].

Here such a program has been executed for the E1 resonance in the hot ^{160}Er nucleus. We have generalized the method of ref. [1] to the case of finite temperature $T \neq 0$ with consistent use of the Migdal force with an established set of parameters.

The system of equations for the quantities $Z^\pm(\omega)$ defining the density matrix which determines the energies of the levels and transition probabilities has been generalized to $T \neq 0$. As a result we have obtained a system of four equations for the quantities $Z^\pm(\omega)$ and $Y^\pm(\omega)$ (instead of two equations for $Z^\pm(\omega)$ at $T = 0$).

$$E_{12}Z_{12}^\pm + \omega Z_{12}^\mp = (1 - n_1 - n_2) \left\{ - \sum_{3,4} [(\eta_{12}^\pm F_{1423} \eta_{34}^\pm + \xi_{12}^\pm F_{1234} \xi_{34}^\pm) Z_{34}^\pm + (\eta_{12}^\pm F_{1423} \xi_{34}^\pm - \xi_{12}^\pm F_{1234} \eta_{34}^\pm) Y_{34}^\pm] - \eta_{12}^\pm e_q V_{12}^{0\pm} \right\}$$

$$E_{12}' Y_{12}^\pm + \omega Y_{12}^\mp = (n_1 - n_2) \left\{ \sum_{3,4} [(\xi_{12}^\pm F_{1423} \eta_{34}^\pm - \eta_{12}^\pm F_{1234} \xi_{34}^\pm) Z_{34}^\pm + (\xi_{12}^\pm F_{1423} \xi_{34}^\pm - \eta_{12}^\pm F_{1234} \eta_{34}^\pm) Y_{34}^\pm] + \xi_{12}^\pm e_q V_{12}^{0\pm} \right\}$$

where $V^0(\omega)$ is an external electromagnetic field, the index 1 stands for the set $\{\lambda_1\}$ of single particle

quantum numbers, and

$$\eta_{12}^\pm = u_1 v_2 \pm v_1 u_2, \quad \xi_{12}^\pm = u_1 u_2 \mp v_1 v_2$$

$$E_1(T) = \sqrt{[\epsilon_1(T) - \mu(T)]^2 + \Delta_1^2(T)}$$

$$n_1 = n(E_1) = \left[1 + \exp \frac{E_1(T)}{T} \right]^{-1}$$

$$E_{12} = E_1 + E_2,$$

$$E_{12}' = E_1 - E_2$$

F and F^ξ are the effective ph(hp) and pp(hh) interactions, respectively. The normalization condition has the form

$$2 \sum_{1,2} \left\{ \frac{\bar{Z}_{12}^\pm \bar{Z}_{12}^\mp}{1 - n_1 - n_2} - \frac{\bar{Y}_{12}^\pm \bar{Y}_{12}^\mp}{n_1 - n_2} \right\} = -1$$

where the quantities $\bar{Z}^\pm, \bar{Y}^\pm$ are the residues of $Z^\pm, Y^\pm$ at the pole $\omega = \omega_s$. ($\bar{Z}^\pm$ corresponds to the quantities $Z^\pm$ of ref. [1].)

The computational scheme is essentially the same as in ref. [1]. We have used a more contemporary set of values for the parameters of Migdal's ph force which has been taken from [2]. The calculations have been performed for $T \leq 2$ MeV; the critical temperature has been obtained as $T_c = 0.55$ MeV for protons, 0.53 MeV for neutrons.

Taking into account that we have not fitted the parameters of the forces, the agreement of the mean energies with the available experimental data at $T = 0$ and $T \approx 1$ MeV [3] is not bad:

$$\bar{E}_{\text{exp}}(0^-) = (12.2 \pm 0.1) \text{ MeV}, \quad \bar{E}_{\text{exp}}(1^-) = (15.2 \pm 0.2) \text{ MeV},$$

and

$$\bar{E}_{\text{th}}(0^-) = 11.3 \text{ MeV}, \quad \bar{E}_{\text{th}}(1^-) = 14.3 \text{ MeV}.$$

The strength ratio $(S_1/S_0)_{\text{th}} = 1.9$ agrees with the experimental value of 1.6 ± 0.5 . The TRK sum rule is exhausted to more than 92%.

In fig. 1 the mean energies $\bar{E}$ and dispersions $\sigma_E = \sqrt{\bar{E} - (\bar{E})^2}$ are given. Only a very weak temperature dependence is seen in the case where both the ph and pp interaction have been used. The corresponding results for the case where the pp force has been ignored in the QRPA look different: here now

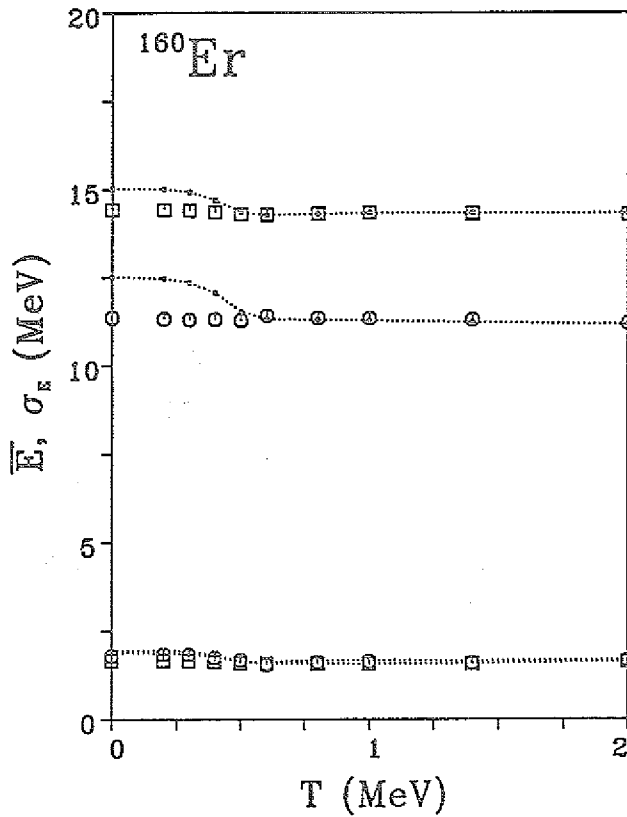


Fig. 1: Mean energies $\bar{E}$ (above 10 MeV) and dispersion σ_E (below 5 MeV) as functions of temperature.

Squares: $|K|=1$, circles: $K=0$ with ph and pp interaction. Dotted lines: the same without pp interaction.

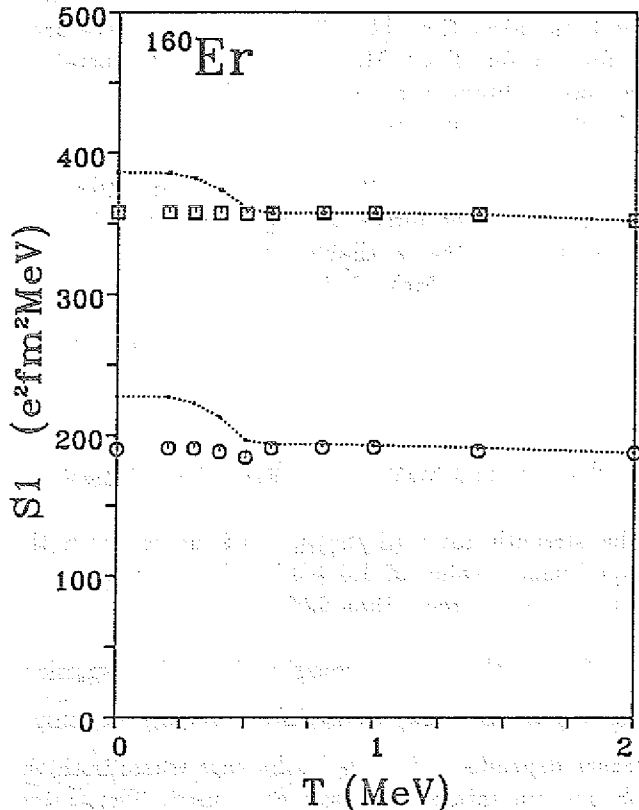


Fig. 2: Energy weighted sum $S1$ as a function of temperature. The notation is the same as in fig. 1.

some more temperature dependence occurs. An analogous and related difference occurs in the energy weighted sum, fig. 2, comparing the full interaction (squares and circles) with the ph-interaction only. Again a noticeable temperature dependence is visible if only the ph-interaction is used. These results demonstrate the importance of a consistent treatment of the ph and pp(hh) forces.

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Low-lying 1^+ levels in spherical open shell nuclei

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In modern nuclear physics experiments it is possible to measure weak 1^+ levels with small $B(M1)$ values, see e.g. [1,2]. As is well known, this led to the discovery of low-lying 1^+ levels in deformed nuclei which are discussed now very actively [1]. There is also much interest in low-lying 1^+ levels in magic nuclei, first of all for ^{208}Pb [3], not to mention the $M1$ resonances in them [4,5].

In spherical open shell nuclei the number of low-lying 1^+ levels studied is increasing rather quickly. There are, for example, experimental data about one or two states in ^{44}Ca , ^{58}Ni , ^{82}Kr and ^{88}Sr . In ref. [2] the energies of six 1^+ levels in ^{96}Zr in the narrow interval (3602—4131) keV have been measured as well as their $B(M1, 0^+ \rightarrow 1^+)$ values. In our opinion, measurements of the spectrum of low-lying 1^+ levels are rather important.

Here we would like to pay attention to the analogy in the microscopic nature between low-lying 1^+ levels in deformed nuclei and in spherical open shell nuclei. In a microscopic approach they should be considered within the QRPA, which takes into account pairing, or within a model using more complex configurations than those being accounted for in RPA, first of all $2p2h$ configurations. One can think that the description of low-lying 1^+ levels in spherical nuclei is simpler than in deformed nuclei. Maybe even the simple comparison of the number of calculated 1^+ levels with the

experiment will help to choose the correct version of the theory.

We have calculated the low-lying 1^+ levels in ^{96}Zr in the framework of:

1. the RPA without pairing,
2. the QRPA [6],
3. a model including $1p1h$ and more complex configuration ($1p1h \otimes \text{phonon}$), but without pairing [7].

The effective particle-hole Migdal's interaction and the local effective charge with the well-known parameters

$$g = 0.05, g' = 0.96, C_0 = 300 \text{ MeV} \cdot \text{fm}^3$$

$$\xi_s^p = \xi_s^n = 0.1, \xi_1^p = \xi_1^n = -0.03$$

have been used in the calculation (see e.g. [5]).

As one can see from the table the RPA gives only two 1^+ levels and does not explain the experiment. Inclusion of pairing (QRPA) improves the result partly (a noticeable part of the strength is shifted to the high-energy region). Taking into account $2p2h$ ($1p1h \otimes \text{phonon}$, to be exact) configurations improves the latter results in the required direction; but still the agreement is not very good, especially for the calculated levels beginning from the third

Experiment		RPA ($\Delta = 0$)		QRPA ($\Delta \neq 0$)		$1p1h \otimes \text{phonon} (\Delta = 0)$	
$E(1^+)$ keV	$B(M1)^\dagger$ μ_o^2	$E(1^+)$	$B(M1)^\dagger$	$E(1^+)$	$B(M1)^\dagger$	$E(1^+)$	$B(M1)^\dagger$
3602(?)	$0.33^{+0.27}_{-0.18} \cdot 10^{-2}$			2261	$0.1 \cdot 10^{-2}$		
3621	$0.51^{+0.69}_{-0.18}$	4131	$0.9 \cdot 10^{-2}$	3772	0.18	3987	1.37
3701	$0.39^{+0.30}_{-0.12}$	4208	2.17	4097	$0.4 \cdot 10^{-2}$	4081	0.75
3947	$0.19^{+0.12}_{-0.06}$			5286	0.48	4303	$0.2 \cdot 10^{-4}$
4038	$0.26^{+0.31}_{-0.11}$			7584	$0.8 \cdot 10^{-2}$	5390	0.07
4132	0.10					5635	$0.4 \cdot 10^{-2}$
						6247	$0.5 \cdot 10^{-3}$
						6779	0.01
						7027	$0.2 \cdot 10^{-4}$
						7616	$0.6 \cdot 10^{-2}$

Table 1: Low-lying $M1$ excitation ($E \leq 8 \text{ MeV}$) in ^{96}Zr .

one. The reason is probably the absence of pairing in our calculation with complex configurations. In any case, the inclusion of pairing should give additional 1^+ levels in the required energy region. One can hope that their $B(M1, 0^+ \rightarrow 1^+)$ values should be more or less comparable with those for the first and the second 1^+ levels. However, so far there is no computational scheme where $1p1h$ and $2p2h$ configurations as well as pairing are taken into account for nonseparable interaction. This work is in progress now.

On the other hand it is very desirable to decrease the experimental errors and to obtain experimental data for the low-lying 1^+ spectrum in other open shell nuclei, e.g. Ni, Sn, Pb.

This work is supported by the German-Soviet exchange program.

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E0 and M1 excitations in ^{56}Ni

E. Migli, S. Drożdż*, J. Speth and J. Wambach

Among the elementary modes of nuclear excitation, the giant resonances are of special importance since they carry valuable information on bulk properties of nuclei and, through extrapolation, of nuclear matter. Of particular interest in this connection is the breathing mode since its excitation energy is directly related to the compression modulus K_A of finite nuclei and, thus, also to the nuclear matter value K_∞ . The most reliable values of K_∞ are extracted from the experimental information involving the giant monopole resonances (GM₀R). It is therefore highly desirable to have a systematics throughout the periodic table. Doubly closed-shell nuclei are especially important because of their simple structure. Some of them, like ^{56}Ni or ^{132}Sn are, however, unstable to electromagnetic or weak decay (^{56}Ni has a half life of 6.1 days) and their measurement is only accessible in inverse kinematics. One purpose of this note is to provide theoretical estimates of the breathing mode characteristics in this nucleus. A second mode of great interest is the magnetic dipole 1^+ excitation which is also accessible in inverse kinematics. The isovector state provides further information on the Landau–Migdal parameter g_0^+ and the quenching of isovector M1 strength in nuclei. The isoscalar state, on the other hand, is related to the isoscalar parameter g_0 which is not well determined so far.

To estimate the total width of the resonance one should solve the full second RPA (SRPA) equations [1] directly in nuclear continuum. This, however, is not possible at present for technical reasons. Furthermore, to achieve convergence of the SRPA solutions with respect to the single-particle basis one necessarily has to use the finite range interactions, for instance those generated by the one-boson exchange models. While such an approach works well in most cases, the centroid energy of the state which is of basic interest here, i.e. the breathing mode, comes out too low by 2–3 MeV in heavier nuclei. Landau–Migdal type zero-range effective interactions, in the framework of phenomenological RPA calculations, are more successful. For these reasons and in order to obtain physically meaningful predictions, we perform the following analysis [2].

First, we solve the continuum RPA equations with the zero-range density dependent interaction [3]. The resulting excitation energy (multiplied by $A^{1/3}$) of the breathing mode as a function of the mass

number is shown in Fig. 1. On average, one finds satisfactory agreement of the calculated results and the experimental data. The predicted centroid energy in ^{56}Ni is about 18 MeV.

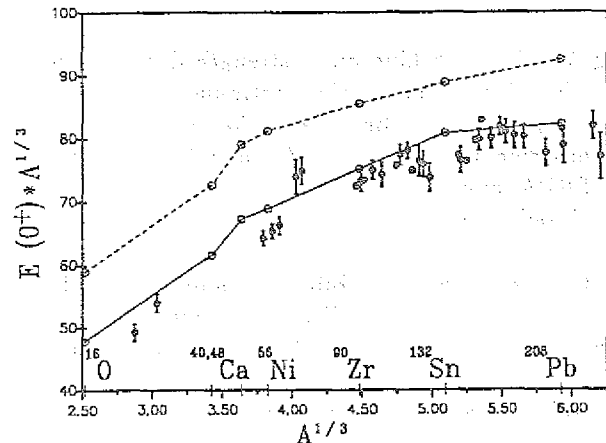


Fig. 1. The excitation energy of the breathing mode multiplied by $A^{1/3}$ as a function of the mass number (A). The solid line corresponds to the continuum RPA results and the dashed line to the free (no interaction) response.

To combine the spreading and escape width we add them incoherently. This means that the underlying subspaces of the Hilbert space are assumed to be orthogonal, which is a reasonable assumption. The effect of the nuclear continuum is to introduce a width on the $1p1h$ level. This effect can be included in calculations with a discrete single-particle basis by taking the ph energies ϵ_{ph} as complex. The accuracy of the procedure can be verified by comparison with exact continuum calculations. Therefore, the most economic way to include both the escape and the spreading width is to solve the SRPA equations with the complex ph energies. Their imaginary part was determined by the free continuum solution for each ph configuration independently. This procedure accounts for an energy dependence in the escape width.

The effect of spreading on the GM₀R in ^{56}Ni is illustrated in Fig. 2. On the level of continuum RPA (dotted line) the resonance has a FWHM of about 2 MeV. Inclusion of $2p2h$ effects with the finite range forces spreads it further (solid line) with a FWHM of 5 MeV. In particular, the distribution develops a high energy tail. As a consequence, about 83% of the total energy weighted sum rule remains in the peak region. Concerning the M1 excitation, in ^{56}Ni we have both $\Delta T = 0$ and $\Delta T = 1$ components since the $1f_{5/2}$

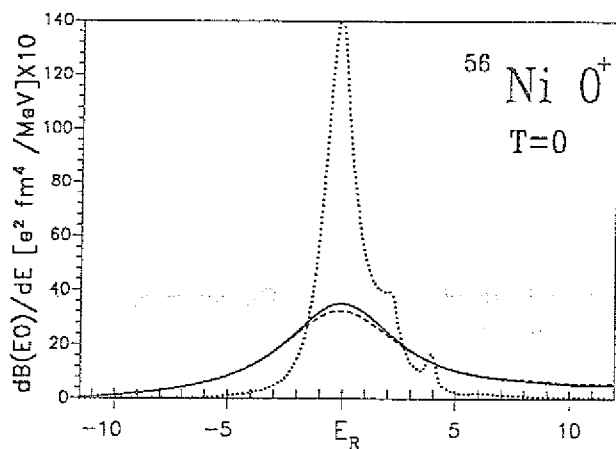


Fig. 2. The breathing mode strength distribution in ^{56}Ni after including the continuum and spreading (SRPA) effects (solid line). For comparison, the dashed line ignores the ladder and bubble diagrams in SRPA and the dotted line includes the escape width only.

proton and neutron orbitals are unoccupied whereas the $1f_{7/2}$ orbitals are filled. The existence of both, isoscalar and isovector M1 transitions is well established in ^{208}Pb and ^{12}C . Their measured centroid energies are correctly reproduced by RPA calculation using the zero-range density dependent interaction.

The M1 strength distribution in ^{208}Pb , ^{56}Ni and ^{48}Ca from our continuum RPA calculations, is shown in Fig. 3. The situation in ^{56}Ni is similar as in ^{208}Pb .

Both, the proton and neutron $1f_{5/2}$ $1f_{7/2}$ configurations contribute to the 1^+ -state. As only the neutron single particle energies are known experimentally, we have chosen the same value for the energy difference of the proton configuration. Therefore, in ^{56}Ni we obtain only one state in the independent particle model at an excitation energy of $E^{1^+} = 6.8$ MeV (middle part of Fig. 3). For the RPA, we predict an isoscalar state at 7.5 MeV with $B(M1\uparrow) = 0.4 \mu_N^2$ and an isovector at 9.6 MeV with $B(M1\uparrow) = 9.5 \mu_N^2$ using the same interaction parameters as in ^{208}Pb . Both states are bound since the continuum threshold in ^{56}Ni starts at 12.2 MeV. The unknown value of the proton spin-orbit splitting introduces some uncertainty of a few hundred keV in the predicted energies. If one increases the proton $1f_{5/2}$ $1f_{7/2}$ ph energy by 400 keV, for instance, the isovector state is shifted up also by the same amount. Also, the transition probability $B(M1\uparrow)$ for the isoscalar state shows quite sizable sensitivity to the proton ph energies. Consequently, its value of $0.4 \mu_N^2$ reflects an anticipated order of magnitude rather than a precise prediction. For completeness, we also show in the lower part of Fig. 3 the M1 distribution in ^{48}Ca . Here only the neutron $1f_{5/2}$ $1f_{7/2}$ configuration can

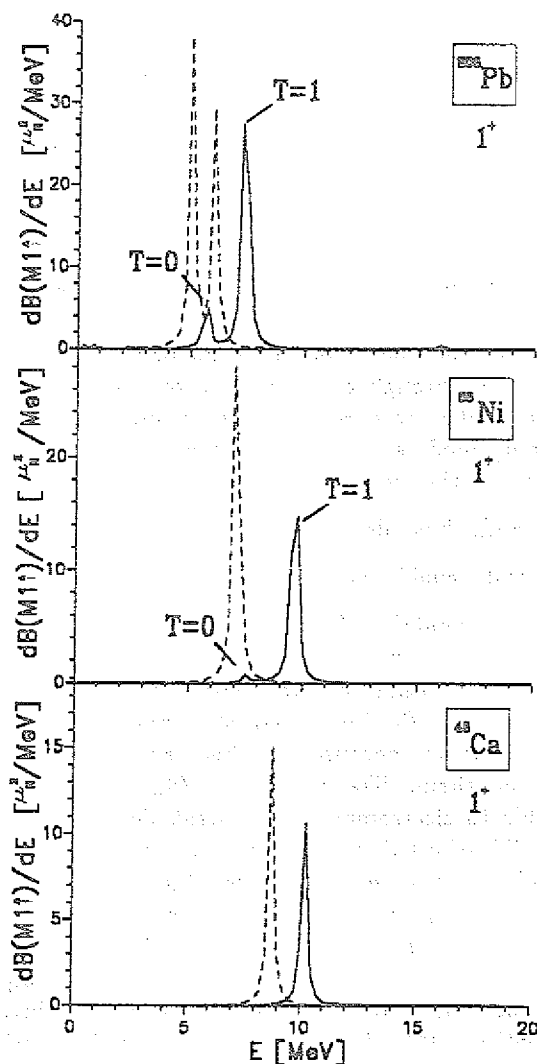


Fig. 3. The M1 strength distribution in ^{208}Pb (a), in ^{56}Ni (b) and ^{48}Ca (c). The solid lines display the RPA results with the density dependent zero-range interaction and the dashed lines denote the free response (no interaction).

contribute and therefore we obtain a single 1^+ -state in the uncorrelated as well as in the RPA case. Its excitation energy of 10.2 MeV and $B(M1\uparrow) = 4.77 \mu_N^2$ agree well with the measurement ($E_x = 10.23$ MeV, $B(M1\uparrow) = 4.0 \pm 0.3 \mu_N^2$). The RPA 1^+ -state is essentially at the same energy as the calculated isovector state of ^{56}Ni .

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Coulomb dissociation of loosely bound neutron rich nuclei and higher order electromagnetic effects

G. Baur and D. Trautmann*

In the Coulomb dissociation reaction of a very loosely bound neutron-rich nucleus



the Coulomb push on the outgoing ^9Li can cause a shift of the final velocity of the ^9Li as compared to that of the neutrons. Such an effect is observed in intermediate energy Coulomb dissociation experiments at MSU [1]. This effect is not explained in approaches where the electromagnetic interaction is treated in first order only. In the post-form DWBA the Coulomb interaction is treated to all orders, and "post-acceleration" effects on (d,p) spectra have been clearly identified [2]. Using this approach, no "post-acceleration" effect on the ^9Li -spectra was observed [3] in a specific calculation at $\theta_{^9\text{Li}} = 0^\circ$. This can be seen in Fig. 1. The

energy distribution of the ^9Li fragment is shown in the forward direction for the reaction $^{11}\text{Li} + ^{197}\text{Au} \rightarrow ^9\text{Li} + ^{197}\text{Au} + 2n$. The cross section peaks around the beam velocity. We study this effect and its dependence on incident beam energy, binding energy and scattering angles in a systematic way.

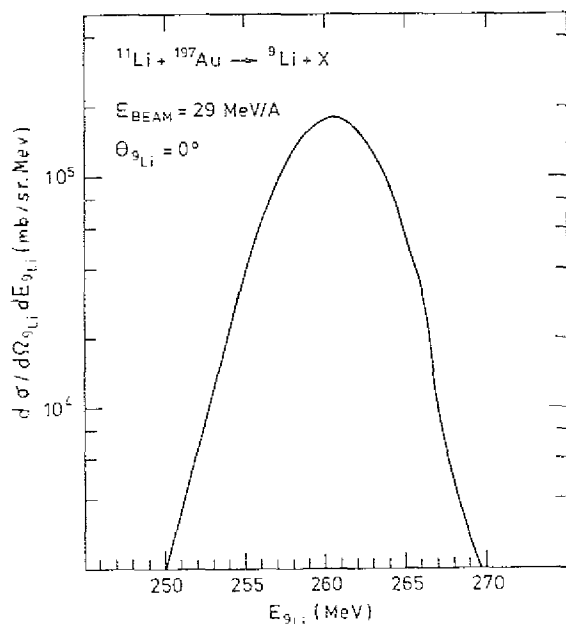


Fig. 1. ^9Li -energy distribution for the emission angle of zero degree in the Coulomb break-up reaction $^{11}\text{Li} + ^{197}\text{Au} \rightarrow ^9\text{Li} + X$ at the beam energy of 29 MeV/A.

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Coulomb dissociation of ^{11}Li and ^{11}Be in a direct fragmentation model

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The Coulomb dissociation contribution to the two neutron and one neutron removal cross section of ^{11}Li and ^{11}Be projectiles respectively has been calculated in a direct fragmentation model [1]. The model is formulated within the framework of the post form distorted-wave Born-approximation. In this theory, the "post acceleration effect" of the ^9Li and ^{10}Be fragments in the final channel is accounted for properly [2]. The Coulomb break-up can explain approximately upto 80% of the experimental $2n$ removal cross-section of ^{11}Li for the ^{197}Au target at the incident energy of 29 MeV/A. The beam energy dependence of the calculated cross sections is in good agreement with the few experimental points which are available upto now. The angular distributions of the ^9Li as well as of $2n$ system is found to be concentrated in the very forward angles.

In Fig. 1 the angular distributions of the ^9Li fragment are shown. The slope of the distributions changes with the beam energy. It is also noted that for the same beam energy the angular distribution of the $2n$ system is flatter as compared to that of the ^9Li . This is because of the lower mass of the former fragment in comparison to that of the latter.

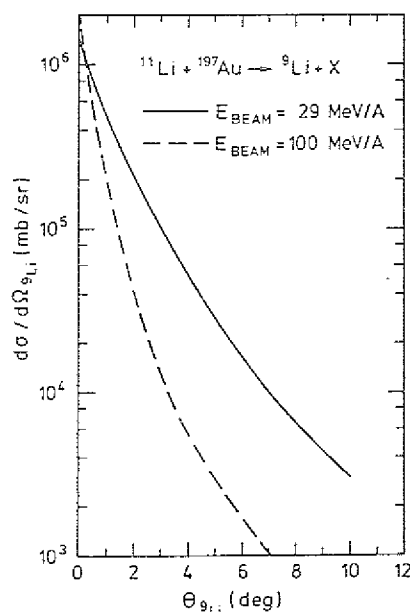


Fig. 1. Angular distribution of the ^9Li fragment in the Coulomb breakup of ^{11}Li .

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The Coulomb field of heavy nuclei provides an intense source of quasi-real photons acting on passing nuclear projectiles. This fact has many applications in nuclear structure and particle physics. In 1984, Coulomb dissociation (or, using an equivalent term, electromagnetic dissociation) of fast projectiles was proposed [1] as a method to investigate radiative capture processes (time-reversed to photodisintegration). This is a new access to a class of radiative capture reactions at low energies, relevant for nuclear astrophysics (for a recent general review see [2]).

In the meantime, many interesting results have been obtained (a recent mini-review is given in [3]). Let us mention the ${}^6\text{Li} \rightarrow \alpha + d$ Coulomb dissociation, leading to the 3^+ resonant state [4] and to the very low energy α -d-continuum [5].

As an application to radioactive beams we mention the ${}^{14}\text{O} \rightarrow p + {}^{13}\text{N}$ Coulomb dissociation (on ${}^{208}\text{Pb}$). The radiative width Γ_γ of the lowest 1^- state is determined. The ${}^{13}\text{N}(p, \gamma){}^{14}\text{O}$ -reaction rate relevant for the hot CNO-cycle is given in terms of this width. The Coulomb dissociation results, $\Gamma_\gamma = 3.1 \pm 0.6$ eV [6] and $\Gamma_\gamma = 2.4 \pm 0.9$ eV [7] are in agreement with the direct measurement, $\Gamma_\gamma = 3.8 \pm 1.2$ eV [8].

For certain cases, Coulomb final state interaction effects ("post-acceleration") can complicate the analysis. We tackle this problem using Glauber theory.

A related problem arises in atomic physics, in the excitation of atomic autoionized states by ions and the ionization of inner shells by collisions with ions, followed by Auger decay. Under the influence of PCI (post-collision interaction), the spectra of the reaction products are distorted. An elegant solution to these atomic physics problems is given in [9].

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Semiclassical and Glauber theories of electromagnetic excitation with high energy heavy ions

G. Baur

A general framework for the study of multiphoton exchange processes in high energy heavy ion electromagnetic excitation is given [1]. By rewriting the Glauber amplitude in a form appropriate for the long range electromagnetic interaction, a connection to the semiclassical coupled channels approach is established in the nonrelativistic as well as the relativistic cases.

Application to two-phonon-giant-dipole resonance excitation in ^{136}Xe (see Ref. 2) is in progress [3].

Quite recently, a broader structure at near twice the energy of the one phonon GDR is observed in the γ - γ coincidence sum energy spectrum, and is identified with the double γ -decay of the two phonon GDR [4]. The γ - γ decay of the two-phonon GDR was first studied theoretically in Ref. 5, from where Fig. 1 is taken.

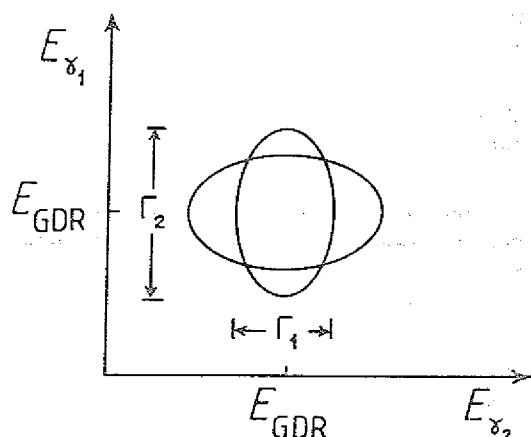


Fig. 1. Correlations of the two γ -energies in the two-photon decay of the $N = 2$ GDR. The widths Γ_1 and Γ_2 of the $N=1 \rightarrow N=0$ and the $N=2 + N=1$ lines are also indicated.

Since the importance of higher order effects decrease with increasing projectile energy, first order (one photon exchange) theories are sufficient in many cases. Higher order effects can conveniently be studied by perturbation theory, instead of using the more elaborate coupled channels approach. This is done now with a group from Milan [6]. Second order amplitudes are calculated in terms of 1st order ones, following the methods developed for non-relativistic Coulomb excitation.

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R. Henneck², M. Jaskola⁴, H. Mühry², I. Sick²

In order to extract the effect of vacuum polarization from ^{12}C — ^{12}C elastic scattering data [1,2], a detailed theoretical study of the low-energy Mott cross section is performed [3]. It is shown that the contributions of nuclear interaction, Coulomb excitation, bremsstrahlung can be neglected, while radiative corrections and nuclear polarizability can be described by a small additional potential. Screening by atomic electrons is corrected by a screening function, which acts on all potentials. Relativistic effects are accounted for using the "Todorov equation". In order to overcome numerical difficulties for very long range potentials a WKB approximation and a semiquantal approach are discussed and compared. The study shows that the first-order vacuum-polarization potential contributes most to the correction of the cross section. All other contributions are at least one order of magnitude smaller. This can be seen in Fig. 1.

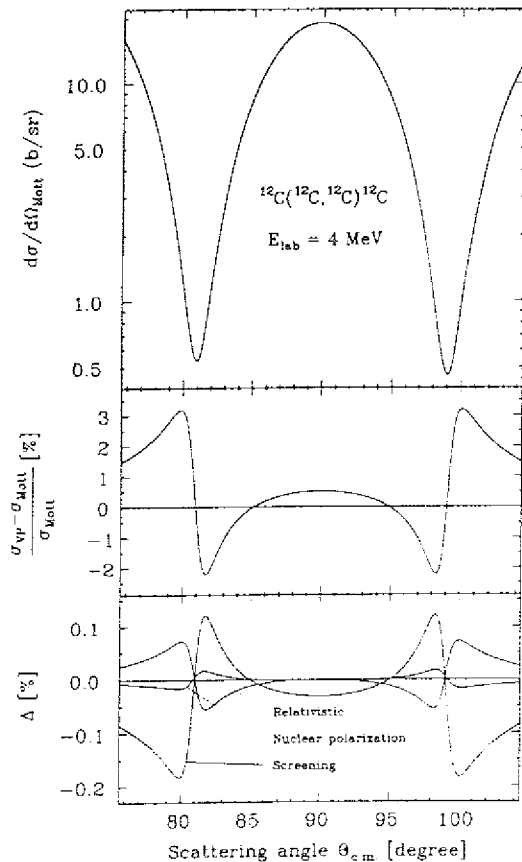


Fig. 1. Angular distribution of the differential cross section (upper part), the relative change of the cross section due to the contribution of vacuum polarization (central part) and additional corrections to the Coulomb interaction.

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III. Accelerator Division

Mathematical Induction

Mathematical induction is a method for proving that a statement is true for all natural numbers. It consists of two main steps: the base case and the inductive step.

Base Case: Prove that the statement is true for the smallest natural number, usually 1.

Inductive Step: Assume the statement is true for a natural number k (the inductive hypothesis). Then, prove that the statement is true for $k+1$.

If both steps are completed successfully, the statement is true for all natural numbers.

Mathematical induction is often used to prove properties of sequences, sums, and divisibility.

Alignment of the COSY-Ring

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T. Pawlat, J. Rook, D. Rosin, Th. Sagefka

Exact positioning of the ion optical elements with respect to the closed orbit is essential for high beam quality. Figure 1 shows the deviation of the beam due to misalignment.

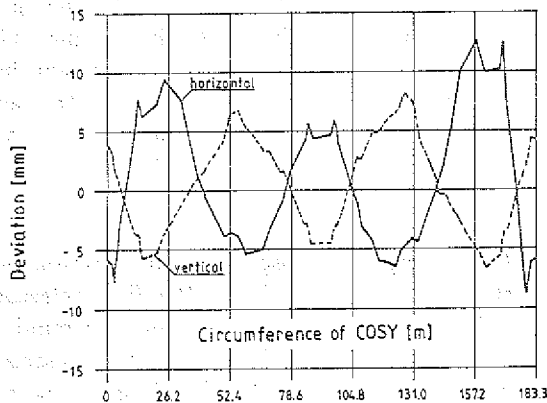


Figure 1: Deviation from central orbit

This simulation assumes a critical deviation of dipoles and quadrupoles by ± 0.2 mm in the vertical plane and a tilt of ± 0.1 mrad for quadrupoles and ± 0.3 mrad for dipoles with respect to the beam axis. Detailed theoretical analysis shows a strong dependence of misalignment of neighbouring components. Thus the alignment tolerances are a factor of 2 to 4 tighter.

With respect to the long ring axis of 70 m the absolute positioning of 0.2 mm is an accuracy of $2 \cdot 10^{-6}$. During measurement the cranes in the accelerator hall have to stay in the parking position due to deformation of the 1 m thick bottom plate. Meteorological influences can deteriorate the measurement, thus all heat sources like heating power electronics have to be switched off.

Due to the weight of building, shielding and (less important) the accelerator differential settlement of the total plant were expected. During the construction of phase several precision nivellements were performed to observe the time dependance of the settlement. The last measurement 14 month after finishing the bottom plate shows that the settlement has decayed to less than 0.5 mm per half year. Figure 2 shows the isolines of the COSY hall.

50 reference points forming a network are the basis for the precise geodetical alignment in the accelerator facility (source injector, COSY-ring, beam lines, experimental areas). Due to the tighter tolerances in the COSY ring, steel pillows are installed as refer-

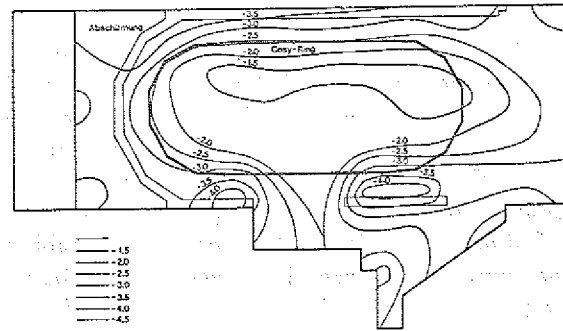


Figure 2: Differential setting of COSY-hall

ence and measurement points. All target points on the components and the reference points are in same observation plane 2.8 m above the floor. The pillows are isolated to prevent bending due to asymmetric heat radiation (cold wall, hot magnet-coils).

Measurement of position and nivellement are completely decoupled for better accuracy and ease in measurement. Figure 3 shows the distribution of the reference points in the COSY-ring. The uncertainty ellipse for each point with a maximum of 0.05 mm for point 1006. With this concept an overall alignment accuracy of better than 0.2 mm and neighbouring tolerance of 0.1 mm can be achieved.

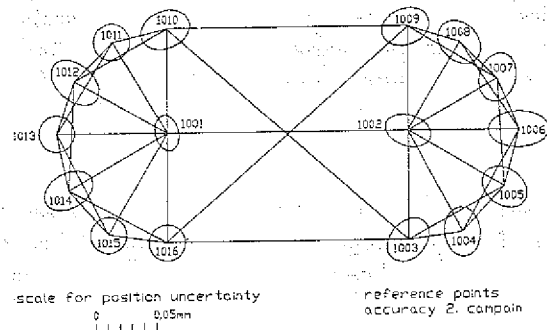


Figure 3: Uncertainty of reference points

All measurement data is stored in a computer and checked directly for confidence and plausibility resulting in short measurement campaigns avoiding time consuming remeasurement. Comparison of different measurement analysis and campaigns shows trends in misalignment and weak points in total plant. In 1991 all dipole magnets were installed and aligned. Sextupoles and parts of the quadrupoles were peraligned up to ± 1 mm. Detailed precision measurements for diagnostic positioning were performed at the injector cyclotron. Alignment of the ring vacuum-chamber and extraction beamline supports were completed.

Assembling of the COSY-Vacuum System

L. Conin, V. Vau

The COSY-vacuum system is specified to operate at an average pressure below $1 \cdot 10^{-10}$ mbar. To match this demand without delay in the assembling period, every component exposed to vacuum is given a final vacuum test. This rather expensive vacuum tests include a base pressure test to the 10^{-12} mbar range for sputter ion pumps, sublimations pumps and beam line vacuum chambers. A base pressure test to the 10^{-11} mbar range is given to valves and gauges. At this low pressures all parts had additional an integral leak test.

Up to now we checked approximately 500 vacuum parts that means ca. 90% of the whole vacuum equipment. In general the defect items are below 5% of the total. Defect rates above the average have been found at the titanium sublimation pump cartridges with 6% and at the vacuum valves with remarkable 25% faults. Keeping in mind that all the vacuum equipment had already a final check at the manufacturer the rate of faults seems rather high. To our experience it pays in time and manpower to maintain an own vacuum test area instead of using the accelerator ring for the final test.

The COSY vacuum ring is parted in eight vacuum sections by gate valves in a manner that each telescope sections and each bending section is divided into two parts. Each section performs an independ vacuum system with an own bakeout control, a roughing system and a gas analysing system.

In the beginning of the year we started to assemble the first telescope vacuum section. This start position was chosen to gather experience on the straight section which is the easiest part to handle

as far as alignment is concerned. During the assembling of the vacuum line the quadrupole magnets have not been installed due to delivery delay. After minor changes of the undercarriage of the quadrupole magnets the latter have been set later into position without disassembling the beam line. The vacuum system of the first telescope section has been leak tight right away at the first try and remained leak tight after minute alignment movements, necessary after putting the quadrupole magnets in position.

Then the bending section opposite the injection side was installed and connected to the first telescope section. The most expenditure of time was required to adjust the vacuum chambers on its undercarriage which are partly attached to the dipole magnets. Several minor changes of the undercarriage and the vacuum chambers had to be done to get a satisfying alignment of the bending section. It is now completely assembled, leak checked and ready for the system bakeout.

At the end of the year 1991 both telescope sections and one bending section are assembled, including all sputter ion pumps, sublimation pumps and roughing pumps. All according thermocouples and heater jackets are complete connected. This part of the ring is kept under vacuum, the bakeout is in preparing.

In parallel to the COSY-ring, the injection beam line has been completely assembled.

In the beginning of 1992 all quadrupole and steerer magnets will be set into position and the assembling of the last bending section will be completed.

Status of the COSY Electron Cooler

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U. Schwarz*, M. Sedlacek**, H.J. Stein, J. Wimmer***, J.-D. Witt, A. Zumloh

In 1991 we were able to start installing the first hardware for the COSY electron cooler. The main magnets, solenoids and toroids were made by SCANDITRONIX-MAGNET and have been delivered on July 1. Before, the electric power cables connecting the cooler magnets with the power supply units in the power supply compartment were prepared in due time. Fig. 1 shows the status in beginning of August, the solenoids and toroids integrated in the support frame. After the installation of the cooling water systems start-up of the magnets with full power occurred in November. The essential investigation of the longitudinal magnetic field including the adjustment of the longitudinal correction coils could be performed until the end of the year. Results of these field measurements are given below. The progress in designing and manufacturing the other components for the cooler systems is briefly described.

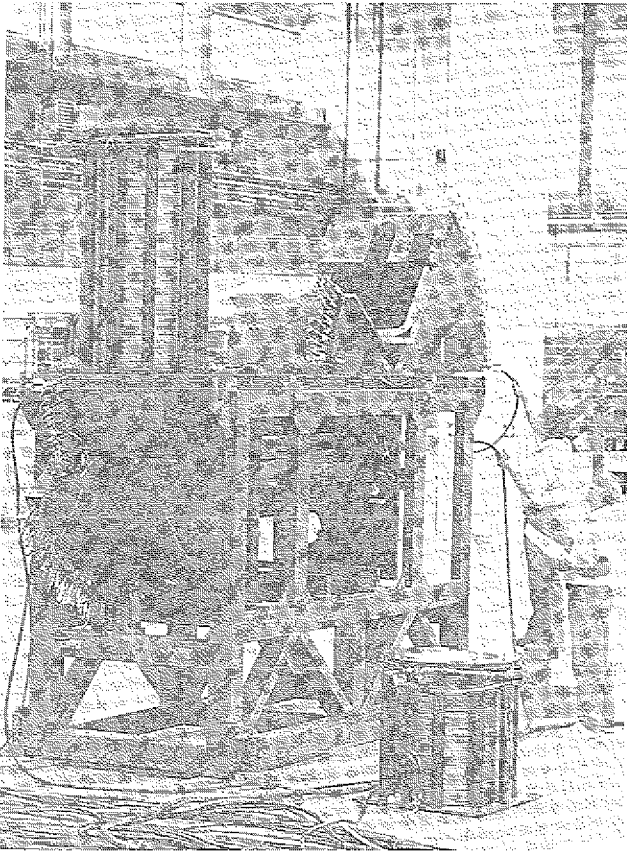


Fig. 1: Main magnets, solenoids and toroids, mounted onto the support frame, status August 1991.

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Gun and collector

The COSY electron cooler has been designed for a maximum electron energy of 100 keV and a maximum current of 4 A in a 25,4 mm diameter electron beam. The acceleration structure consists of acceleration tubes from NEC¹⁾. The company NTG²⁾ has completely designed gun and collector and is also manufacturing and assembling these units until March 1992. The cathode will be delivered by SPECTRA MAT³⁾.

Vacuum system

All components have been delivered and are under test in a separate test set-up. The electron cooler will form a complete vacuum section in the ring, and is, therefore, equipped with turbopump set, ion getter, Ti evaporation, and NEG pumps, as well as vacuum control instrumentation.

High voltage system

The building structure for the Faraday cage has been completed, Fig. 2. Due to local building regulations the structure had to be made earthquake proof. For this reason the 100 kV high voltage platform will also be mounted on a stable support made of special wood instead of power line insulators. The position of the Faraday cage was chosen such that it serves for the test position and the ring position of the electron cooler as well. Completion of the high voltage system is scheduled for April 1992 in order to start first electron beam experiments in May.

Proton beam correction magnets

The PBCMs have been designed and are being manufactured by SCANDITRONIX-MAGNET. The compensation solenoids are of standard design. The steerers on each side for compensating the dipole component of the toroids are integrated in one magnet with an X and a Y coil. Special care was taken to reduce the sextupole component of the Y coil by splitting the winding in two parts⁴⁾.

Control

We have access to the overall COSY control systems via a workstation at the moment placed nearby the test site. This work station is located on the lowest command level. VME and VXI stations, e.g. at the power supply sites or nearby the position pick-ups are loaded on request by a computer one level above our workstation. At the moment only the main power supply is under this control. In May we will have full control on all components, well-timed for start-up of the electron beam.

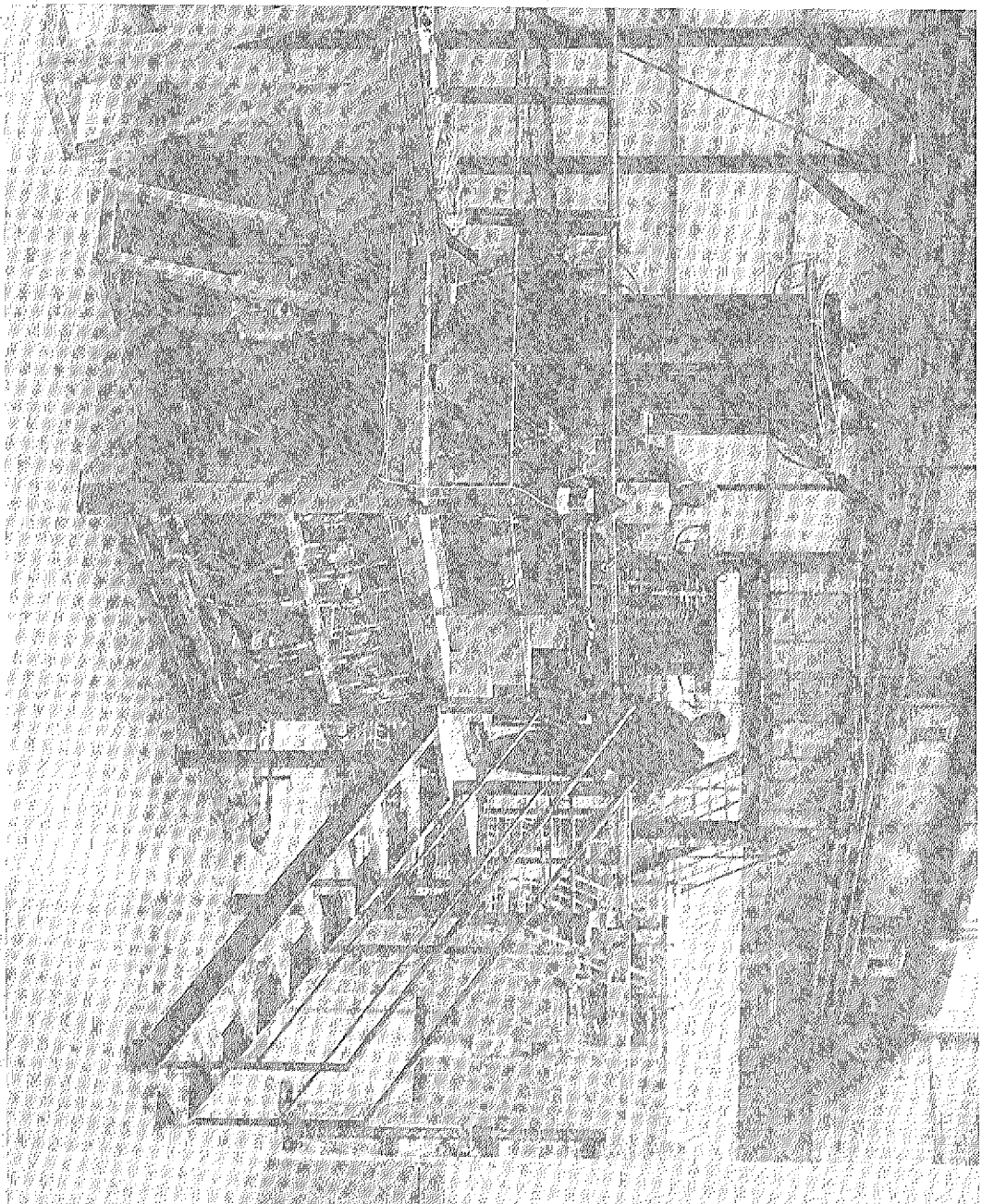


Fig. 2: The COSY Electron Cooler at its test site in the free space inside the ring. All magnetic field measurements are being done in this position. Later, the cooler will be completely assembled here and then lifted into the ring position as a whole unit. The Faraday cage behind containing the high voltage supply systems serves for the test bed and the ring position as well.

Longitudinal magnetic field measurements

Due to design principles which are not explained here, the currents for producing the same field is higher in the toroids than in the solenoids. Since we have all main magnets in series supplied by one main power supply, active shunts for reducing the currents in the solenoids have been chosen⁵⁾. The first reason for longitudinal field measurements is, therefore, to adjust the active shunt currents in gun and drift solenoid, and in the two toroids in between. (We will not take care for an equal field in the collector solenoid.) In a second step the currents for the longitudinal correction coils filling up the field depressions at the transitions of gun solenoid, toroids, and drift solenoid had to be determined.

For these longitudinal field measurements we have chosen the simple and elegant method which was also applied for the CELSIUS electron Cooler magnets. A

closed rail of Al joining the center line of solenoid and toroid was mounted into the magnets⁶⁾. Alongitudinal Gaussmeter probe mounted in a "mouse" was drawn through the guide rail with constant velocity by a stepping motor. The measured field values were continuously registered on an X-Y-line recorder.

First measurements were done with 875 A main current corresponding to 1200 Gauss and with/without the shunts adjusted at calculated values. Then five correction coils were included. Alternately the correction coil and the shunt currents were changed to minimize the residual field fluctuations. Thereafter, the longitudinal field was measured at 1100 and at 1350 Gauss, with all correction coils scaled linearly.

The measurements show that two active shunts (for gun and drift solenoid) are sufficient to equalize the

field up to the drift region. The field depression at the transitions toroid/solenoid (80 — 150 Gauss) was only about 70% of the calculated worst case. Within the measured range the correction coils can be set linear with the main current. The field 20 mm off-axis was of same quality as on-axis. The remanent field (going down from 1100 Gauss) was lower than 4 Gauss, in the drift region even lower than 3 Gauss. Comparing the corrected transition regions with (collector side) and without (gun side) iron shunts we conclude that with iron shunts the field increase in the toroid is much higher and higher correction currents are needed. On the other hand the effective cooling length is longer. A final decision will be made after having measured the transverse field component.

References:

- 1) National Electrostatics Corporation, Middleton/Wisconsin, USA
- 2) Neue Technologien GmbH, Gelnhausen-Hailer, Germany
- 3) Spectramat, Watsonville/California, USA
- 4) K. Halbach, S. Martin, G. Müller, U. Sieling, H.J. Stein, publication in preparation
- 5) Main power supply and active shunts made by Földi AG, Dietlikon, Switzerland
- 6) Supplied by M. Sedlacek, Kungl. Tekniska Högskolan Stockholm, Schweden

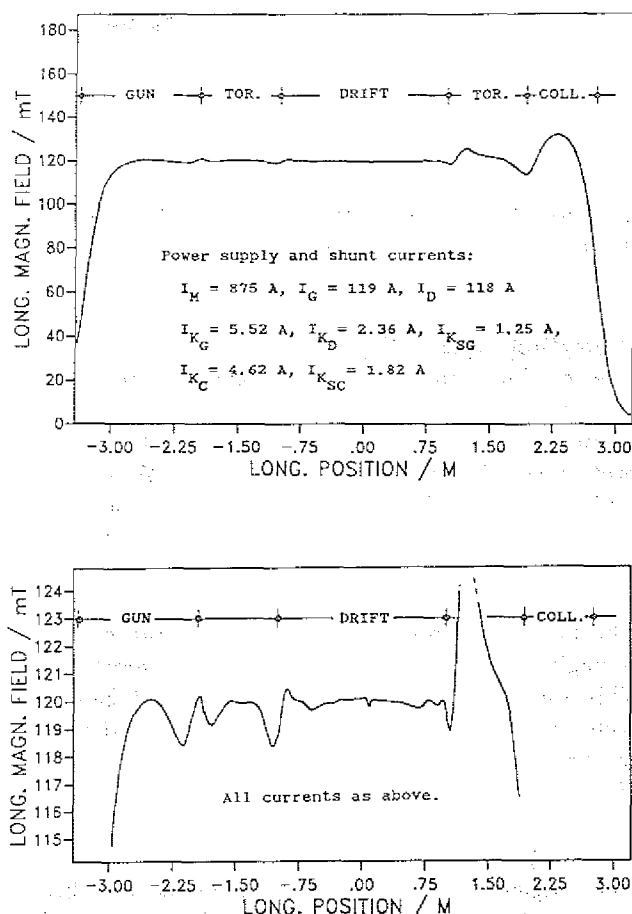


Fig. 3: Longitudinal magnetic field in the COSY electron cooler after adjustment of shunt and correction coil currents.

The Stochastic-Cooling System of COSY

P. Brittner, H.U. Hacker R. Maier, U. Pfister, D. Prasuhn, G. Schug,
H. Singer, W. Spieß, R. Stassen, H. Stockhorst, A. Zumloh

A global description of the system, the locations of the tanks in the COSY ring, and the cooling performance were given earlier^{1,2)}.

Theoretical investigations

The cooling times, the amplification and phase conditions have been recalculated based on the final technical design and the measured data of the prototypes. A microwave power of 30 W per RF feedthrough of the kickers will be sufficient for cooling times less than 30 s for 1.E10 stored protons.

The betatron phase advance over the tank sections is neglectable for the chosen optical parameters¹⁾. This allows the orientation of band-I and band-II sections within the tanks in that way, that identical mechanical RF elements can be used in the PU and kicker tanks.

Cooling tanks

The engineering of the cooling tanks has been completed in October 1991. Fabrication of the tanks will start early in 1992, in collaboration with the KFA central workshop. Table 1 lists the main mechanical tank components. Figure 1 shows a partial view of a vertical stochastic-cooling tank. The thermal parameters of the cooling tanks are given in Table 2.

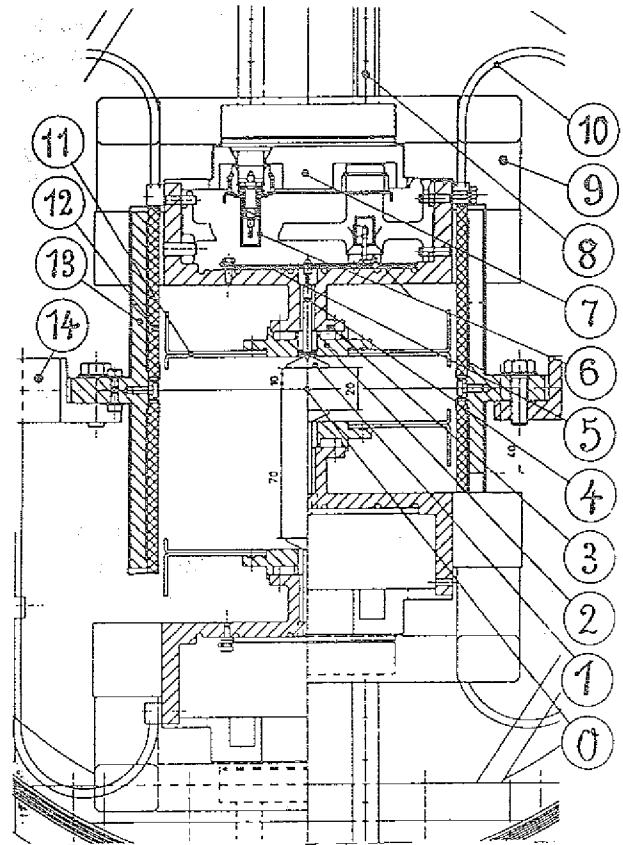


Figure 1:

Stochastic-cooling tank, partial cross-section view around the beam. Three quadrants show the arrangement for the cooled beam; the lower left quadrant indicates the case of uncooled beam.

- 0: beam center,
- 1: stripline beam coupler,
- 2: stripline-coupler fixing plate,
- 3: stripline-coupler support bar,
- 4: 50Ω transmission line,
- 5: power-combiner/splitter board,
- 6: connectors to 5, intermediate lines (rectangular coaxial), connectors to 8,
- 7: traverse for coaxial lines,
- 8: coaxial line connected to movable stripline,
- 9: support-bar pivot connected to actuating rams,
- 10: heat-conducting band (pickup only),
- 11: RF screening sheets,
- 12: ferrite tile for attenuation of waveguide modes,
- 13: ferrite-tile support bar, springs for fixing the tiles,
- 14: fixed ferrite-bar pillar.

Table 1 : Tank components and dimensions

PU kicker		
2	1	tank sections
4	2	beam-coupler support bars
112	56	stripline beam-coupler pairs
3	2	X seals (ø500 * 8)
	2	flange transformers 500 <-> 150
	3	adjustable feet
4	2	Taylor-Hobson sphere supports
beam aperture, V tank /mm ²		140* 20...140*140
beam aperture, H tank /mm ²		20*140...140*140
tank lengths, PU kicker /mm		4355,2 2355,2
tank section length /mm		2050
length b.c. support bar /mm		1950
tank inner diameter /mm		500
beam pipe diameter /mm		150
moving-cycle duration /s		2
expected life time /cycles		10 ⁶
moving tolerance /mm		± 0,1

Table 2 : Thermal parameters of the tanks

temperature of bar, PU Ki	/K	25		300
water-cooling power, Ki	/kW	1		
bakeout temperatur, tank pipe	/°C	100		300
Pickup refrigeration				
temp. of cryo stages, 1 2	/K	25		45
heat input per sect., 1 2	/W	6,5		43
refr.power per sect., 1 2	/W	20		240
cooling-down time	/h	~ 40		

RF-signal processing

The scheme of one signal path is shown in Fig.2. 4 paths are required for 2-band horizontal (H) and vertical (V) operation.

The signal combining and distribution within the cooling tanks will be performed using microstrip lines on 2.5-mm alumina (97.6%) substrates. The strip metallization will be made using thick-film pastes. Ag-Pd will be used for the ground plane in order to get good thermal matching to the alumina sheet. Test boards with ring resonators of different strip widths were measured and lead to a permittivity value of 9.8 and attenuation values around 1 dB/m at room temperature. This value is small enough for both the noise figure in the PUs and the power losses in the kickers. The final layout of the combiner/splitter boards will be adjusted to these measured values and a medium beam velocity of $0.92 \cdot c$. By this, the largest phase error causes an amplitude decrease of less then 0.2 dB between 1.1 and 2.5 GeV and 2 dB at 0.8 GeV.

The electronic components and measured data of prototypes are listed in Table 3.

The variable-delay network has been finally designed. It will enable optimum phase adjustments and amplitude losses of less than 5% for cooling energies between 0.8 and 2.5 GeV using a minimum number of RF switches.

The prototype power amplifiers have been fabricated using Fujitsu FLL100 FETs leading to about 20-W output power. In the next step, a redesign using FLL120 instead of FLL100 will increase the output power to 30 W and smooth the group delay. 2 additional booster modules inside each power amplifier could raise the power if necessary.

Table 3 : Components of RF-signal processing

component	type		number	
preamplifiers	VV I	VV II	32	32
low-power amplifiers	KLV		16	
medium-power amplifiers	MLV		12	
high-power amplifiers	HLV I	HLV II	16	16
program. attenuators	PA		4	
program. delay lines	PV1a,b,PV2,PV3		92	

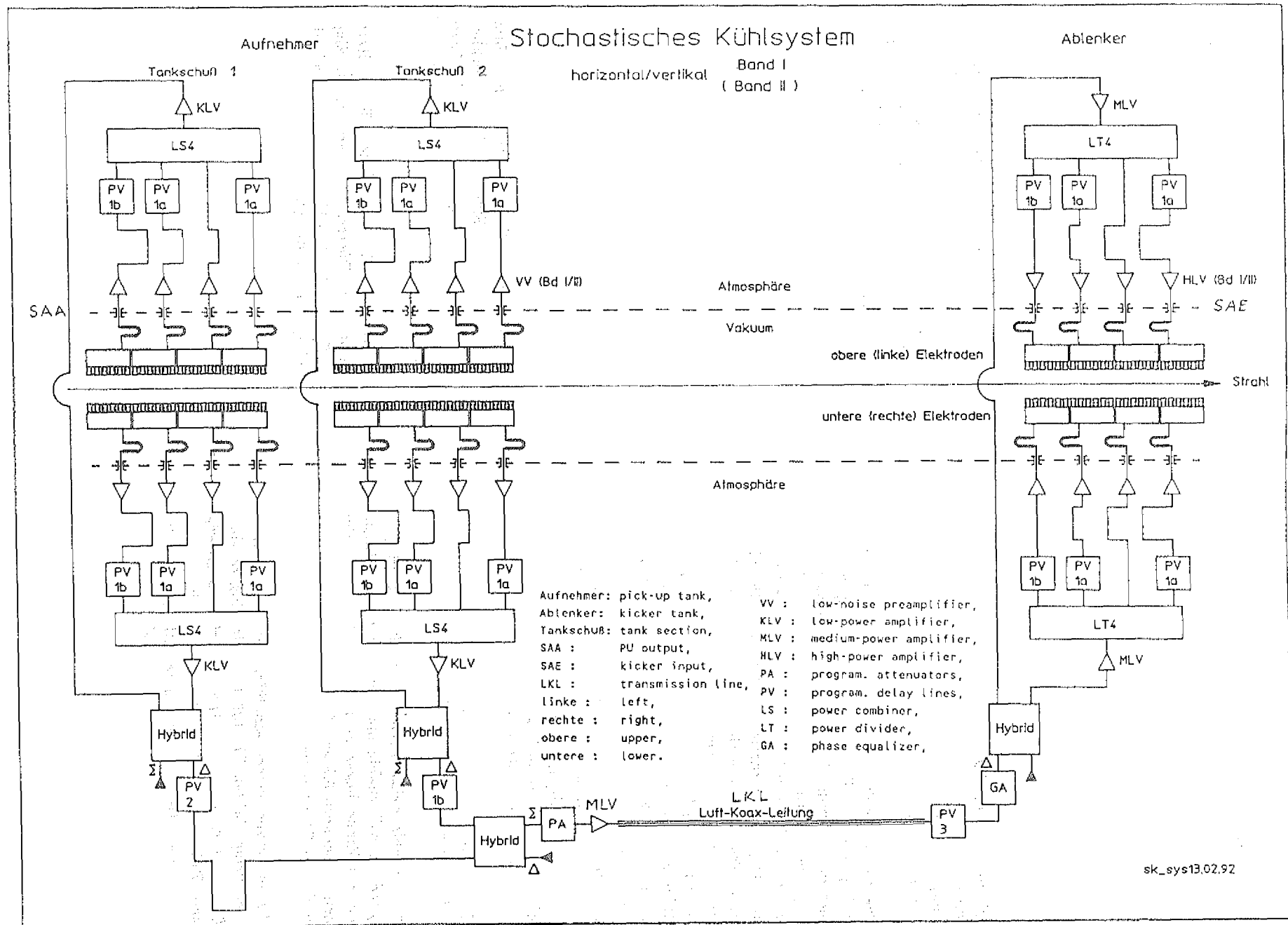
Measured characteristics							
component	VV	VV	KLV	MLV	HLV	HLV	PA
band	I	II	I&II	I&II	I	II	I&II
piece nb.	1	1	3	1	1	1	2
gain /dB	31	29	29,5	33	38	32	4,3
/dB	± 1	$\pm 1,5$	$\pm 0,5$	$\pm 1,1$	$\pm 1,5$	± 2	$\pm 0,5$
S ₁₁ /dB	>10,5	>11	>15	>14	>10	>15	>13
S ₂₂ /dB	>13	>11,5	>12	>11	>16	>15	>13
S ₁₂ /dB	>65	>45	>45	>57	>90	>90	>37
delay/ns	1,0	1,0	0,93	1,0	8,3	7,1	1,42
/ps	± 200	± 200	± 30	± 150	± 1200	± 800	± 55
noise/dB	<0,75	<0,8	<6,7	<5,5	<4,8	<6,3	15
power/dBm	>+6	>+6	>+6	>29,6	+44	+44	>29
alim./V	18-30	18-30	18-30	15-25	18-25	18-25	15-30
/A	0,1	0,1	0,1	0,7	18	18	0,17

Programmable delay lines				
type	PV1a	PV1b	PV2	PV3
nb. of stages	1	2	2	13
el.lengths/mm	20	20, 40	80, 160	2,5..10240
requir.number	48	32	8	4

References:

- 1) R. Danzglock, H.U. Hacker, R. Maier, H. Singer, W. Spiess, H. Stockhorst, D. Prasuhn, K. Bongardt, The Stochastic-Cooling System of COSY, Ann. Rep. 1990, IKP, KFA Jülich
- 2) P. Brittner, R. Danzglock, H.U. Hacker, R. Maier, U. Pfister, D. Prasuhn, H. Singer, W. Spieß, H. Stockhorst, The Stochastic-Cooling System for COSY-Jülich, Part. Acc. Conf. 1991, San Francisco

Figure 2: Scheme of a vertical (horizontal) signal path, only one band is shown



Power Supply Status and First Tests on Original Load

H. Borsch, I. Jannakos, K. Kruck, M. Sauer, J. Schmitz,
H. Schneider, U. Schwarz, W. Schalt, E. Veiders

1. Delivery and mounting:

The power supplies for ring-dipoles, quadrupoles (ring and U-sections), septum magnets and steerers are delivered and placed. For the E-cooler system the setup was made on a testing area next to the final place where the powering could be performed with final power cable length. Nearly all of the power supplies for the extraction beamline are delivered.

The power supplies for sextupoles are ordered and will be delivered in February 1992. Several substation transformers and secondary distribution units are mounted, connected to subsequent units and are tested.

All major cabling for the COSY ring (power cable, interlock and interface cables including cable-tray's are mounted as well.

Ordering, delivering and mounting of crossbar distribution panels to be used with the quadrupole-sets are finished successfully.

A 500 mm² copper bar system throughout the COSY facility is installed in order to achieve a suitable distributed ground potential system by avoiding ground loops. Connections to the dedicated components are made.

2. Measurements of the ring dipole system

The original setup of dipole magnet-system and power supplies for the COSY ring were made and measurements on stability, dynamic behavior and permissible deviation are completed. Proper insulation of the whole dipole powering system was achieved and for earth-leakage detection and indication the connection to a 24 V potential on it was installed.

The dipole magnet power converter consists out of 3 different units (Fig. 1). The first unit is built out of 2 six-pulse thyristor bridges. The 2nd and 3rd unit are forming a very fast parallel current source (called as PE unit) where the 2nd unit consists out of 4 switched fixed current sources and the 3rd unit is an analog transistor controlled current source. This design with a PE unit does not have any tracking error and the time-lag between the set value and the actual value of the current is equal to zero in every instant of time. The actual value is

always within the permitted tolerance band of the width Δ_{xi} surrounding the current set value.

a) Long term stability measurements

The long term stability of the dipole power supply was tested by applying full power over a period of 4 hours. The set value given was 5000 A via remote control. For measuring of the actual current the signal of the internal DCCT was recorded. To control even this device an external DCCT was used as well. Due to the set value of 5000 A the readout was:

actual current value internal DCCT:

$$I_{t=0} = 5000.09 \text{ A} \quad I_{t=0} = 5000.11 \text{ A}$$

actual current value external DCCT:

$$I_{t=0} = 5000.03 \text{ A} \quad I_{t=0} = 5000.06 \text{ A}.$$

Out of this values the drift were calculated to $8 \cdot 10^{-6}$ for the internal DCCT and $6 \cdot 10^{-6}$ for the external DCCT. The absolute current deviation with respect to the set value is $2 \cdot 10^{-5}$. This is quite well within the permissible deviation at 5000 A of $1 \cdot 10^{-4}$.

b) Dynamic measurements on dipole-power supply

The tests were carried out also with connection to the overall control system so the controlling of time dependent measurements (function generated dipole current) could be done. A typical dipole current function over time for COSY was programmed and the deviation of measured dipole current and set value are recorded. Every 40 μs a new set value was applied to the power supply. The final goal is to achieve a time step of 4 μs for updating of the set value in order to meet the envisaged 19 bit resolution.

As a result in Figure 2 one can see the dipole current measured via the analog output signal on the power supply front panel (Channel 1). For better resolution of ripple and stability channel 2 indicates the deviation. It is seen that the deviation does not exceed the range of $\pm 120 \text{ mA}$ with respect to the programmed function. The requirements during ramping the dipole current are indicated in Fig. 2 also. They are within the limits of $8 \cdot 10^{-4}$ during ramping the dipole current. Channel 3 is the output signal of the fast power-unit providing the required tolerances. While switching back and forth

in combination with the analog stage the exact output value will accord to the analog input signal. This gives an good valuation of the self controlled switching of this part.

c) Power supply for E-cooler

After delivering and set up of the power supplies for the electron cooler system, the whole system was tested with the original load of the cooler magnets. Special emphasize was given to the proper alignment of the feedback-loops of the electron-

ically controlled bypass shunt resistors (transistor serial regulated; se Fig. 3). They are used to reduce individual coil-currents. Due to the required range of the magnetic field from 800 G to 1500 G and to the relation between main current and shunt current the linear behaviour was proved (Fig. 4). Exceeding this range will result in a power dissipation for the used transistor array and on the lower limit a nonlinearity behavior of shunt current due to low transistor voltage (U_{CE}). The measured linearity was within 10^{-3} .

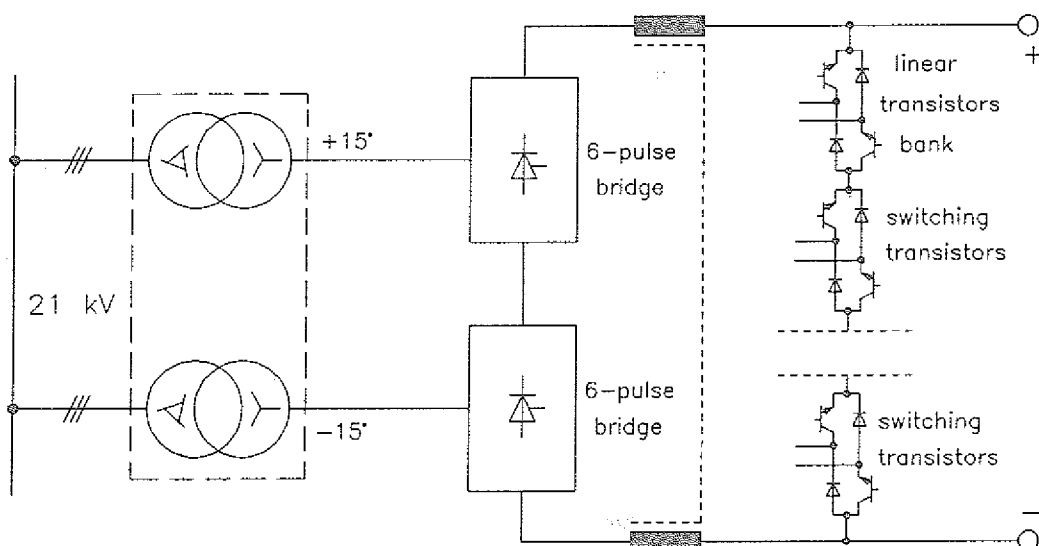
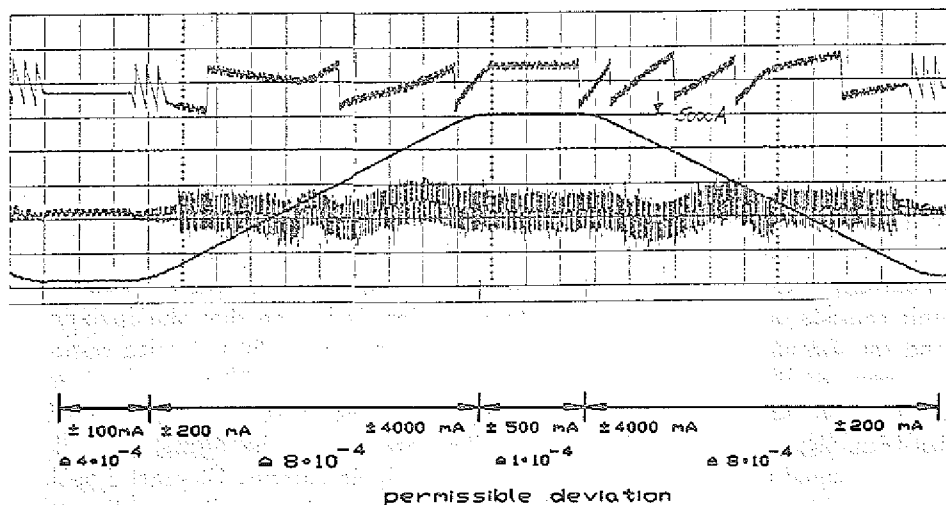


Figure 1: Dipole power supply with PE-unit



DATE: Dec 11/91
TIME: 15:48:24

CH3: 5.00 V
± Wu control of analog stage

CH2: 200 mV
 $\Delta I = 120 \text{ mA/D}$

CH1: 2.00 V
 $I_D = 1000 \text{ A/D}$

$\frac{dI}{dI} = 2500 \text{ A/s}$

ΔI glitchdetect with 30 kHz filter

I_D - magnet current

ΔI - deviation

Figure 2: Measured signals of dipole power supply

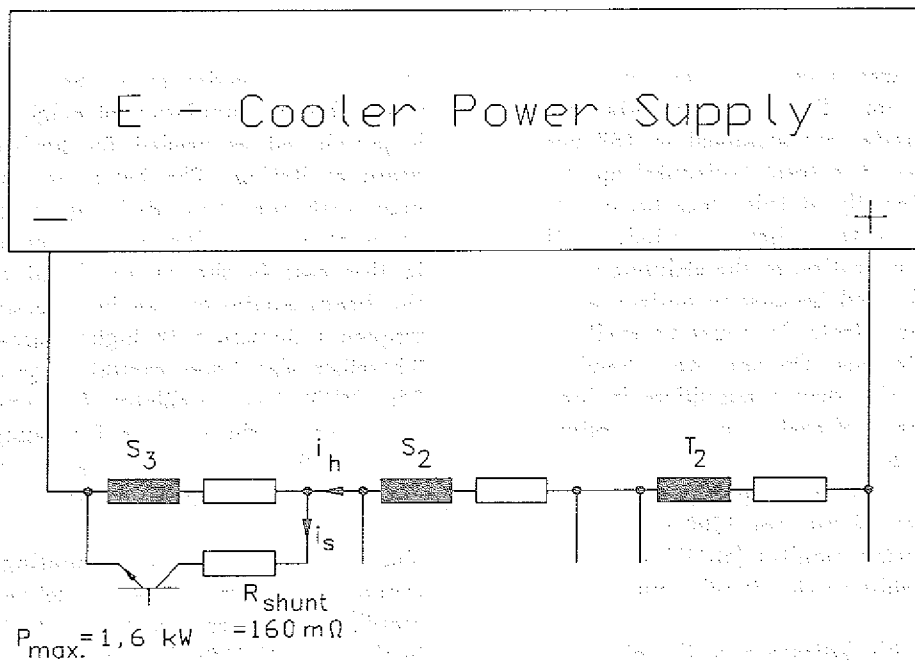


Figure 3: Transistor controlled bypass

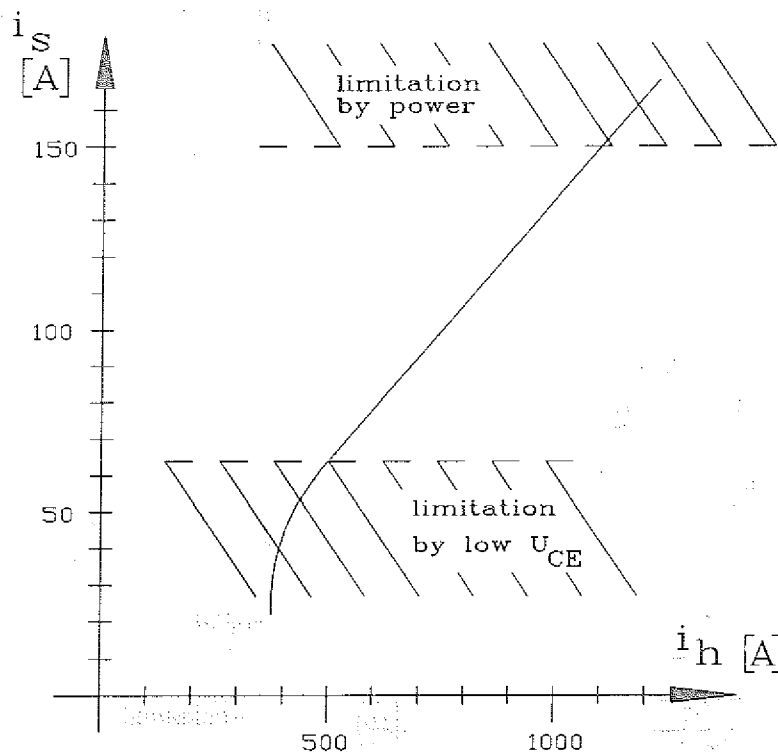


Figure 4: Range of bypass current

Stripline Unit

J. Bojowald, K. Bongardt, J. Dietrich, R. Maier, H. Stockhorst

The stripline unit was meanwhile manufactured. The preparations for UHV-requirements were initiated. The electrodes are arranged in 45° -positions in order to save the total horizontal aperture which is needed especially at this place because the beam is shifted during the injection periode by the bumpers. Thus, the operation of the stripline unit is slightly more complicated because in horizontal and vertical mode all four electrodes must be excited in proper phase-correlations. On the other hand the necessary power of the driver-amplifiers is lower. For 30W output power of each amplifier an estimation¹⁾ yielded a beam deflection of 10^{-7} rad per turn at 2.5 GeV. This corresponds to a beam position shift of about 3 mm for 1000 turns, measured in a beam position monitor (BPM) with about 90° betatron phase shift to the stripline unit.

Figure 1 shows the blockdiagram of the electronics for tune measurements by resonant excitation of beam oscillations using the stripline unit as exciter. The bandwidths of all components allow measurements in the 1-60 MHz frequency range. The signal of the computer controlled synthesizer is properly adjusted in amplitude by the programmable attenuator and distributed by splitters as well as combiners. In this way, two opposite electrodes are stimulated by signals with 180° phase shift. By

means of a coaxial relay the in phase or 180° -phase shifted stimulation of neighbouring electrodes is preselected as needed for horizontal or vertical beam excitation. The beam oscillations are measured with the same BPM-systems as used many times at COSY. The BPM-electronics is operated in this case in the narrow-band mode and hence the beam excitation can be selectively analyzed in frequency domain with higher signal to noise ratio. Therefore also lower excitation power is necessary. The BPM-local oscillator frequency f_{LO} must be correlated to the exciting frequency f_{exc} by $f_{LO} = f_{exc} + 110.7$ MHz. From f_{exc} at resonance the tune can be determined (see Ref. 2).

The components for the stimulating electronics were ordered and already partly delivered. The power amplifiers were ordered with 100 W output power in the 0.1-100 MHz frequency range.

References:

- 1) H. Stockhorst et al., An Analytical Approach to Transverse Tune Measurement by Continuous Wave Excitation, this Annual Report
- 2) J. Bojowald et al., Stripline Unit and Diagnostic Kicker, Annual Report 1990

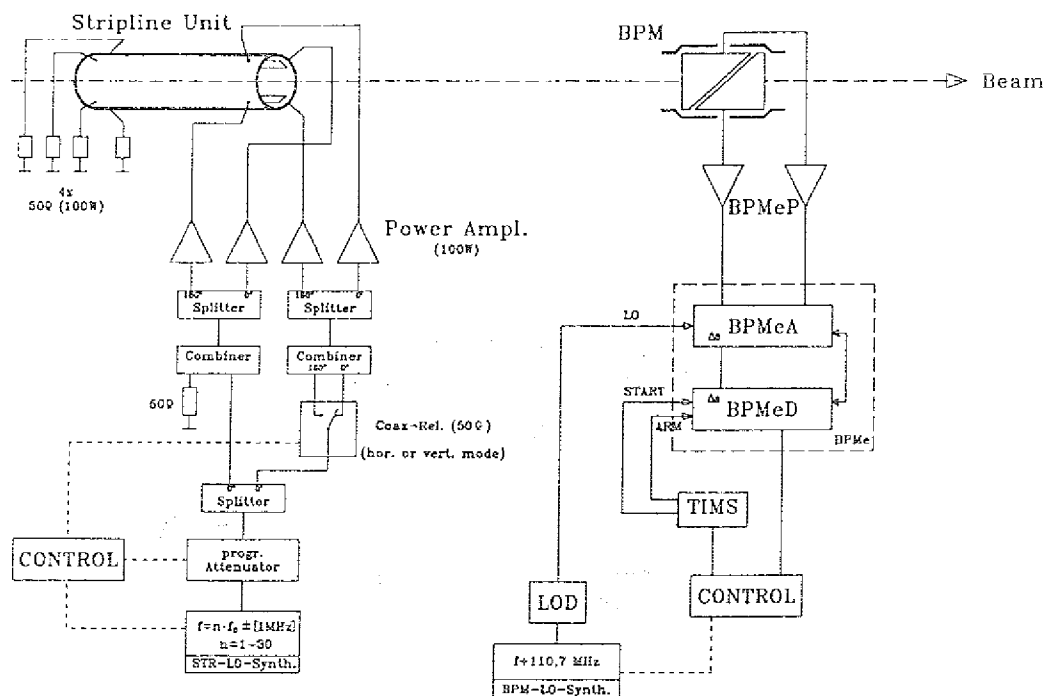


Figure 1:

Electronics for tune measurement by resonant excitation of beam oscillations. The BPM-system for measurement in one plane is shown only.

J. Bojowald, K. Bongardt, J. Dietrich, H. Labus, G. Lürken, R. Maier, H. Stockhorst

Most of the beam position monitors¹⁾ (BPMs) were already installed in COSY. The BPMs were checked by RF-measurements for symmetry of electrode mounting and quality of electrical contacts; evident irregularities were not noticed. Before installation in COSY one BPM was baked out 10-times and checked up after cooling down; no evident alterations could be found.

For measurement of the transverse sensitivity of the BPMs a test equipment using the BPM-prototypes (round and rectangular) was installed. By means of the inner conductor, which was precisely adjusted by mechanical position mechanism and was to simulate the beam and his position, a 2 MHz-signal was injected as bunch current. The quantity difference $\Delta(x)$ over sum Σ of the electrode signals, which is expected to be proportional to the position x (derivation from center position), was measured in dependence of position x . Due to the diagonally slit construction of the BPM-electrodes the expected relation is

$$\frac{\Delta}{\Sigma}(x) = \epsilon_x \cdot (x - x_0). \quad (1)$$

In this equation x_0 is the mechanical offset (for $\frac{\Delta}{\Sigma} = 0$) and ϵ_x the position sensitivity of the BPM Σ which is determined by

$$\epsilon_x = \frac{2}{D} \cdot \frac{1 - \frac{C_c}{C_e + C_c}}{1 + \frac{C_c}{C_e + C_c}} \quad (2)$$

with D = distance of the electrodes (i.e. diameter for the round BPM), C_e = capacitance of BPM-electrode and preamplifier input, C_c = coupling capacitance between both BPM-electrodes. The second factor is a correction term due to the non-vanishing C_c and amounts e.g. to about 0.7 for the round BPM.

The results of the measurement with the round BPM are shown in Fig. 1. The linear relation (1) holds almost over the total BPM-aperture. From this the position sensitivity can be determined to be $9.6 \cdot 10^{-3} \text{ (mm}^{-1}\text{)}$, which is the same in both planes

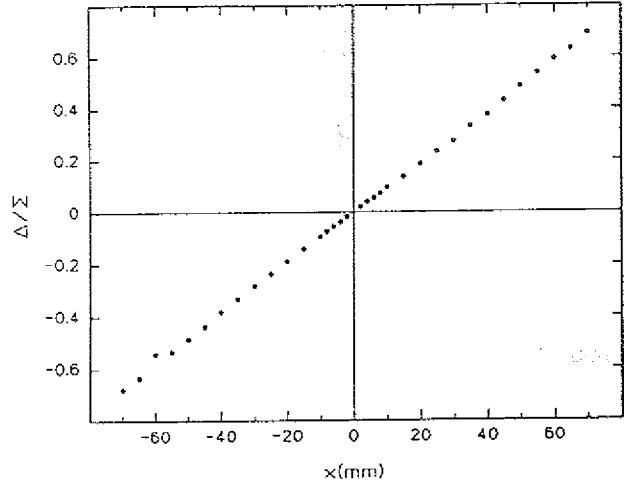


Figure 1:

Position dependence of the quotient $\frac{\Delta}{\Sigma}$ for the round BPM (position x is deviation from center position)

of the round BPM. This value is in good agreement with the theoretical according to equation (2). For the rectangular BPM the position sensitivities are $7.4 \cdot 10^{-3}$ and $19.6 \cdot 10^{-3} \text{ (mm}^{-1}\text{)}$ in the horizontal and vertical plane, respectively.

The amplitude sensitivity or transfer impedance of the round BPM was measured in a 50Ω -matched coaxial system shown in Fig. 2. This system allows to inject a well defined current for beam current simulation in a large frequency range. Figure 3 shows the results of the measurements with high impedance preamplifiers (a) and 50Ω -impedances (b) at the BPM-connectors. In the case of high impedances (i.e. $\geq 500 \text{ k}\Omega$ input impedance as projected for the COSY-BPM preamplifiers) the transfer response is nearly constant in the 0.01-100 MHz frequency range. The transfer impedance Z_{BPM} for one electrode results to 3.2Ω in this range. It corresponds to the transfer impedance of particles with $\beta=1$. This agrees well with the theoretical values of 9.6Ω for $\beta=0.28$ (40 MeV) and 2.8Ω for $\beta=0.96$ (2.5 GeV), which result from the relation

$$Z_{\text{BPM}} = \frac{1}{2C_e \cdot \beta \cdot c} \quad (3)$$

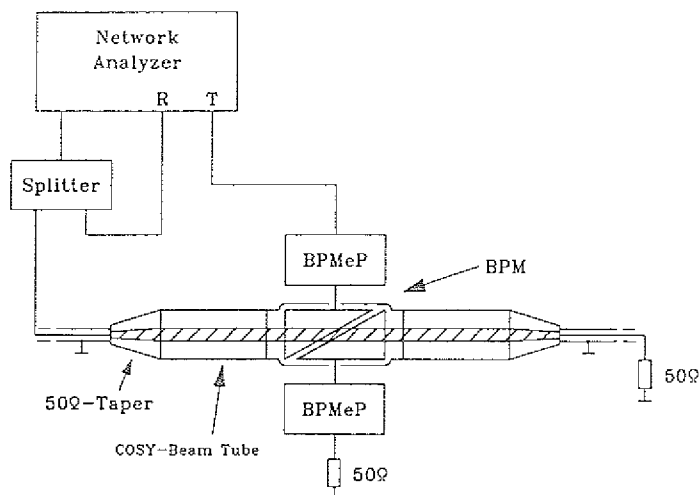


Figure 2:
Teststand for measuring amplitude sensitivity or transfer impedance of the round BPM. The inner conductor diameter is realized to match the 50 Ω -characteristic impedance.

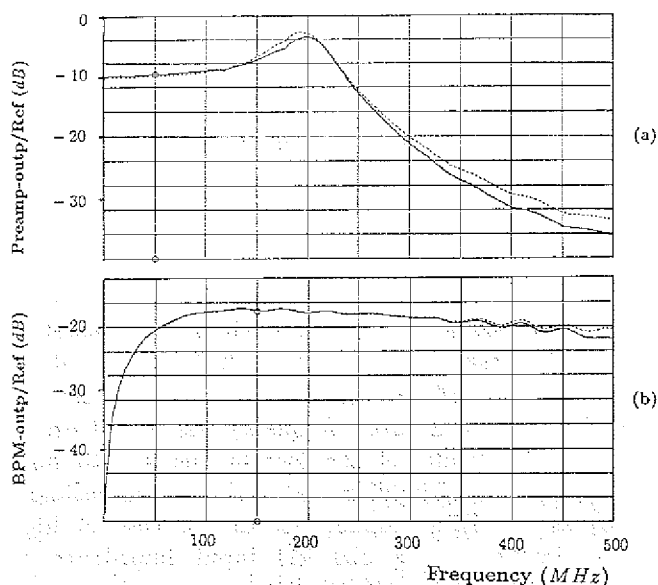


Figure 3:
Frequency dependence of BPM-signals (both electrodes) with high impedance preamplifier (a) and 50 Ω -impedance (b). The ratio of the output signals to the RF-reference signal feeding the inner conductor current is plotted.

with l = length of BPM-electrodes (130 mm) and C_e as above. The factor $1/2$ stems from the geometry factor, because each electrode covers $1/2$ of the azimuthal angle. With 50 Ω -impedances the transfer response is that of a high-pass filter with low frequency cut off (3dB) near 70 MHz. The amplitudes of the output signals in the 150-300 MHz frequency range are to be compared with those of the signals in Fig. 3(a) in the 0.01-100 MHz range; they differ by the gain of the preamplifier and the difference in capacitance C_e .

With the same test equipment the sensitivity of the two BPMs in the e-cooler was measured where 1 m long coax-cables are needed between the electrodes and preamplifiers. The strong resonances due to these cables can be damped by 50 Ω -series resistors as RF-terminators (in connection with the electrode capacitances) inserted at the electrode-end of the cables. Unfortunately the sensitivity is reduced by the capacity of the inserted cables by about 6dB.

In two series 38 BPM-electronics systems (each consisting of 2 matched preamplifiers, analog- and digital module) manufactured in KFKI Budapest were delivered. As a consequence of extensive tests essential improvements are implemented: The preamplifiers were redesigned and exceed the requirements now (e.g. noise $1.5 \text{ nV}/\sqrt{\text{Hz}}$ or $5 \text{ } \mu\text{V}_{\text{rms}}$ for 10 MHz-bandwidth; balanced operation within 0.2 dB and 2° until 100 MHz) and the analog modules were modified.

The distribution of the local-oscillator signals (LOD) for the BPM-electronics at COSY is already installed. Figure 4 shows schematically the arrangements. The individual signals are detached from the branch-lines by directional couplers with -20 dB coupling. By this way a distribution system with passive components only can be realized. But because of the power losses in cables and couplers a separation in two branches is necessary to guarantee the power requirements at the LO-inputs of the BPM-electronics. The same LOD-system as in Fig. 4 is installed twice, identical for horizontal and vertical BPM-systems. Thus position measurements in both planes can be carried out independently and with different harmonics in the frequency spectrum.

Reference:

- 1) J. Bojowald et al., Beam Position Monitors, IKP-Annual Report 1989, p. 189 and R. Maier et al., Non-beam disturbing diagnostics at COSY-Jülich, EPAC 1990

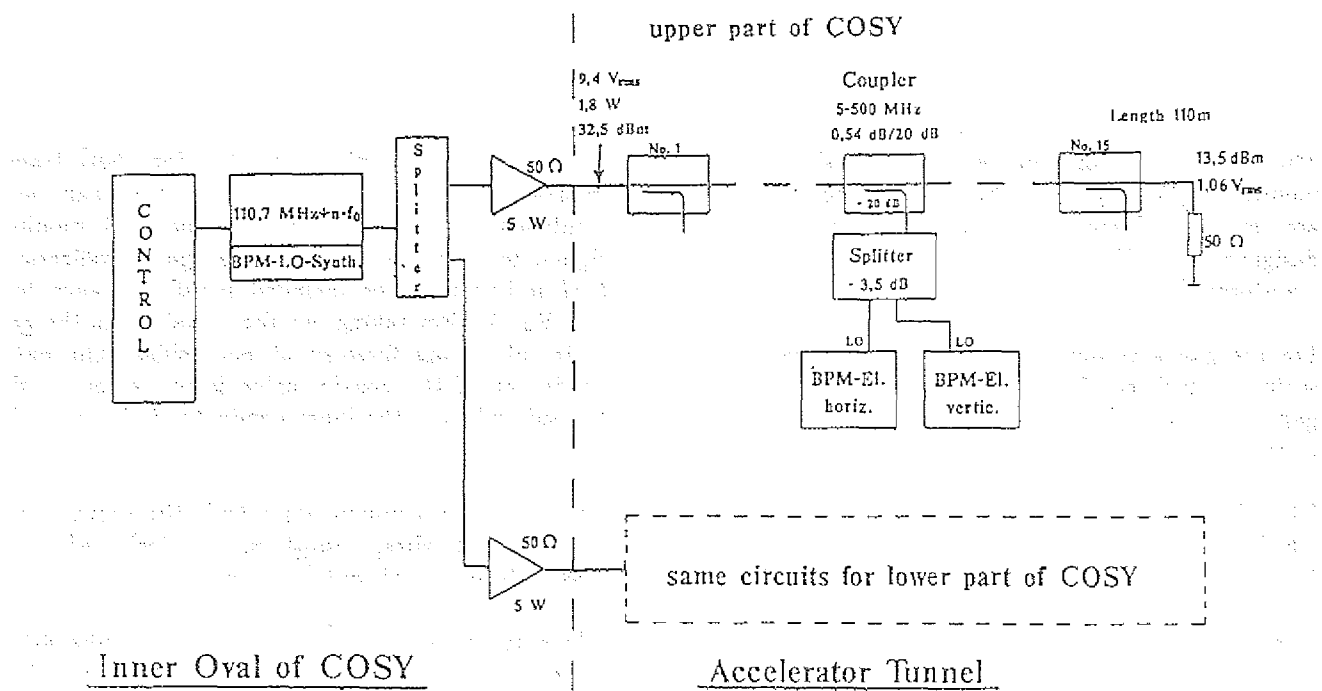


Figure 4:

Local-oscillator distribution (LOD) for the BPM-electronics at COSY. The same system is installed twice, for horizontal and vertical BPM-systems

Wall-Current Monitor

H. Lawin, J. Bojowald, K. Bongardt, J. Dietrich, U. Engels, R. Maier, H. Meuth, U. Probst

The mechanical construction of the wall-current monitor¹⁾ (WCM) is finished and all components are ordered. Identical inner beam tubes were designed for WCM and BCT (beam-current transformer).

For the gap a ceramic UHV-proofed and bakeable sealing is realized. The axial length of the ceramic gap is 20 mm, the radial width 7 mm yielding a capacitance $C_g \leq 60$ pF.

The shielding of the gap against external fields, which also provides a d.c. shunt path, is connected to the beam tube by RF spiral contacts.

The ferrites are located outside the vacuum which are needed to lower the cut-off frequency by their inductance and for damping of waves reflected back to the gap.

The beam signal is picked-up by eight 50Ω transmission lines distributed around the gap and combined to give a position independent monitor signal. In order to be insensitive against reflections 50Ω resistors can be included parallel in each line (see Fig. 1). For taking out the signals from the gap metal pieces are foreseen at both sides. The outer conductor of the coaxial cable is connected to the left side whereas the inner conductor is fixed to the right one.

All mechanical components outside the vacuum (i.e. shielding, ferrites, resistors, cables) will be reassembled for bakeout procedures.

Three types of ferrite materials (Ferroxcube 8C11, 4M2, 4E1 manufactured by Philips) characterized by different initial permeabilities and frequency responses will be applied.

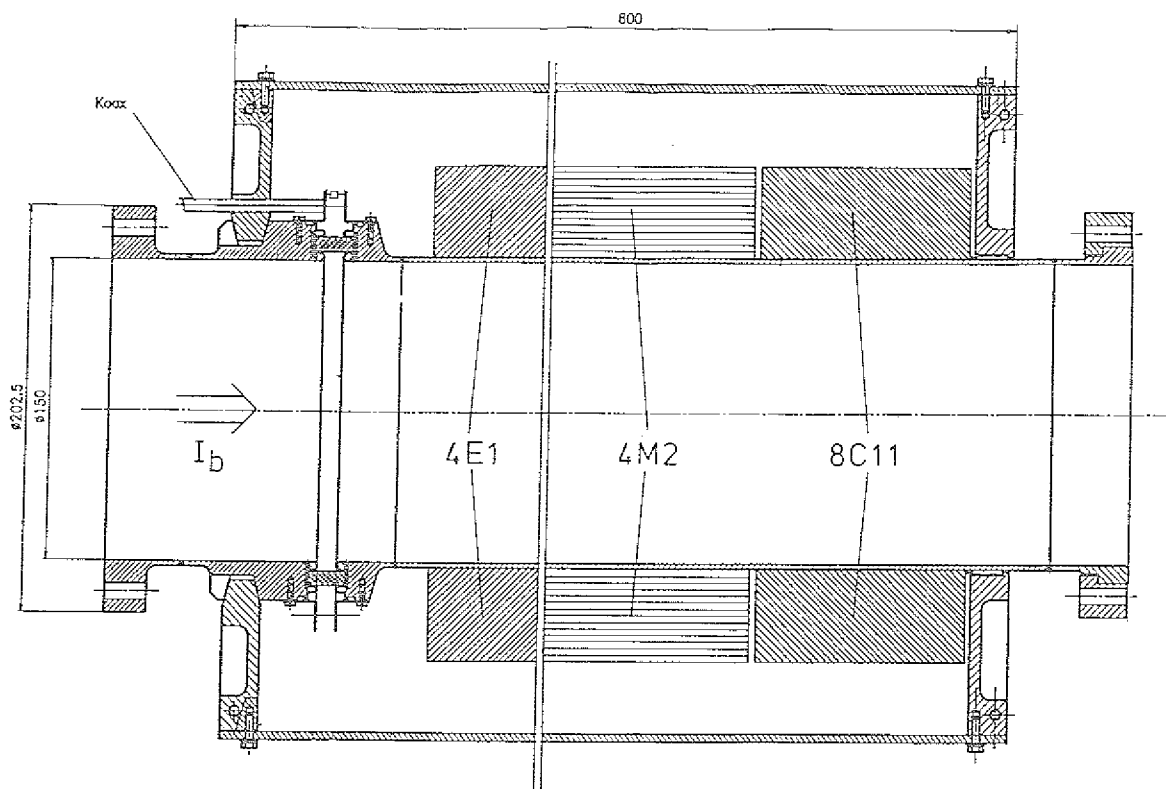


Figure 1: The mechanical lay-out

The Ferroxcube 8C11 with highest permeability ($\mu_{ra} \approx 1400$) mainly produces the inductance at low frequencies whereas the other ferrites provide the damping of RF-resonances by their losses at higher frequencies. Tests to find the optimum arrangement of the ferrites will be performed on a teststand.

The shielding chamber together with the beampipe forms a coaxial line filled with ferrites which has a short circuit at its end. The impedance at the gap $Z_s(\omega)$ is deduced by transformation theory.

Due to the equivalent circuit (Fig. 2) the sensitivity $S(\omega)$ of the monitor is

$$S(\omega) = \frac{1}{Y(\omega)} \quad (1)$$

with
$$Y(\omega) = j\omega C_g + \frac{1}{R_g} + \frac{1}{Z_s(\omega)} \quad (2)$$

whereas $C_g = 60$ pF, $R_g = 3.2 \Omega$ (6.4Ω) with (without) parallel resistors at the gap, respectively.

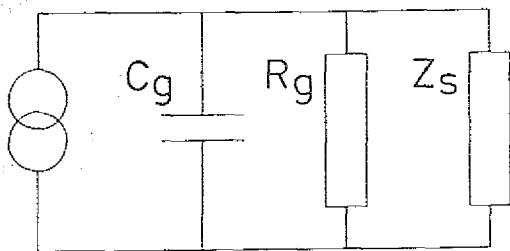


Figure 2:
Equivalent circuit

The signal picked-up by each transmission line is

$$U(\omega) = I_b(\omega) \cdot S(\omega) \quad (3)$$

where $I_b(\omega)$ is the bunch current.

In order to calculate $S(\omega)$ the transformed impedance $Z_s(\omega)$ of the ferrite loaded coaxial line (sump) has to be known.

As an example for 6 ferrite rings of each material (length 12 cm) Z_s is calculated for frequencies in the range 0.1-100 MHz (Table 1).

f(MHz)	$Z_s[\Omega]$
0.1	7.4 j
1	13+99j
10	373-321j
100	22-0.6j

Table 1:
Impedance of the sump Z_s dependent on frequency

For low frequencies Z_s is inductive with an inductance of $L_g = 10 \mu H$.

The imaginary part of Z_s is changing from inductive to capacitive around 5 MHz. At frequencies above 100 MHz Z_s is a complicated function dependent on the mixture of the different ferrite materials.

The lower cut-off frequency of the monitor $f_1 = R_g/(2\pi L_g)$ is 100 kHz ($R_g = 3.2 \Omega$) or 200 kHz ($R_g = 6.4 \Omega$), respectively. The upper cut-off frequency mainly is determined by the gap-capacitance C_g and the resistance R_g resulting in $f_u = 1/(2\pi R_g C_g) = 800$ MHz ($R_g = 3.2 \Omega$) and 400 MHz ($R_g = 6.4 \Omega$), respectively.

As the bunch frequencies are the higher harmonics of the revolution frequency (0.46 MHz at an energy of 40 MeV and 1.56 MHz at 2.5 GeV) in the measurement range of the monitor the relation $|Z_s| \gg R_g$ is fulfilled. Therefore the sensitivity $S(\omega)$ is determined by R_g .

In order to measure a signal independent of beam position the outputs of the eight lines are summed by means of two 4-fold and one 2-fold combiner. The frequency range of the combiners goes from 0.5 up to 100 MHz. The resulting total voltage is

$$U_\Sigma = I_b \cdot S_\Sigma \quad (4)$$

where the sensitivity including 1 dB losses is $S_\Sigma = 7.5 \Omega$ (15Ω) for $R_g = 3.2 \Omega$ (6.4Ω), respectively.

The combined signal is amplified by means of the 100 MHz (≤ 2 dB) preamplifier of the BPM (beam position monitor) electronics.

For detailed investigations of the bunch shape the amplified signal can be analyzed either by a fast oscilloscope and/or by a spectrum/network analyzer. The averaged beam current is deduced by integration of the bunch-pulses after baseline restoration.

The routine measurements within the COSY-control system will be performed by using the standard BPM electronics. At this moment, in time domain the response is limited to 10 MHz (3 dB). In frequency domain the amplitude information can be measured up to 60 MHz, but no phase information is available.

Reference:

- 1) J. Bojowald, H. Labus, H. Lawin, R. Maier, H. Meuth, G. Schug, A. Zumloh, Annual Report 1989, Jül-Spez-562 (1990)

Both Schottky-pickups were already inserted into COSY. Before installation measurements of the pickup-sensitivity (also called transfer impedance) were carried out in a special test equipment shown in Fig. 1. A well defined RF-current is coupled in to an inner conductor (ϕ 1 mm) for simulation of beam current. For that purpose a good RF-adaptation of the coaxial structure and the signal coupling is necessary. At the signal input (top in Fig. 1) this is realized by a matching network transforming the characteristic impedance $Z_0 \approx 200\Omega$ of the structure to 50Ω -input impedance, at the output (bottom) by a symmetrically distributed termination. The quality of the signals in Fig. 2 and 3a demonstrate the sufficiently good adaptation obtained.

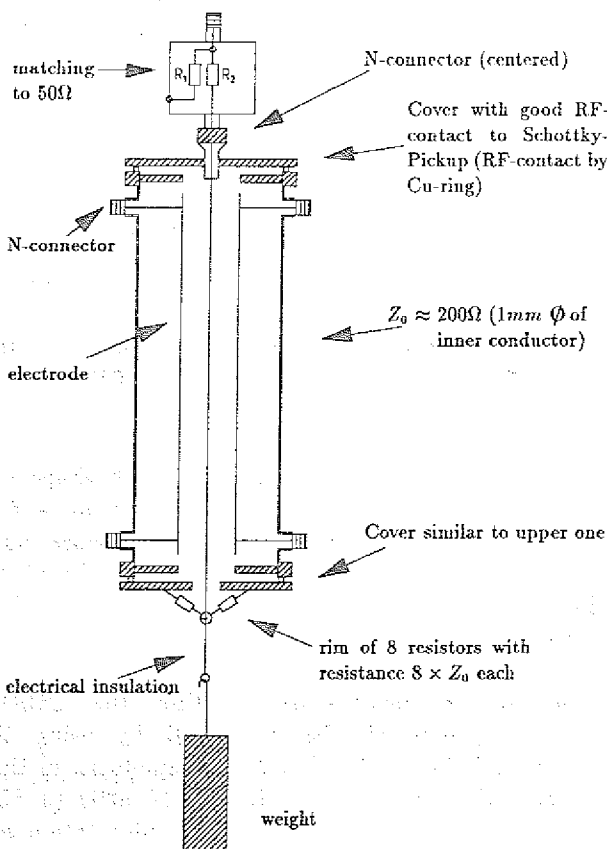


Figure 1:
Test equipment for frequency response measurement of the Schottky-pickups. Signal input is at top.

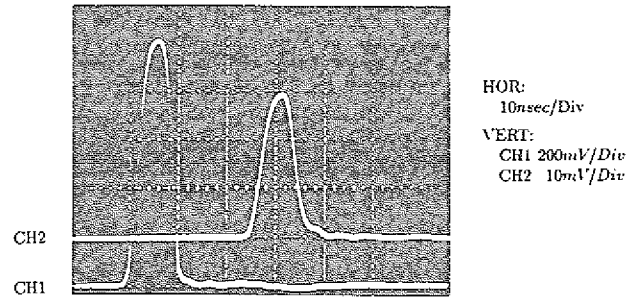


Figure 2:

Transmission of the coaxial structure (quality of RF-adaptation).

CH1: input signal; CH2: output signal, coupled out at bottom. Scale: horiz. 10 nsec/Div, vert. 200 mV/Div.(input) and 10 mV/Div (output).

The output signal is attenuated because of damping in the matching networks (the network for coupling out at the bottom is not shown in Fig. 1 for clearance).

In Fig. 3 the response of the pickup is shown in time (a) and frequency (b) domain. The frequency response of the sensitivity (Fig. 3b) is in good agreement with theoretical predictions¹⁾: amplitude and phase show sinusoidal and linear shape, respectively. The maximum sensitivity of one electrode amounts to 15 mV/mA at 75 MHz (the pickup-length corresponds to $\lambda/4$ at this frequency). Figure 4 shows the sensitivity enhancement of about 10 dB (factor 3) obtained by resonant tuning using coax-cable RG 58 C/U. At the same time the bandwidth is strongly reduced. A larger enhancement of sensitivity is expected if cable with less transmission loss will be used. Further investigations with resonant tuning can be performed also after installation of the pickups in COSY (i.e. without inner conductor or beam) by signal coupling in via the second electrode. Tests showed that in this case the sensitivity enhancement is similar in amplitude and shape.

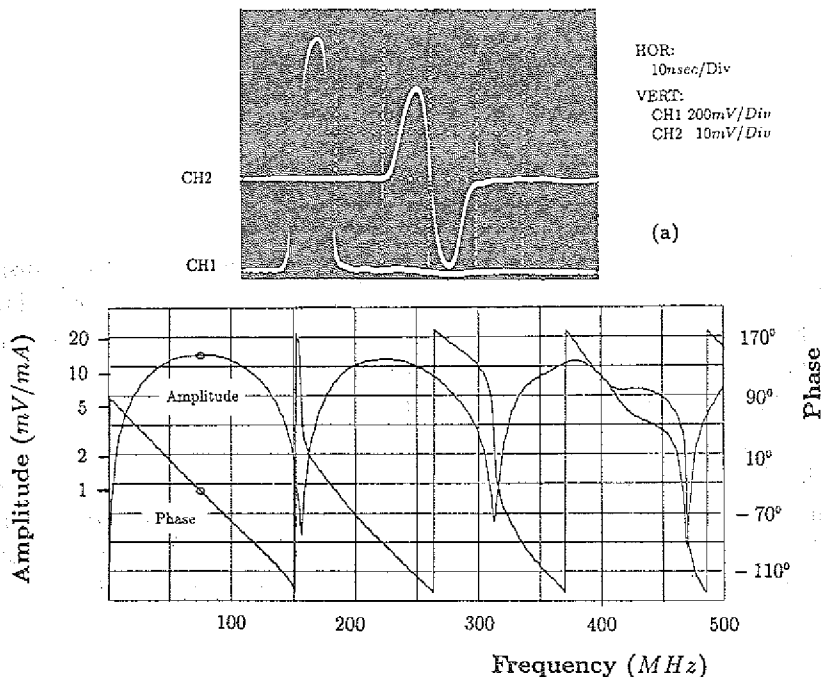


Figure 3: Response of the Schottky-pickup electrode signal in time (a) and frequency (b) domain. Signal at top contact of electrode with bottom contact terminated. Without resonant tuning. Scale in (a) as in Fig. 2.

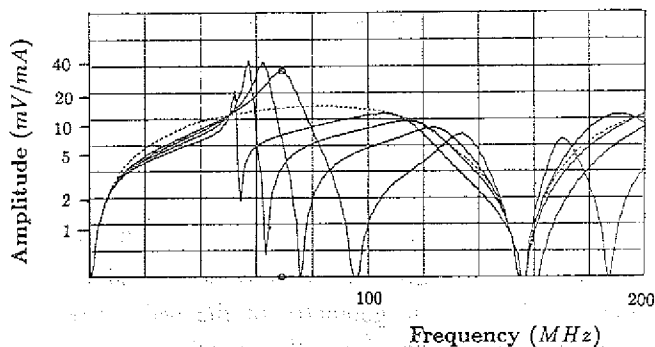


Figure 4: Frequency dependence of the sensitivity (amplitude only) with different resonant tuning. For comparison the dotted line shows the sensitivity without resonant tuning.

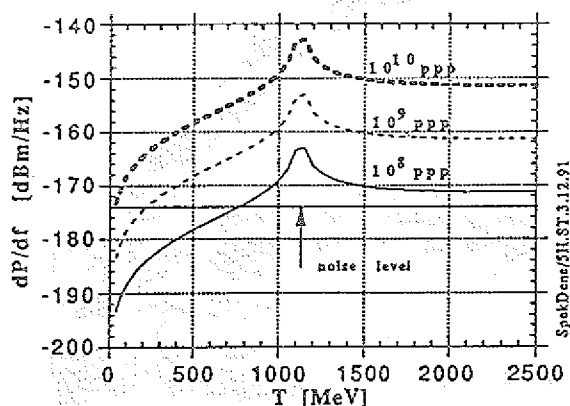


Figure 5: Longitudinal Schottky densities for different particle numbers (ppp=particles per pulse).

A theoretical estimation¹⁾ of longitudinal Schottky densities for an uncooled beam versus kinetic energy T is shown in Fig. 5 at different proton numbers. At injection we assume an initial relative momentum spread $\delta = 4.5 \cdot 10^{-3}$. At higher energies the relative momentum spread is calculated assuming adiabatic damping. The densities are determined for the highest sensitivity of the pickup. Depending on energy this occurs for $f = 75$ MHz at $T = 2.5$ GeV and $f = 35$ MHz at $T = 40$ MeV. The lattice was tuned to $(Q_x, Q_y) = (3.38, 3.38)$ with transition energy at $T = 1032$ MeV.

Averaging over successively measured spectra amplitudes with 0.4 dB above noise level will be visible²⁾. This corresponds to a signal-to-noise power ratio $S/N = -10$ dB at room temperature ($N = -174$ dBm/Hz). The figure clearly shows that spectra of a beam with $N \geq 10^9$ ppp can be easily observed. For $N = 10^8$ ppp a resonating network is necessary to enhance the signal strengths at energies below 200 MeV if an increase in measuring time is to be avoided. Depending on transverse emittance, the corresponding transverse power densities will be of the order of 20 dB less than in the longitudinal case. In this case, even at moderate statistics a resonating network will be necessary if the beam energy lies below 1 GeV. Higher energies ($T \geq 1.5$ GeV) are seen to be less critical.

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- 1) H. Stockhorst, Internal Report
- 2) K. Beckert, U. Schaaf, (GSI Darmstadt), private communication

Schottky Beam Position Monitor

R.D. Neef, Ch. Reul

When a beam bunch of a certain time distribution and the velocity βc travels along the axis of the beam pipe, it is accompanied by image currents on the walls of the beam pipe. For a centered beam, the azimuthal distribution is uniform. Furthermore, the longitudinal distribution of relativistic beams is the same as and coincident with the beam bunch itself. If each stripline subtends an azimuthal angle ϕ , then a fraction $\phi/2\pi$ of the total image current will flow on its inner surface. As the stripline is not grounded, the net induced charge on it is zero. As the image currents pass the upstream gap between the stripline and the ground plane, they flow across the gap, representing the parallel impedance of the stripline and the transmission line at the upstream part. The signal in the transmission line travels at a velocity determined by its permeability and permittivity. This signal has the same time distribution as the beam bunch and half the charge of the image current on the inside surface of the electrode. The beam bunch when it passes the downstream gap, induces similar signals, which again are equal amplitude signals, but of opposite polarity to the signals induced at the upstream port. They propagate out the downstream port, and upstream along the outer surface of the stripline.

The net signal seen at each port is a bipolar doublet with each lobe having essentially the same time distribution as the beam bunch itself. In particular there is complete cancellation at the downstream port if the beam velocity and the signal velocity are the same. For this reason, the upstream port is the signal port.

The longitudinal impedance in an accelerator is the longitudinal voltage per turn seen by the particle, divided by the beam current. As the longitudinal electromagnetic field is not necessarily in phase with the beam current, the longitudinal impedance can be complex. The absorbed power absorbed by the striplines represents the real part of the longitudinal impedance. Additional power will be absorbed by the striplines as the beam is moved off-axis.

The transverse impedance is the lateral electromagnetic field seen by the particle, integrated around the ring, and divided by the beam current,

and divided by the lateral displacement. The important part of the transverse impedance is that part, which is in phase with the displacement, as it can cause a shift in the betatron sideband frequencies.

The signal power scales linearly with the characteristic impedance (or shunt impedance) of the pickup, while the noise power does not. Hence, a high impedance position pickup will provide better resolution.

The three-dimensional computer code system MAFIA, which produces a set of finite-difference matrix equations for the electric and magnetic field vectors in the structure under study yields as solution of these equations the frequency-domain solution or the time-domain solution of Maxwell's equations. MAFIA delivers the values of stored energy, the dielectric loss, and the wall current losses and permits so to evaluate the quality factor or the shunt impedance value.

MAFIA has been installed partly in a preliminary version on IBM ES/9000. This version is able to describe half of the geometry of the wall current monitor only. Therefore, the results are very preliminary, but it was possible to test codes and make parameter studies. The implementation of the complete MAFIA code on CRAY-XMP is almost terminated.

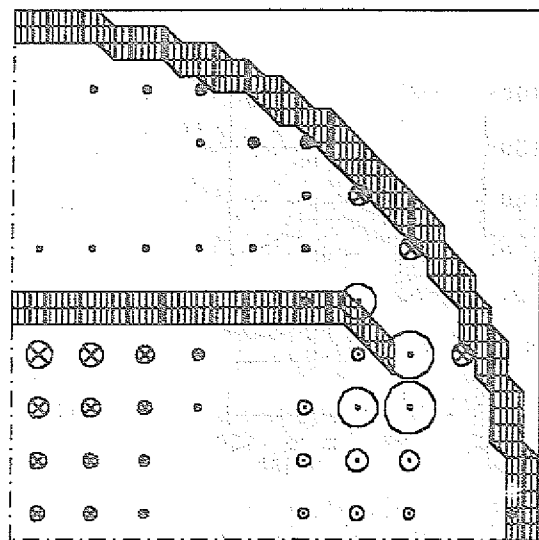


Fig. 1: Real part of time harmonic electric field in a quarter geometry of Schottky monitor

Phase-Switching of the COSY ($h=1$) Cavity for the γ -Transition

Ch. Hamm, H. Meuth, A. Schnase

The cooler synchrotron COSY, accelerating protons from ca. 40 MeV to a maximum energy of 2.5 GeV, features a symmetric (push-pull), ferrite loaded re-entrant accelerating cavity of the SATURNE type¹⁾, with a corresponding tuning range from 460 kHz to 1.6 MHz. Due to lattice-optics constraints, COSY may have to pass through the gamma transition. For this end, the ($h=1$) acceleration system must be capable of undergoing a rapid change of the acceleration phase of about 180° (see Fig. 1) at this transition point. While the phase-jump procedure at the gamma-transition is well known and commonly used, the situation at COSY is considerably more difficult to implement. Since the harmonic number h stands for the integer ratio of the rf frequency over the (proton) revolution beam period, the use of the rf fundamental, $h=1$, implies one turn per rf period (or one beam bunch circulating in the ring). With a cavity quality factor Q of some tens (a measure for the speed with which transients can be forced in and out of the cavity), the operation at the fundamental $h=1$ requires thus many (i.e., on the order of Q/h) particle revolutions for a phase change to take place. This may well lead to errors in the acceleration phase and bunch de-focussing. The method of simply switching the phase at the main amplifier input is thus ruled out for COSY, while it generally is employed with other accelerators going through the gamma transition, for those, as a rule, have $Q/h \leq 1$ by use of a high harmonic.

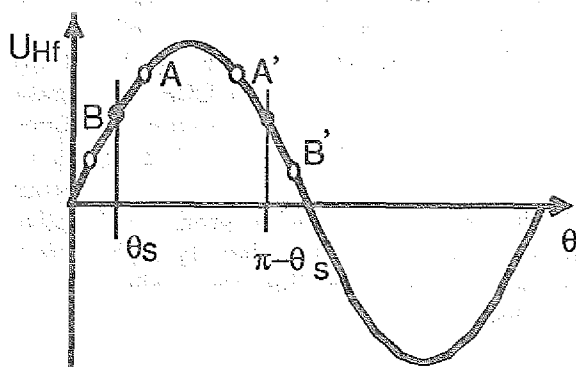


Figure 1:
Acceleration voltage waveform and stable phase values for particle energies below (Θ_s), and above ($\pi - \Theta_s$) γ -transition

Before exploring alternatives in detail, it was felt necessary to test the cavity response under opera-

tion conditions. Amongst those, correct power levels and frequencies are especially crucial in light of the nonlinear ferrite properties. It was our first goal to establish a workable circuit model incorporating model parameters accurately describing the circuit properties as they are accessible to measurement. With these, simulations can be performed, that allow examining various scenarios, before tests on the cavity itself need to be made. This is especially advisable, if the method of overvolting mentioned below is used.

For phase switching, two basic approaches are being contemplated:

(a) direct switching at the amplifier input as discussed, however with a short-time substantially increased power flow into and out of cavity and amplifier, in order to compensate for the long characteristic time constants. But this may require to raise operation parameters above maximum rating.

A second approach (b) would attempt to shift the acceleration phase less rapidly, but very precisely, such that the time-integrated acceleration errors would be minimized, taking into consideration the detailed longitudinal particle dynamics. Such an accurately controlled phase shifting should pose no problem with the novel, fully-digital frequency signal synthesis and phase control, that is being implemented on COSY.

The approach (b) involves considerable investigations of the applicable particle dynamics, as was done first by Johnsen²⁾. Examining approach (a), the knowledge of the cavity's and the amplifier's electrical parameters is of prime importance, as, e.g. quality factor, maximum gap voltage and maximum rating for main power amplifier and cavity, under pulsed and cw conditions. These parameters enter the analytical calculations, and the numerical simulations, using a circuit solver like PSPICE.

The complete cavity-amplifier combination³⁾ would be very complex for detailed circuit modelling studies. For a tractable, simple model, the two power tubes (Thomson-CSF TH120) are replaced by voltage sources with subsequent in-line resistors, each with a value characteristic for the tube's inner resistance. The symmetry of the push-pull scheme of both the two amplifier channels, and the two $\lambda/4$ -cavity segments allows the further reduction

into an equivalent half-scheme (Fig. 2). The tuning provision of the inductance via ferrite biasing is also indicated.

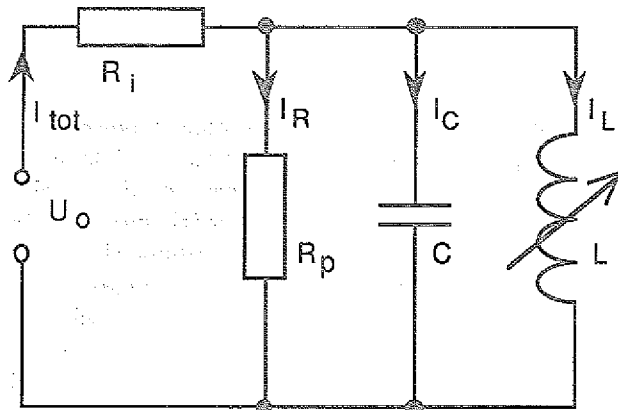


Figure 2:
Simplified circuit diagram

The current through the capacitor follows the differential equation:

$$\frac{1}{R_i} \frac{d^2 U_0(t)}{dt^2} = \frac{d^2 I(t)}{dt^2} + \frac{\omega_0}{Q} \frac{dI(t)}{dt} + \omega_0^2 I(t), \quad (1)$$

$$\text{with } Q = \frac{R_p R_i}{Z_0 (R_i + R_p)} = \frac{Q_0}{1 + \frac{R_p}{R_i}}$$

The characteristic frequency is given by $\omega_0 = (LC)^{-1/2}$, and the effective quality factor Q for the parallel circuit of Fig. 2 follows with the characteristic impedance $Z_0 = (L/C)^{1/2}$. Choosing the voltage $U = C^{-1} \int I dt$ as variable of interest, the general homogeneous solution of Equ. (1) is thus of the form

$$U(t) = A \exp(-\alpha t) \sin(\Omega t + \psi) \quad (2)$$

where $\alpha = \omega_0/(2Q)$, and $\Omega = \omega_0[1+(2Q)^{-2}]^{1/2}$. The constants A and ψ still remain to be determined by boundary conditions in time. The homogeneous solution accounts for all transient characteristics of the circuit. For a harmonic driving voltage of the form $U_0(t) = V_0 \sin(\nu \omega_0 t + \phi_0)$, the stationary solution is of the form $B(\nu) \sin(\nu \omega_0 t + \phi_0 + \varphi)$, with

$$\tan \varphi = Q \left[\frac{1}{\nu} - \nu \right]; \quad (3)$$

$$B(\nu) = \frac{Q Z_0 V_0}{R_i \sqrt{1 + Q^2 \left[\frac{1}{\nu} - \nu \right]^2}}$$

Maximum response, i.e. resonance, is encountered at $\omega = \omega_0$, or $\nu = 1$. Tuning into resonance is achieved by suitably biasing the ferrites. At resonance, the angle between drive and cavity response vanishes, $\varphi=0$, indicating a purely real load.

Leaving aside for the time being all other parameters and nonessential details, we will be most interested in the quality factor, which basically limits the phase switching speed. Now, to determine Q , there are a number of possibilities:

(i) Measuring in steady state the frequency width of the resonance line of Equ. (3)¹⁾. For the 3dB-frequency width of $B(\nu)$, or similarly for the 45°-width of φ , we get $\omega_0/[2\pi Q] = f_0/Q$ (full width). This is done by driving the circuit at fixed tuning condition with frequencies in the vicinity of the resonance, $f = f_0$.

(ii) Recording the transient decay, after switching the input signal to zero at the amplifier. From the observed exponential decay $\propto \exp(-\alpha t) \sin(\Omega t + \psi)$, both the resonance frequency, and the quality factor can be extracted from the decay constant α , and the damped frequency Ω , bearing in mind that $\Omega \approx \omega_0 + 1/(8Q^2) \approx \omega_0$, valid already for moderate Q .

(iii) Last, with the extremely agile rf-signal synthesizer to be used for COSY⁴⁾, the phase of the driving input can be changed at will virtually instantaneously. This method comes the closest to the actual scenario of phase switching, and is therefore used here. We thus switch the phase by 180° at a zero-crossing point of $U_0(t)$, i.e., for reference, at time $t=0$, with $\phi_0=0^\circ$ for $t<0$, and $\phi_0=180^\circ$ for $t>0$. The driving frequency be $f \approx f_0$. After some calculation, this yields the solution, comprising now both transient [Equ. (2)] and stationary [Equ. (3)] contributions:

$$U(t) = V_0 (Z_0/R_i) Q \sin(\omega_0 t) [1 - 2 \exp(-\alpha t)] \quad (4)$$

The value for Q can directly be determined from the time t_c of zero-crossing of the envelope of $U(t)$, i.e. $Q = \pi f_0 t_c / \ln 2$. Figure 3 shows an actual measurement of this last procedure, following the relationship of Equ. (4), for both $U_0(t)$ and $U(t)$. The waveform was taken on the COSY acceleration tests there. Here we have $t_c \approx 8.5 \mu s$, with $f = 1 \text{ MHz}$, and conclude from it (in good agreement with the results from Saclay groups) a relatively high value $Q \approx 38$. This relatively high Q factor poses a challenge regarding the realisation of a rapid phase switch.

Further circuit analysis work will study the remaining parameters in detail, and perform simulations along the lines discussed above.

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- 2) K. Johnsen, Transition, in Proceedings CAS, Gif sur Yvette, 1984
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- 4) P. Shirkey, H. Meuth, H. Rongen, A. Schnase, A fully digital RF-synthesis and phase control for acceleration in COSY, RF-EXPO East, 1990

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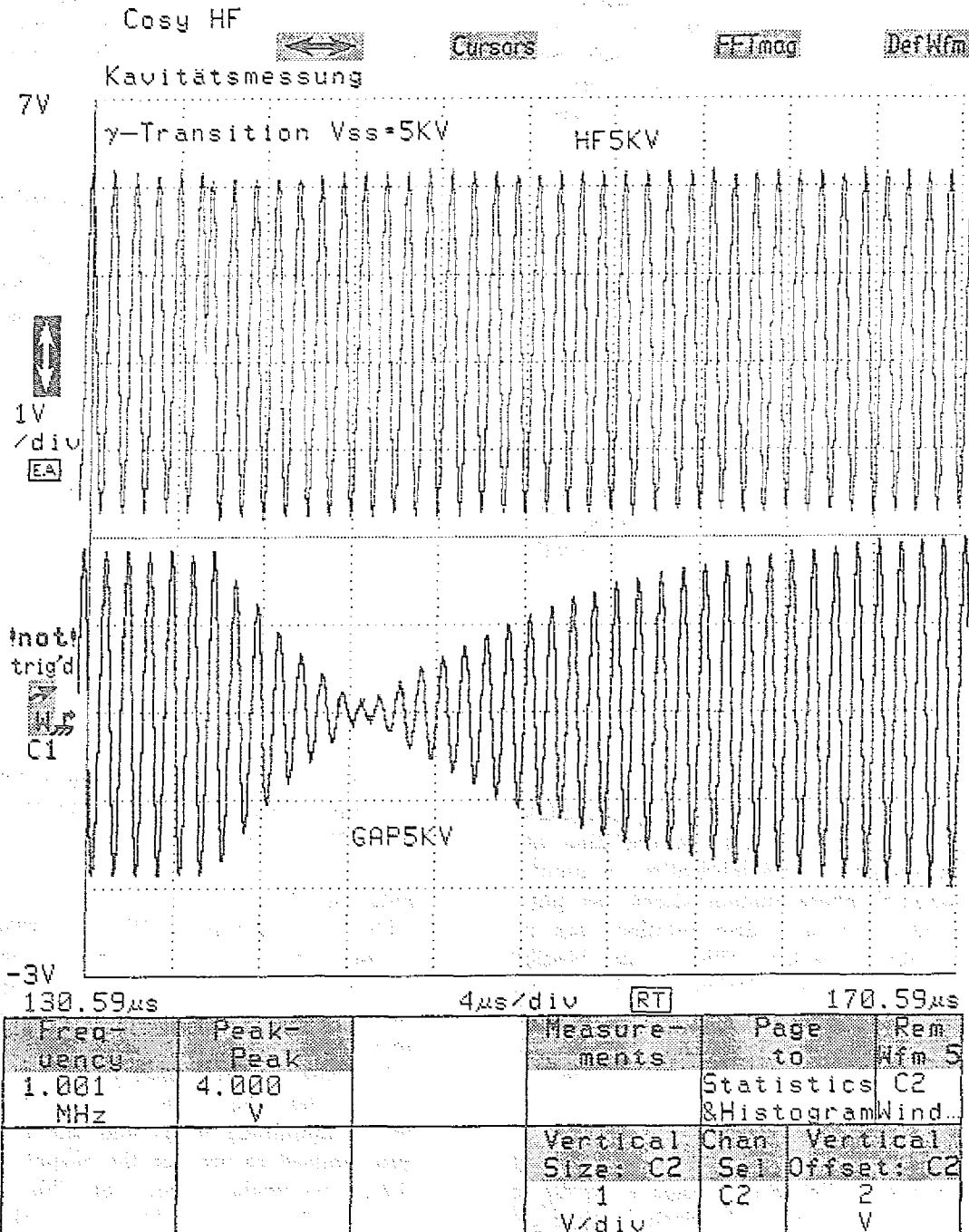


Figure 3: Waveforms for amplifier input (top) and resulting gap voltage (bottom)

A Digital Amplitude Demodulation Circuit for Signals in Quadrature

J. Hagedorn*, H. Meuth, A. Schnase

Introduction

A fully digital RF synthesis and control of the low-level acceleration voltage signal for the ($h=1$) acceleration system in the COSY ring is presently being implemented. Its synthesizer part was recently used for the system's acceptance tests under actual COSY-ramp conditions at the Laboratoire National Saturne at Saclay. Frequencies for $h=1$ at COSY range from 450 kHz to 1.6 MHz. This (relatively low) frequency range permits to attack the task of frequency generation and synchronous phase control, common to any synchrotron for protons or heavy-ions at non-relativistic energies, by an entirely different technology. A first conceptual introduction to our new approach was presented at rf-expo east in November 1990¹⁾.

The performance of this digital signal processing system at unprecedented speed and precision is achieved by a hardware-programmed hybrid scheme, based on an integer algorithm. Digital real time processing at 16 bit value resolution necessitates maximum clock and sampling rates of about 25 MSPS, and heavily relies on parallel pipelining. The algorithm exploits trigonometric calculus, using function values in quadrature. By this, a virtual instantaneous (i.e. within a few clock cycles) phase evaluation is possible, without the usual integration over many rf periods, as it is for analog phase detectors of conventional design.

The control part of the system, with the additional peripheral hardware for beam phase monitoring, is being designed, developed, and built by A.S. as part of a doctoral dissertation. In spite of the system's complexity, the digital approach allows the subdivision into function boards due to a modular design, that facilitates separate optimization and sequential development and fabrication. A number of computational phase manipulations (as phase shifting, modulation and demodulation) are required for a phase lock loop. While it was possible to unite all these manipulations by one board design, the (phase independent) detection and elimination of amplitude variation, required for phase extraction, has to be solved with a different concept.

In the framework of a graduate study project, J. Hagedorn and A. Schnase explored a number of concepts, resulting in two alternatives, that can ensure maximum amplitude dynamic (value) range:

(a) A commercial demodulation and modulation VLSI circuit back to back. This requires a relatively simple circuit board, but is at the expense of a fall through of some 50 clock cycles.

(b) For very short fall through, use can also be made of the trigonometric relationship $[\text{Asin}(x)]^2 + [\text{Acos}(x)]^2 = A^2$ for any A and x , converting the problem into a multiplication, with subsequent root extraction and inversion. The latter two functions can be cast into one and implemented on a set of nested look-up tables with fall through of only a few clock cycles, requiring, however, substantially higher circuit complexity.

The detailed problem

In our phase-control-system, a quadrature scheme is applied. Every signal is represented in the sin and in the cosine signal path. Given are two signals: $S1 = A \sin(\phi)$ and $S2 = A \cos(\phi)$. These are 16 bit numbers in Two's complement notation. The amplitude A is unknown and of no direct interest. The phase ϕ is what we like to measure for phase detection and correction. Both are used in the rf-phase-control-loop.

The pair of signals $T1 = B \sin(\phi)$ and $T2 = B \cos(\phi)$ better fits our needs, if the amplitude B is known and can be normalized to "1". Then one can compute the phase by the following set of conditions:

$$\phi = \begin{cases} 180^\circ - \arcsin |T1| & \text{for } |T1| \leq |T2|, T2 < 0 \\ \arcsin |T1| & \text{for } |T1| \leq |T2|, T2 > 0 \\ \arccos |T2| & \text{for } |T2| < |T1| \end{cases} \quad (1)$$

To get $T1$ from $S1$, and $T2$ from $S2$, we have to multiply both signals with (B/A) . We may do it in two different ways.

Solution (a)

VLSI-chips of type TRW TMC 2330 are able to perform a coordinate transformation between rectangular and polar coordinates and may be used up to 25 MHz clock rate. One chip is programmed to convert the incoming signals $S1$ and $S2$ to the intermediate polar signals A and ϕ . This takes 23 clock cycles. The second chip gets the fixed value B as an amplitude input and the phase ϕ and is programmed to produce the output signals $T1$ and $T2$ in rectangular coordinates. This operation needs another 23 clock cycles. Every clock cycle this design gives us a new update of the signals. But it

takes 46 cycles for every new information to fall through the pipeline. This is approx. 3 turns of the beam and correspondingly limits the control bandwidth of the phase control.

Solution (b)

We present here a more sophisticated, but faster solution. The basic problem was to find the value (B/A) to multiply the incoming signals $S1$ and $S2$ with. We know B and we can calculate A from the well-known formula:

$$[A \sin(x)]^2 + [A \cos(x)]^2 = A^2 \text{ for any } A \text{ and } x. \quad (2)$$

In other words:

$$A = \sqrt{S_1^2 + S_2^2}$$

Now we have to create the factor (B/A) using only integer multiplication, addition, and table-lookup.

First we generate

$$SQ = S_1^2 + S_2^2 \quad (3)$$

This is a 31 bit digital word. We search for the highest nibble and then cut out the 16 most significant bits containing the highest nibble. This is equal to a division by (4^n) where n is in the range $(0..6)$. This value points to a memory location in a look-up-table where the necessary factor

$$F = (B/A) = \frac{B}{\sqrt{SQ}} \quad (4)$$

is stored in a normalized binary fixed format. The denormalization factor is a power of two, since this is the square root of the power-of-four division. Then the two input signals are multiplied with the factor F and denormalized with another circuitry which cuts out the 16 most significant bits of this multiplication. This is controlled by the same n as the input decimation.

To ensure timing integrity, some delay registers have to be added to compensate for the fall through cycles of the multipliers.

Acknowledgements

We are much indebted to Prof. Dr. H. Halling and Dipl.-Ing. H. Rongen, ZEL, KFA, for considerable help, to Professor Dr. rer. nat. H.J. Schmitt, RWTH Aachen, and to Professor Dr. Hill, RWTH Aachen, for their continued interest in this work.

*RWTH Aachen

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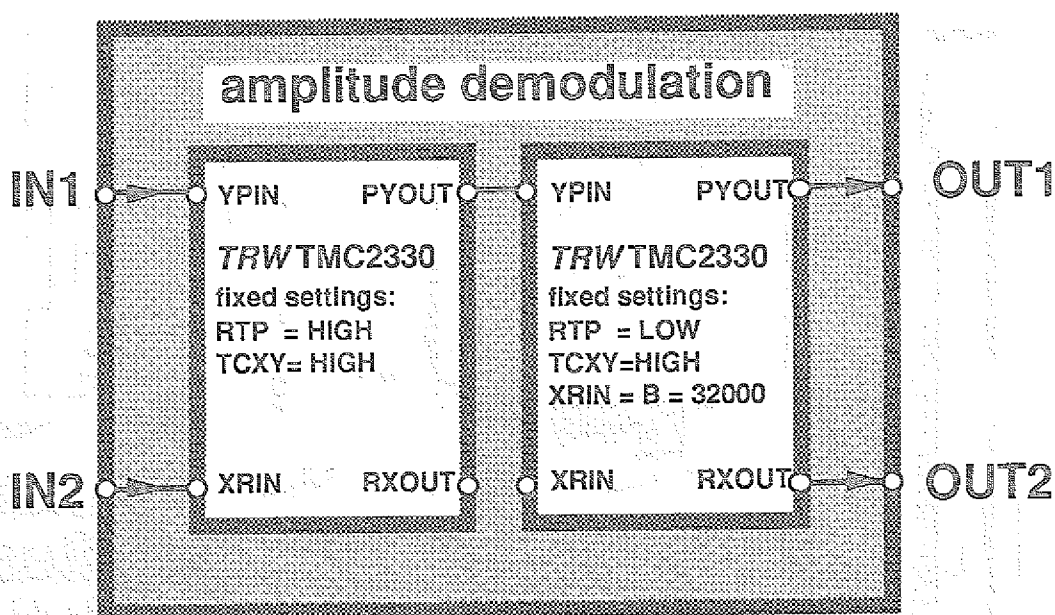


Figure 1: Solution (a) using two coordinate transformer chips

An Analytical Approach to Transverse Tune Measurement by Continuous Wave Excitation

H. Stockhorst, J. Bojowald, K. Bongardt, J. Dietrich, R. Maier

One of the most important machine parameters is the number of betatron oscillation per turn, called tune Q ¹⁾. Its fractional part q can be determined by observing transverse Schottky spectra made up of side bands at each multiple of the revolution frequency, $n \cdot f_0$. By measuring a betatron frequency

$$f_\beta = (n \pm q) \cdot f_0 \quad (1)$$

together with the revolution frequency f_0 the fractional tune q is found. The betatron lines appearing naturally as an incoherent sum of individual particle signals can be strongly enhanced if the beam is excited by a kicker. The method discussed below consists in exciting a steady-state coherent betatron motion by applying a sinusoidal voltage to the kicker. A transverse pickup positioned downstream will deliver a continuously increasing position signal if the driving voltage is in resonance with a betatron line. In general, the tunes are measured for the horizontal (x) and the vertical (y) phase space, yielding values (Q_x, Q_y).

As a first attempt we assume a bunch with duration τ much less than the revolution time $T_0 = 1/f_0$. Also, we neglect effects due to the finite momentum spread of the protons in the bunch. Further, the discussion of one, the horizontal or the vertical phase space plane alone is sufficient if there is no coupling between them. In this way, the beam bunch can be treated as a whole, behaving as it were a single circulating proton. Thus, each time the bunch passages the kicker it receives a kick according to

$$\Delta x'_K(k) = \Delta x' \cdot \cos(2\pi q_D k) \quad (2)$$

where $k = 0, 1, 2, 3, \dots$ is the turn number. Note, that the position is not altered during the kick. The amplitude $\Delta x'$ is derived from the kicker transfer function and the applied kicker voltage. The driving frequency normalized to the revolution frequency f_0 is denoted by q_D . On successive N turns a bunch initially centered at $(x, x') = (0, 0)$ in transverse phase space will then start turn $N + 1$ at the kicker with co-ordinates

$$\begin{bmatrix} x_K^{(N+1)} \\ x'_K^{(N+1)} \end{bmatrix} = \sum_{k=0}^N M^k \begin{bmatrix} 0 \\ \Delta x'_K^{(N-k+1)} \end{bmatrix} \quad (3)$$

The transfer matrix M describes the motion of the co-ordinates in phase space for one complete revolution in COSY. The initial condition for the first turn results from eq. (3) with $N = 0$.

Using known relations²⁾ for the transfer matrix M the sum in eq. (3) can be expressed in closed form. Resonance occurs if the driving frequency equals the betatron frequency (eq. (1)).

In suppressing small oscillating terms, eq. (3) reduces in the resonance case to easy to handle expressions for the position

$$x_K^{(N+1)} = \frac{\beta_K \Delta x'}{2} N \sin(2\pi q N) + \text{small osc. terms} \quad (4)$$

and for the angle co-ordinate

$$x'_K^{(N+1)} = \frac{\Delta x'}{2} \cdot N \cdot \{ \cos(2\pi q N) - \alpha_K \sin(2\pi q N) \} + \frac{\Delta x'}{2} \cdot \left\{ \cos(2\pi q N) - \frac{\cos(2\pi q N) - \cos(2\pi q (N+2))}{1 - \cos(4\pi q)} \right\} + \text{small osc. terms} \quad (5)$$

if $q \neq 1/2$. For $q = 1/2$ two side bands contribute and the amplitudes must be multiplied by two. The parameter α_K denotes the orientation of the phase space ellipse at the kicker.

In resonance the amplitudes of the bunch position as well as of the divergence at the kicker will increase with the number of turns. Also, the position amplitude is proportional to the betafunction at the kicker. In both cases the amplitudes will grow with increasing driving voltage.

An example of formula (4) is illustrated in Figure 1 for 200 turns. The fractional tune $q = 0.38$ and the normalized driving frequency $q_D = 1.38$ were chosen. The excitation is calculated with the transfer function of the stripline unit³⁾ as kicker. The stripline kicker consists of four plates rotated by 45 degrees around the beam direction. It is therefore possible to produce horizontal and vertical deflections simultaneously. An input power $P = 30$ W per plate will result in a kick amplitude $\Delta x' \approx 2 \cdot 10^{-7}$ rad at kinetic energy $T = 2.5$ GeV. The lattice parameters for the horizontal phase space are $\alpha_K = -0.7$ and $\beta_K = 26$ m at the position of the stripline unit ($s = 98.4$ m). Eq. (4) easily allows an estimation of the position amplitude at any N .

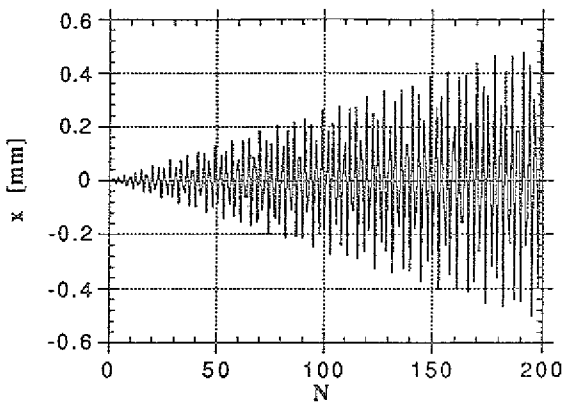


Figure 1:
Position signal at the kicker versus turn number. The kicker was driven in resonance with the betatron motion at $q = 0.38$.

In Figure 2 the phase space evolution is shown for 200 turns after switching the kicker on. The bunch centered in $(0,0)$ at the beginning of the kicks spirals out to larger amplitudes.

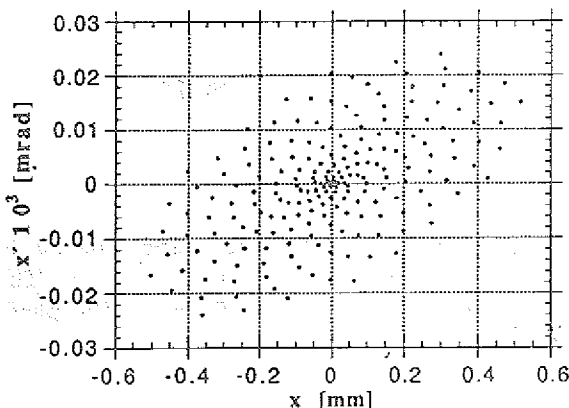


Figure 2:
Phase space evolution in the resonance case

The amplitude of a fast Fourier transform (FFT) of the signal (4) clearly shows up a peak at $q = 0.38$ (Figure 3). Note that the width of the line is due to the finite signal duration.

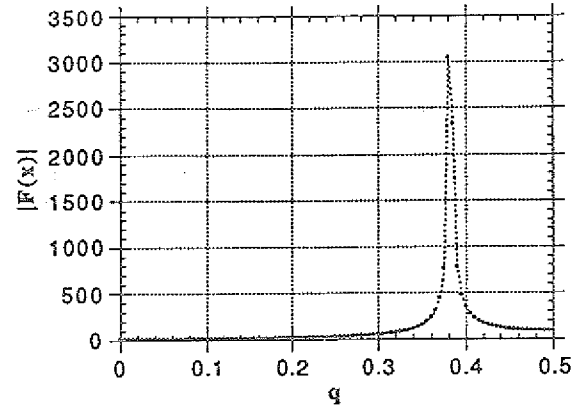


Figure 3:
Amplitude of a FFT of the position signal at the kicker versus fractional tune. The kicker was driven in resonance with the betatron motion at $q = 0.38$

The off-resonance case is characterized by an oscillating solution as shown in Figure 4 for a normalized driving frequency $q_D = 1.39$. As compared to Figure 1, the maximum amplitude has now dropped down significantly.

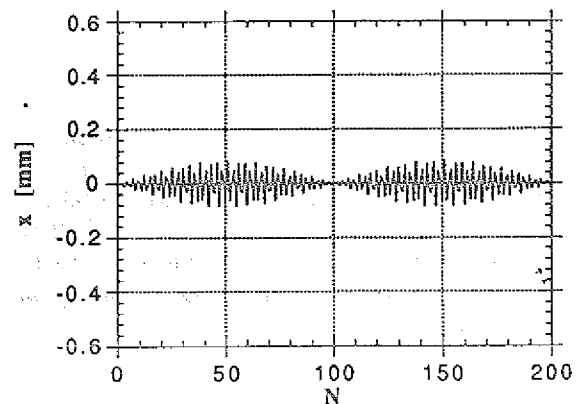


Figure 4:
Position signal at the kicker versus turn number in the off-resonance case for $q_D = 1.39$

Up to $N \approx 1/(2 \delta q)$ turns with $\delta q = |q - (q_D - 1)|$ the amplitude growth is rather similar but then it is seen to diminish, crossing zero at $N \approx 1/\delta q$. The corresponding FFT in Figure 5 shows predominantly the fractional part of the driving frequency.

Also, the betatron motion at $q = 0.38$ is visible, albeit slightly reduced in amplitude. Both lines can be well separated if the measuring time is lengthened (≈ 1000 turns). Observe that the amplitudes are about a factor of 6 less than in the resonant case. Similar results are found for the x' coordinate.

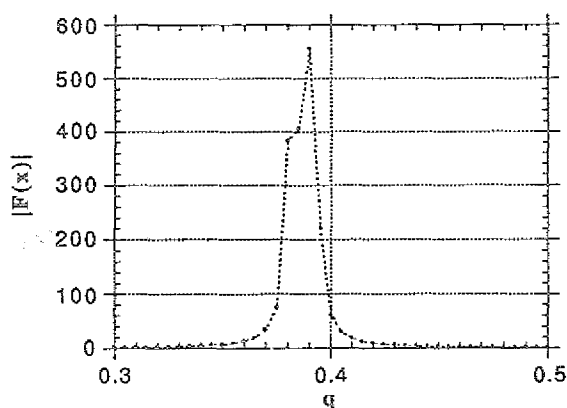


Figure 5:

Amplitude of a FFT of the position signal in the off-resonance case. Predominantly visible is the fractional part of the drive, $q_D = 1.39$

So far we have treated a beam bunch of small duration. If the bunch duration τ is increased, bunch shape and width become important. It can be shown⁴⁾ that in general the excited spectrum will contain all lines falling in the range d.c. to $1/\tau$. However, the amplitudes are reduced at larger frequencies⁵⁾. For small τ a large number of lines are excited. This makes it possible to shift the observation frequency to an appropriate line if desired (e.g. to avoid cross talk between excitation (30 W) and pickup signal (pW)).

From eqs. (4) and (5) the resulting co-ordinates in phase space at any place in the ring can be found if the transport matrix from the kicker to that point is known. In Table 1 we compile the results for the position pickup DPOS16, located at $s = 154.76$ m in COSY tuned to $(Q_x, Q_y) = (3.38, 3.38)$. A square shaped bunch with duration of $\tau = 0.1$ msec⁶⁾ was used at kinetic energy $T = 2.5$ GeV ($T_0 = 0.6$ μ sec). The excitation frequency was 17.3 MHz, corresponding to the betatron line at harmonic number 10. For this frequency the maximum absolute value of the kicker transfer function is achieved. The frequency extend around this betatron line is about ± 4 MHz @ -3 dB level. The vertical position amplitude is only half the horizontal value due to the different lattice optic in the vertical plane.

turn no. N	τ [μ sec]	horizontal x_p [mm]	vertical x_p [mm]
50	30	0.1	0.05
100	60	0.2	0.1
500	300	1	0.5
1000	600	2	1
1500	900	3	1.5

Table 1

In the preceding discussion we have neglected all momentum dependent effects. A quantity which should be examined more carefully is the change in tune with momentum spread, $\Delta Q/Q = \xi \cdot (\Delta p/p)$, in combination with the synchrotron period T_S in longitudinal phase space. A finite

chromaticity ξ of the machine will then lead to a tune spread ΔQ which is modulated in time caused by the synchrotron motion. If the chromaticity is large and the observation time becomes comparable with the synchrotron period (app. 1 ms for COSY) a large fraction of the bunch particles possibly runs out of phase leading to a partial decoherence of the resonant signal.

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The COSY Control System

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Architecture

The COSY control system is hierarchically organized with distributed intelligence and autonomous operation units for each accelerator component. Data communication is performed via LAN and over fieldbus. The hostsystems are Un*x-based, whereas the field controllers are running a modular realtime kernel RT/OS which has been developed at the COSY project. The computer hardware consists of RISC mini-computer from Hewlett-Packard, VME-computers in the field and G64 equipment-control-modules in geographical expansion of the VME-bus. The man-machine interface is built using object based methods on top of the X-window system. A cross development environment including a high-level-debugger hosted on the Un*x machines is used for the VME-target software development. Together with the software development system for the Un*x applications a uniform view of the system appears to the programmer.

The device connectivity

To describe the technology of dataflow the 7 layer ISO/OSI model is transformed for device connection. This model fits in a design for all the protocol handling software and the interaction between different places in the control system. This model helps to find the position in the implementation process where different layers interoperate and so different modules are built and connected. This model helps also to identify the positions where protocols will be used and which are the contents of the protocols (Fig. 1). This is true for protocol features like address resolution, data validating, handshaking and so on. A second major point is the chance to find with this understanding of the protocol states the places in the implementation process where automatic protocol creation tools can be used.

SCSR/UState

For the realtime part of the control system a special module is implemented. This module called single command single reply handles the communication between the control application side and the process picture on one hand and on the other hand the accelerator equipment. The SCSR is built from a definition written in bison and flex input for the protocol decoding and a couple of pDDs (Fig. 2).

Device Connectivity protocol components

Layer

- 7 Application Program
 ➡ GUI
- 6 Device Management, Configuration
 ➡ target management
- 5 Generic Interface
 ➡ SCSR/UState
- 4 Message Interface
 ➡ pDD
- 3 Protocol Interface
 ➡ Sci / p2pc
- 2 Hardware Driver
 ➡ device driver

- 1 Hardware

Figure 1

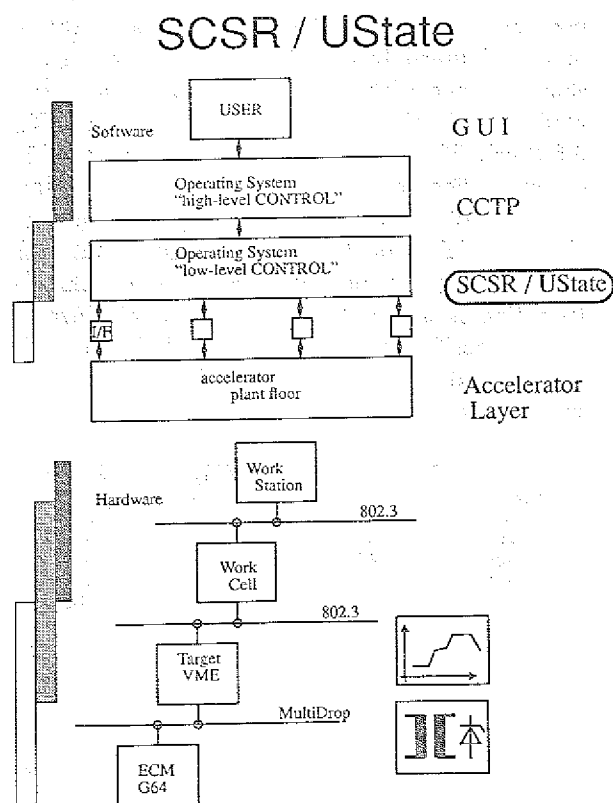


Figure 2

PDDs, pseudo device driver, are abstract classes of a model which knows to handle the equipment I/O. SCSR creates a concrete class in a process step, we call connection. The connection builds an instance of the class and opens all file descriptors to get access to the device driver. The communication between different instances of the same object is done by messages (Fig. 3). The messages are

SCSR - pDD

abstract classes for device access

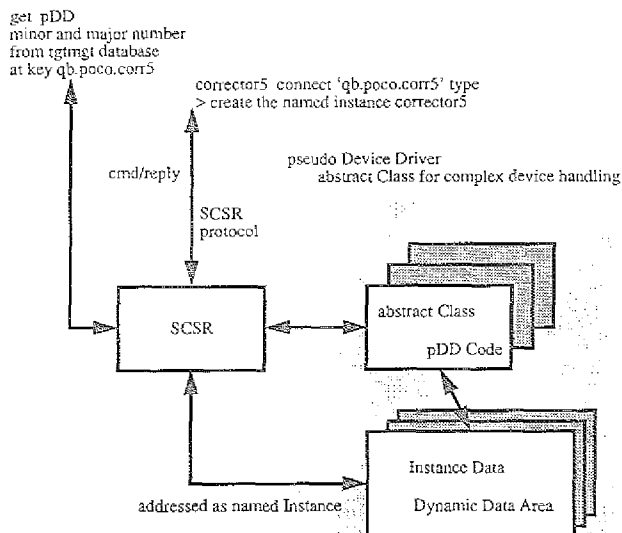


Figure 3

represented as ASCII strings. This frees us from the low and high endian problem. The messages carry only the new information for the object. The COSY control transaction protocol handles the data path finding and updates the logbook and the process picture. PDDs get the messages already decoded. The concrete class of pDDs acts like tasks (Fig. 4).

SCSR - pDD

instance creation

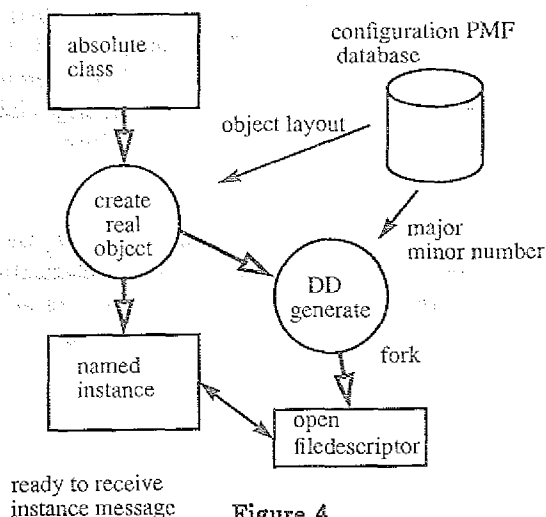


Figure 4

The normal communication via SCSR is synchronous. For pseudo devices there is also an asynchronous communication channel called unrequested state (UState). The messages send here will be handled like normal replies.

Every other component can receive these messages on its UState part. For faults in the control system or on devices there is an asynchronous notification system based on the syslogged protocol coming along with the Un*x operating system.

CCTP analyses the path information and sends the messages to the receiver (Fig. 5). In extension to this it feeds the information to the logbook and to a decoder which updates the state and data in the process picture. The process picture is built on top of the online data base PMF. PMF, process message format, is a single key data base implemented using the buffercache. The buffercache in junction with nfs provides the mechanism to distribute the data. In addition PMF has a state monitor. The state monitor can be triggered if records are created or changed. The events are sent to the destination by the RPC protocol.

CCTP

COSY Control Transaction Protocol

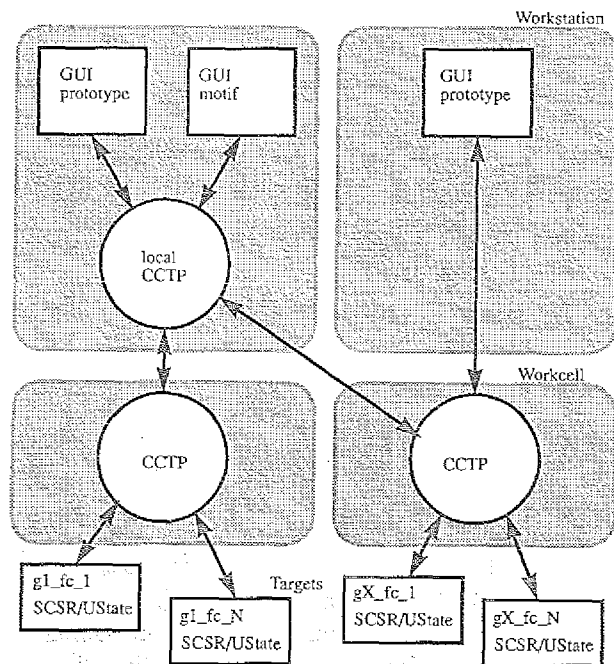


Figure 5

The present status of theoretical investigation concerning the behaviour of a polarized beam in COSY itself is presented here. The basic physics of spin resonances in COSY, notation and resonance energies are given in¹⁾. The discussion of linear effects including different aspects of tune modulation due to synchrotron oscillations and sensitivities of correction methods is rather finished now. Within this context also the realizability of a "partial snake" concept in COSY is considered, which would cure the polarisation losses caused by imperfection resonances in a rather easy to handle way. The physical specifications for the design of the pulsed quads are basically influenced by the requirement of both tune- and γ_T -jump in COSY.

Calculations concerning tune-jump are finished now whereas the study of γ_T -jump schemes is still in progress.

Imperfection resonances

During acceleration in a synchrotron the usually linear rise in spin-tune, which is directly proportional to the energy of the particle, is modulated by synchrotron oscillations. This causes a highly nonlinear crossing of resonances giving rise to a complicated interference pattern in contrast to the Froissart-Stora relation. This effect was first observed at SATURNE²⁾ and qualitatively explained in³⁾. The modulated crossing speed reads now:

$$\alpha(\Theta) = \frac{\Delta\gamma G}{2\pi} + \left[\left(\frac{\gamma_0^2 - 1}{\gamma_0^2} \right) \gamma_0 G - \xi \right] \nu_s \left(\frac{\Delta p}{p} \right)_{\max} \sin(\nu_s \Theta + \delta_s)$$

where G is the gyromagnetic anomaly, ξ the chromaticity and ν_s the synchrotron tune. The nonlinear behaviour is increased with increased synchrotron tune. Only for small energies with a large synchrotron tune ($\nu_s \approx 2 \cdot 10^{-3}$) one expects dramatic effects in COSY. For energies above 200 MeV synchrotron modulation can be neglected. In Fig. 1 the polarisation after crossing the $\gamma G = 2$ imperfection resonance as a function of the excitation strength is shown.

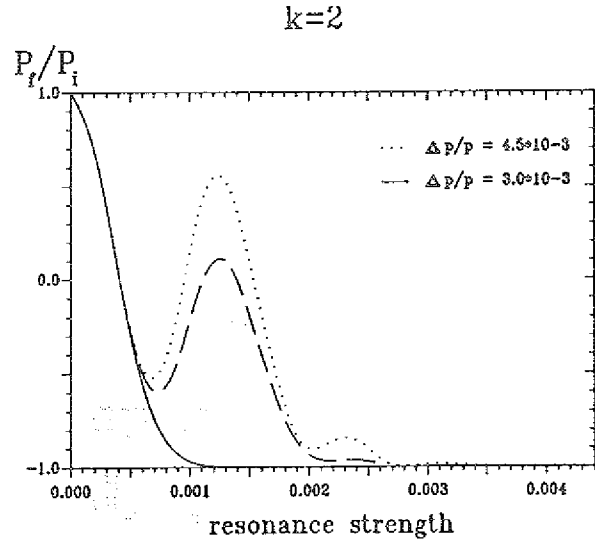


Figure 1:

Depolarisation after crossing the $k=2$ resonance as a function of the excitation strength for different initial momentum spread at injection (40 MeV)

A parabolic line charge density is assumed for different maximum momentum spread at injection. The parameters used will be applied in operation of COSY for maximum possible crossing speed. The maximum modulation amplitude of α is here $\Delta\alpha_{\max} \approx 1,7 \cdot 10^{-6}$, four times larger than the time independent α at maximum acceleration. Adiabaticity is now suppressed even for a strong resonance. To reach a resonance strength of more than $\epsilon_R = 2,5 \cdot 10^{-3}$ the averaged closed-orbit distortion must be larger than 9 mm with maximum distortion up to 24 mm. The vertical aperture of COSY is 27,5 mm. Compensation of the resonance strength therefore is inevitable in order to avoid particle losses. But due to the small crossing speed the requirement on stability for the applied steering elements is rather high ($\Delta B/B \approx 10^{-3}$). Especially remanent fields can produce nonreproducible polarization loss.

Therefore a "partial snake" seems to be more practical for conservation of vertical polarisation especially at the $k=2$ resonance. But there are other reasons in COSY which would favour this concept.

In COSY the cooler solenoid which already exists, is strong enough to flip the spin even at the $k=6$ resonance. Therefore a "partial snake" can be realised only with the main solenoid of the electron cooling devices, switching off all toroids and compensation solenoids. Starting with an integrated longitudinal field strength of 0,063 Tm at 108 MeV ($k=2$), which rotates the spin about the longitudinal axis by $\delta = 6,5^\circ$, and keeping the field constant during acceleration, the final angle at 2201 MeV ($k=6$) would be $\delta = 1,0^\circ$. The latter angle might be too small but an appropriate closed-orbit correction and an increase of the solenoid strength during acceleration by a factor of two at most would cure any polarization losses even at the highest lying imperfection resonance. Here it should be mentioned that the "partial snake" concept is insensitive against field fluctuations because the "partial snake" produces an adiabatic spin-flip. At extraction energy the solenoid can be switched off adiabatically.

The problem of additional phase space coupling due to the "partial snake" device during acceleration and at extraction and the excitation of weak coupling resonances is still under investigation.

Intrinsic resonances

The COSY synchrotron will operate normally with a three parameter lattice, with only two identical cells left. In this situation four strong depolarizing resonances will exist. With non-zero dispersion on the straight sections a lattice with 6 identical cells is possible with only one resonance left¹⁾ ($\gamma G = 8 - \nu_z$). The rather large strength of this resonance makes it impossible to flip the spin or conserve the polarisation¹⁾. In Fig. 2 the polarisation loss is plotted as a function of the rms-emittance (Gaussian distribution assumed in transverse phase-space). The solid line shows the polarisation mirrored on the abscissa for comparison. Only for very high emittances the polarization loss will be less than 5%. In general even in the case of the 6-parameter lattice a tune jump must be applied. The dotted and dashed lines show results of simulations assuming a tune-jump per turn of $\Delta\nu_\tau = 0,002$ (resp. $\Delta\nu_\tau = 0,004$) applied on 10 turns. The polarisation loss is reduced to less than 10% (< 5%) for all emittances. The parameters for the unipolar tune jump are summarized in Tab. 1. The same tune jump applied to all the other three resonances will add only 2% (1%) polarisation loss.

Different schemes for positioning pulsed quads have been discussed. In the arcs there are six possible positions for pulsed quads of a length of about 30 cm. With equal strength and an average betatron function of 15 m one needs a strength of $k=0,03 [\frac{1}{m}]$ in order to get the wanted tune jump of Tab. 1.

The situation in COSY is further complicated, if one tries to jump γ_τ too. Here the scarcity of empty space in COSY forces us to use some of the quads for both tune- and γ_τ -jump. Especially the combined jump in the same cycle is a rather bothering constraint. Calculations concerning these problems are in progress.

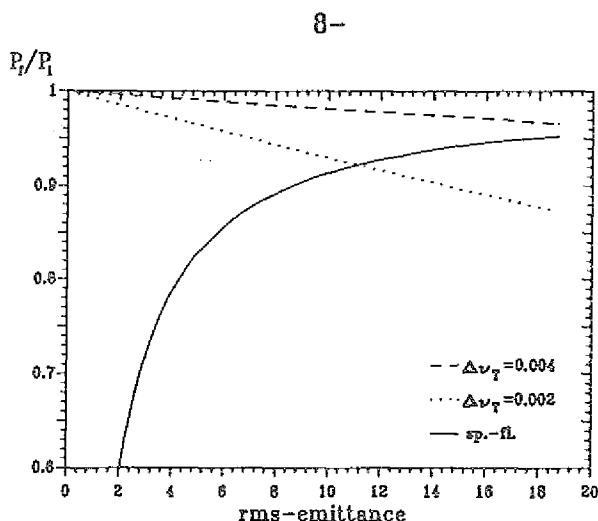


Figure 2:

Polarisation loss as a function of the rms-emittance for crossing the (8-) resonance. A Gaussian phase-space distribution is assumed: spin-flip (solid line), tune jump with $\Delta\nu_\tau = 0,004$ (dashed line) and $\Delta\nu_\tau = 0,002$ (dotted line)

Pol.-loss	$\Delta\nu_\tau$	2δ	Δt_1	Δt_2
<10%	$\approx 2 \cdot 10^{-3}$	$\approx 4 \cdot 10^{-2}$	13μs	11ms
<5%	$\approx 4 \cdot 10^{-3}$	$\approx 6 \cdot 10^{-2}$	10μs	16ms

Table 1:

Parameters for unipolar tune jump at (8-) resonance (Δt_1 = rise time, Δt_2 = fall time)

References

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COSY Stripping Target

J. Pfeiffer, G. Riepe

For the use as large-area stripping targets the fabrication of thin carbon foils by electron beam evaporation proved to be a relatively reproducible process, at least as compared to the procedure for the preparation of Al_2O_3 -films by electrolytic oxidation.

Targets in thicknesses of $20 \mu\text{g}/\text{cm}^2$ were prepared by attaching two carbon foils of half the thickness to both sides of a frame, thus compensating the internal stress. Afterwards one side of the frame is taken off, and a C-shaped target with a free area of 40 by 25 mm^2 is obtained.

Carbon targets of this type survived baking up to 300°C in vacuum.

For the injection system of the CELSIUS-ring in Uppsala a series of carbon targets had been made following these procedures. According to information from Uppsala they proved to last very long, i.e. at least several weeks.

S. Martin, C.A. Wiedner*

The results of the ion optical calculations with the TRANSPORT code for the COSY external beam lines 2 (medical area), 3 (TOF) and 4 (BIG KARL) are described.

Beam parameters at extraction

For the TRANSPORT calculations the COSY tune with $Q_x = Q_y = 3.38$ and the corresponding TWISS parameters of Table 1 have been used.

Table 1: TWISS Parameters at Extraction for the COSY - Tune ($Q_x/Q_y = 3.38/3.38$)

$\beta_x = 7.2$ m	$\beta_y = 12.7$ m
$\alpha_x = 0.14$	$\alpha_y = 0.07$

The properties of the ultra slow extracted beam is basically determined by the mechanism of the resonance extraction, but not by the properties of the circulating beam. Here the geometric phase space properties are taken from the ring optic as an estimation for the first runs. It has been assumed that the angular dispersion vanishes at extraction point ($R_{26} = 0$); the radial dispersion $D_x = R_{16}$ was varied in a rather large range: $0 \leq D_x \leq 5$ m. Extraction simulation by tracking leads us to estimate a dispersion of less than 0.4 m.

From the values of the β -function at extraction the TRANSPORT beam parameters are readily obtained from the assumed - momentum dependent - beam emittance $A = \pi \cdot \epsilon$. They are listed in Table 2.

Table 2: Beam Parameters for the TRANSPORT Runs.

ϵ [mm mrad]	x [mm]	x' [mrad]	y [mm]	y' [mrad]
5	6.0	0.840	7.97	0.629
2.5	4.243	0.595	5.635	0.445
1	2.683	0.376	3.564	0.281

Since the TWISS parameters $\alpha_{x,y}$ are finite at extraction point, we assumed for the TRANSPORT calculations there the corresponding matrix elements $R_{12} = -0.14$ mm/mrad and $R_{34} = -0.07$ mm/mrad. For the beam a momentum of $p = 3.3$ GeV/c was assumed, for the momentum spread $dp/p = \pm 0.25\%$ because this is the maximum momentum acceptance of COSY for the working point presently under discussion.

External beamline 2 (medical area)

The beam acceptance of beamline 2 was assumed to be $\epsilon_x = \epsilon_y = 2.5 \cdot \pi \cdot \text{mm} \cdot \text{mrad}$. The beam dimension at the target position was specified to be:

$$x = y < \pm 2 \text{ mm.}$$

Therefore as optical solution a doubly achromatic waist in both coordinates has been envisaged. Compared to the proposed design in the COSY

Report No. 84 we have therefore added two further quadrupole singlets - 3MQ21 and 3MQ22.

The resulting beam dimension at the target position is contained in Table 3, again for different values of the beam dispersion at the extraction point. The remaining dispersion at the target point 2 is less than 0.17 m.

Table 3: Beam Radii at the Target of Beamline 2. The divergence is $x' = 1.34$, $y' = 1.24$ mrad

D [m]	Mom. spread dp/p [%]	x [mm]	y [mm]
0	0.5	1.86	2.01
1.25	0.5	2.97	2.46
	0.1	2.27	2.46
2.5	0.5	3.41	0.88
	0.1	3.25	1.88
5.0	0.5	2.21	1.46

Beamline 3 (Time - Of - Flight)

Also for this beamline we have assumed an acceptance of $\epsilon_x = \epsilon_y = 2.5 \cdot \pi \cdot \text{mm} \cdot \text{mrad}$. Compared to the previous concept two additional vertical bending magnets - 3MBE34 and 3MBE35 - had to be included to enable lifting the beam by 80 cm at the TOF - target position. An ion optical consequence is a finite vertical dispersion. The beam dimensions, radially and axially at the target position, were specified by the experiment to be

$$x, y < \pm 1 \text{ mm.}$$

The beam dimension at the target position is given in Table 4, for the different values of the beam dispersion at extraction. The remaining dispersion gives a contribution to the beam size at the target if the dispersion at the extraction point is as large as 5 m in the worst case. The angular dispersion is negligible, so that the divergence at the target is almost constant.

Table 4: Beam Radii at the TOF Target for different Dispersions at Extraction. The divergence is $x' = 4.90$ mrad, $y' = 3.47$ mrad

D [m]	Mom. spread dp/p [%]	x [mm]	y [mm]
0	0.5	0.51	0.72
1.25	0.5	0.75	0.81
2.5	0.5	1.38	0.58
	0.1	0.56	0.52
5.0	0.5	2.18	0.58
	0.1	0.50	0.52

Beamline 4 - BIG KARL spectrograph

For this beam line again a beam acceptance of $\epsilon_x = \epsilon_y = 2.5 \cdot \pi \cdot \text{mm} \cdot \text{mrad}$ has been assumed. The beam dimensions and divergences at the target are contained in Table 5, for different values of the beam dispersion and for a beam momentum of $p = 3.3$ GeV/c.

Table 5: Beamradii and - divergences at the BIG KARL target for different dispersions at extraction.

D [m]	Mom. spread dp/p [%]	x' [mrad]	y [mm]	y' [mrad]
0	0.5	0.50	5.76	0.50
1.25	0.5	0.52	5.05	5.94
2.5	0.5	0.53	5.44	7.04
5.0	0.5	1.43	7.15	6.92
	0.1	0.45	7.14	6.92

*MPI Heidelberg

The general configuration as well as many details have already been described (ref. 1,2,3,4,5,6). The delay in the completion of the rf system caused by unexpected technical problems forced us to rearrange the sequence of activities at the cyclotron. In especial the installation of the new trim coils inclusive the correspondent fieldmapping, the realization of new cooling systems and the construction of the transfer beamline between cyclotron as COSY-injector and the COSY-ring interrupted the work at the rf. But in September 1991 the new rf-system was principally ready for operation.

Figure 1 is a simplified scheme of the total system. Operation of the rf is controlled by a special dedicated computer system based on a Motorola series 68000 processor. It not only controls more than 20 tuning elements but also monitors the proper operation of signal detection and conditioning, regulation loops etc. Data of the rf-system, either calculated or experimentally evaluated for a number of frequency settings, are stored in EEPROM's (the computer has no mass storage device). The rf-system is tuned according to this data; intermediate settings are calculated by use of splining methods.

The frequency is generated in a programmable master oscillator, which was selected for a very good short term stability. The main signal path to the accelerating system appears from the figure. The power amplification is achieved in a two stage setup. Two solid state wideband amplifiers of 1 kW each in parallel serve as preamplifier for a power stage equipped with the SIEMENS tetrode RS 2058 CJ. The nominal output power of this setup is

100 kW in the total operating frequency range of 20 to 30 MHz. Forward and reverse power is monitored by directional couplers at the output of the transmitter. These signals are used for fail safe (trip circuits) and for regulations purposes.

The rf power is transferred to the cyclotron via a 6 1/8" coax-line over a distance of about 21 m. After passing a tuned impedance transforming network, the power is coupled inductively into one accelerating sector (one of three dees of the cyclotron). Several closed control loops are installed to control

- resonance frequency of the accelerating system;
- phase stability;
- amplitude stability.

Multiple closed loop control was necessary for amplitude and phase regulation, to achieve the specified long and short term stability. The resonance frequency of the accelerating system (RESONATOR in Fig. 1) is preset within the operating range by motor driven panels in the dees. Simultaneously also the adjustable coax line on top of the machine is pretuned (ref.1,5). This combined coarse frequency setting is done in such a way, that the peak rf voltage (and hence minimum rf current) appears in the center of the cyclotron. The coax line is terminated with two watercooled variable vacuum capacitors in parallel, which are the active elements of the resonance frequency fine tuning loop arranged around the resonator. This loop controls the phase offset between the signal in the feeding line (leading phase) sensed by directional coupler 3 and the rf voltage in the resonator, which is decoupled by a small loop probe. The phase of

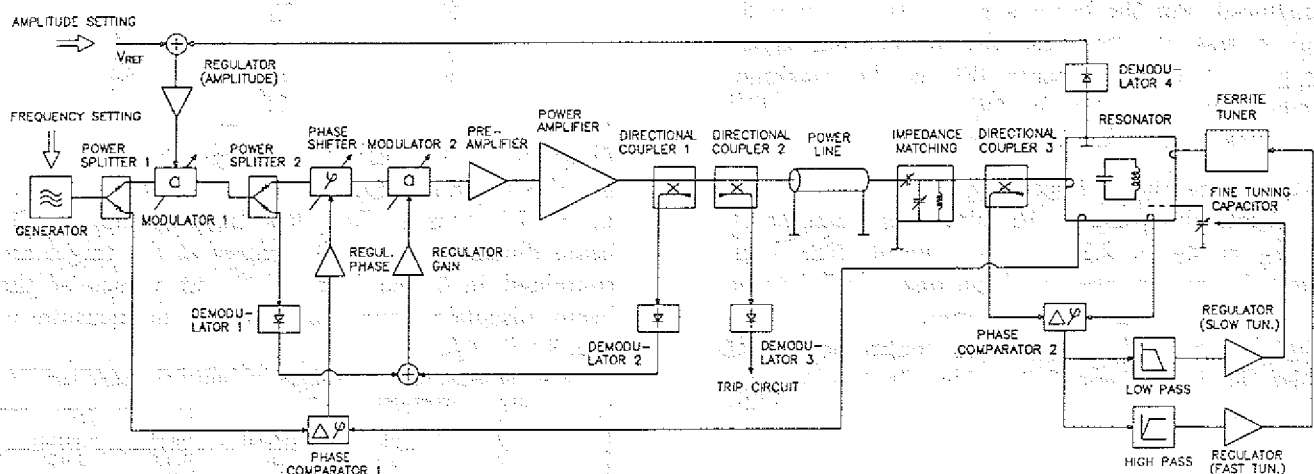


Figure 1: Simplified circuit diagram of the rf-system

these two signals is detected in phase comparator 2, then splitted into a low pass and a high pass channel and used for the fine tuning of resonance frequency. So far only the slow branch is used. The fast operating part with a ferrite tuner as active element is installed to compensate rapid changes in resonance frequency, which are caused by mechanical vibrations of the large dee system. It turned out during first tests, that the specified figures of stability could be achieved without this option. It is planned to take this subsystem into operation at a later date with the idea, to further improve the performance of the system. The operation of the analog resonance frequency fine tuning loop is continuously surveyed by the computer. If for example the tuning C during variation of frequency, caused by the operator or by warming up effects in the resonator, approaches the limits of a predefined range, then the computer will generate an incremental step of the coarse tuning elements for correction.

Amplitude regulation is performed by means of two superimposed closed control loops. Because of the stability criterion an overall feedback loop, which includes the high Q resonator, can hardly be made fast enough to reduce effectively modulation contents caused by harmonics of the mains frequency. For this purpose an inner loop was installed around the amplification stages without this stability constraints. This loop stabilizes the amplification factor between power splitter 2 and directional coupler 1, hence suppressing amplitude modulation contents introduced by the non ideal power amplifier. By this method the initial amplitude stability of the master oscillator is transferred to the output of the power stage. Amplitude setting is achieved by use of an overall regulation loop. The characteristics of this loop define long term stability. An amplitude signal is decoupled from the resonator by a capacitive probe and rectified by demodulator 4. As the specs of this unit are essential for the stability factor, here a modul is used, which was especially developed some years ago by ZEL (Zentrallabor für Elektronik des Forschungszentrums Jülich). By use of compensation techniques the temperature coefficient is trimmed to $t_c \ll 1 \cdot 10^{-4}/^{\circ}\text{C}$. Up to three probe/demodulator paths can be used in parallel to reduce instabilities and noise.

The regulation scheme is completed by an overall phase regulation loop. As appears from Fig. 1, the phase of the acceleration voltage in the resonator is directly locked to the generator.

Figure 1 shows the situation in normal mode operation. Because of the influence of the multipactor effect⁷⁾, for the switching on procedure the system configuration is modified automatically by the control computer. In this case trip- and safety-circuits as well as parts of the regulating loops are

switched to an inactive state. Amplification factors for signal conditioning etc. are adapted, to respond on the very low levels during start up. After pre-setting all tuning elements in the system, rf power of maximum only 1 Watt is applied to the resonator, to be safely free from the influence of multipactoring. When the resonator has been precisely tuned in closed loop mode to the operating frequency, then a rf burst is generated for start up. For approximately 10 μs the maximum available rf power is applied to the system, in order to increase the voltage with a slope of $> 1\text{ kV}/\mu\text{s}$ in the resonator, which is necessary to overcome the multipactor effect. In case of failure this process is periodically repeated up to 16 times, before the computer generates an error message. If the start up is successful, then after a transition time of about 500 μs the system is switched to normal mode operation. Figure 2 illustrates this procedure.

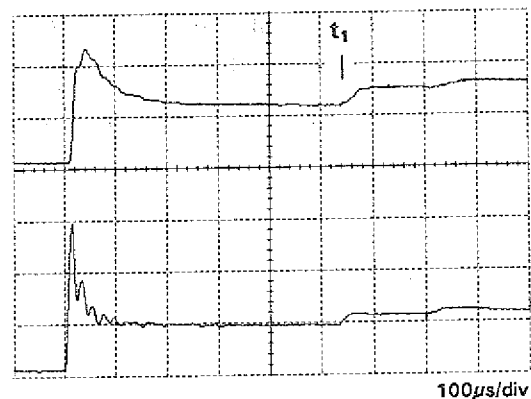


Figure 2:
Switching on procedure in presence of multipactoring,
upper trace: rf voltage in the resonator
lower trace: rf voltage at the output of the transmitter
 t_1 : 540 μs after start the system is switched to normal mode closed loop operation

After completion of the new rf system in September 1991, the important features were tested and optimized. The following items were of prime interest:

- stability of amplitude and phase;
- variation of frequency under load conditions;
- specified accelerating voltages;
- stable operation over the total frequency range of 20 to 30 MHz;
- interfacing between the rf system computer and COSY control.

As the improvement in stability of amplitude and frequency/phase had been a first order argument for the new rf system, especially these figures were checked carefully during the acceptance tests. Amplitude stability in normal operation, referred to the maximum accelerating voltage of $\hat{u}_{\text{max}} = 45\text{ kV}$

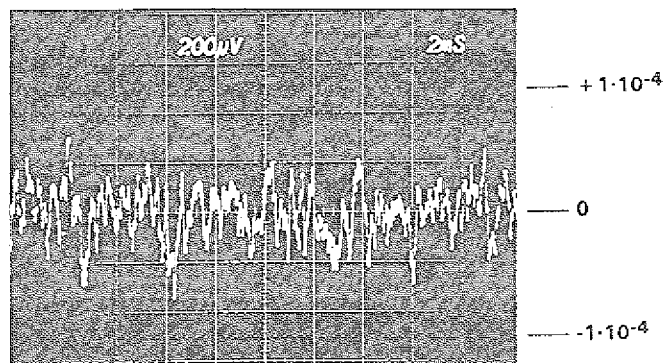


Figure 3:

Short term amplitude stability, scale referred to $\hat{u}_{\max} = 45$ kV, detection bandwidth 0-1 kHz, system settings: $f = 28$ MHz, $P_{\text{out}} = 36$ kW, $\hat{u} = 38$ kV

in our cyclotron was well below the specified figure of $\hat{u}/\hat{u}_{\max} \pm < 1 \cdot 10^{-4}$ for all settings of voltage and frequency in the tuning range. A typical result is given in Fig. 3. Even during frequency variation under load conditions the stability was $\hat{u}/\hat{u}_{\max} < \pm 4 \cdot 10^{-4}$ (specified amplitude stability during variation $\hat{u}/\hat{u}_{\max} < \pm 1 \cdot 10^{-3}$).

Phase stability referred to the output of the master oscillator proved to be $\Delta\varphi < \pm 0.3^\circ$ ($\Delta\varphi < \pm 1^\circ$ during frequency variation) and hence is well within the specified value of $\Delta\varphi < \pm 1^\circ$ (Fig. 4).

The maximum output power of 100 kW in the new system by far exceeds the requirements for normal operation. For ions of $Q/A = \frac{1}{2}$, accelerated to the maximum energy of $E/A = 45$ MeV/A, a voltage of $\hat{u} = 45$ kV is needed, which correspond to rf power of only 45 kW. But the above described switching on procedure is significantly simplified, when excessive power is available to overcome the multipactor influence.

The cyclotron was back for real beam acceleration in November 1991 after a shut down period of slightly more than two years. Experience with the new rf system under realistic conditions showed a number of weak points. Among others the adjustable coax-line of the frequency fine tuning system was damaged by overheating. Both watercooled vacuum capacitors, which terminate this line, failed and had to be replaced. The control computer system needs some debugging and modification. - So the contractor HERFURTH GmbH, Hamburg, is still in charge with the task.

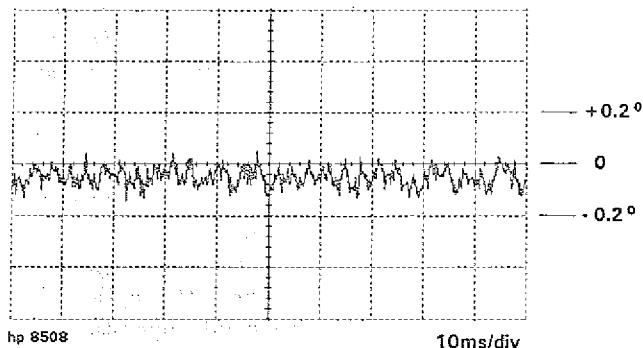


Figure 4:

Phase stability, system settings see Fig. 3

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Measurement of the Magnetic Field of the Cyclotron for the H_2^+ -Beam (40 MeV/A)

H.G. Böge, W. Bräutigam, R. Brings, F. Gehlen, H. Hadamek,
U. Rindfleisch, N. Rotert, P. Wucherer

Introduction: Starting 1992 the cyclotron will be needed as injector for the ring COSY-Jülich. For the first time only H_2^+ -ions of 80 MeV will be injected and stripped to 40 MeV p^+ . The trim coils of the cyclotron had to be redesigned for technical reasons, e.g. the number of coils has been reduced from 12 to 9. Therefore, also the physical design has been upgraded¹⁾. After repair of the failures in manufacturing²⁾ they have been installed in the cyclotron and measurements of the main field and trim coil fields carried out.

Geometry and accuracy: Since the geometry of the coils is changed new sets of currents are necessary to trim the field for isochronism which ideally means maximum energy gain and RF beam-phase (Offset) 0° on all 450 turns. The accuracy of field measurements is sufficient to get beam at extraction radius and to perform phase measurements and correction³⁾ for a better beam quality⁴⁾. The required accuracy of field measurements may be estimated in the following way: There are 12 phase probes in this cyclotron representing about 40 turns each. A tolerated phase deviation in between two probes of $\pm 20^\circ$ needs an accuracy in the integral field on the orbits of 10^{-3} . For the trimmed field one measurement for the main and two for each trim coil are necessary (see below).

The geometry of the yokes is not symmetric and causes a small field gradient in the direction "North-South" (see Fig. 1) which depends from the field level. A special control of the three largest trim coils (No. 1-3) allows to compensate this field gradient without changing the isochronous field shape. Therefore, the gradient can be neglected in the field measurements and the magnetic field taken as essentially 3-fold. The map ranges over 120° , Radius $R < 1.6$ m, the azimuths are defined to $\pm 60^\circ$ from the center of the hill sector "South". For the measured field level the hill (H) field is 1.8 and the valley (V) field 0.65 T. The hill gap is 8.4 and the valley gap 24 cm and the field gradient up to 0.25 T/cm. At the hill-valley edge deviations in the position of the field probe, therefore, can be measured easily (Fig. 2). The accuracy in reproducing the position resulting from 3 contributing measurements must be less than 0.05° . For a systematic valuation of beam profiles and their errors it is reasonable to achieve 0.005° (as in Fig. 3).

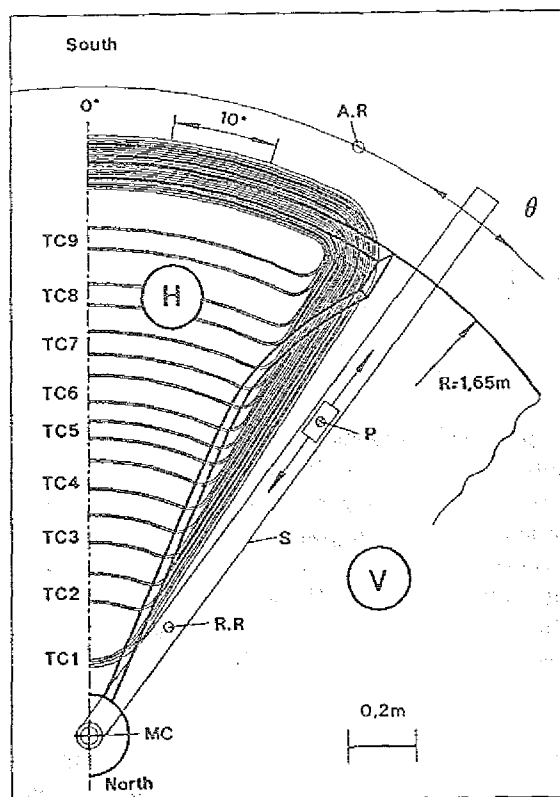


Figure 1: Part of the geometry of the field measurement, MC = machine and magnetic center, TC: trim coil, H: hill, V: valley sector, RR: radial reference, AR: azimuthal reference, S: support, azimuthally movable, P: probe on support, radially movable

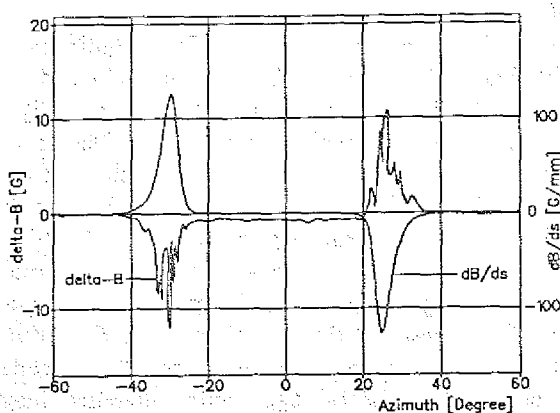


Figure 2: Azimuthal field gradient dB/ds (right hand axis) and field differences δB vs. azimuth (left hand axis) of two measurements with identical field. Example for position errors for the azimuthal resolution of 0.005°

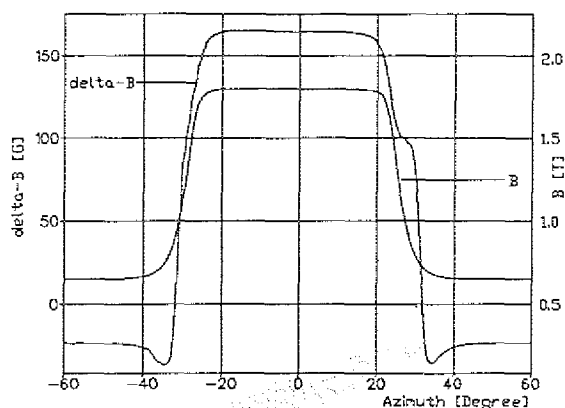


Figure 3:

Field differences ΔB ("azimuthal trim coil profil") for coil 1, excitation difference 960 Amp. turns and main field B vs. azimuth, $R = 0.8$ m

For practical reasons the mechanical system is adapted to polar coordinates. An azimuthal profile ($R = \text{const}$) is almost perpendicular to the hill-valley edge and clear in the graphical representation. Since the center of revolution coincides with the Machine Center (MC in Fig. 1) the azimuthal profiles can be evaluated e.g. in Fourier series without applying interpolation procedures which mix data errors and cannot be related to a singular set of measurements.

Concerning the grid of the map a compromise between the simple but time consuming $\Delta\theta = \text{const}$ and the best adapted $\Delta\theta \cdot R = \text{const}$ (constant gap!) is the used $\Delta\theta = 4, 2, 1, 0.5, 0.25$ degree. For $R = 1.5$ m the minimum grid constant is 8% of the hill gap.

Mechanics⁵⁾: Figure 1 schematically shows a part of the field mapping area with the hill pole face profile (H) which in reality is asymmetric (not shown). It generates the main field configuration which has to be corrected by the trim coils (TC 1 - TC 9). The Hall probe P is positioned radially in the azimuthally movable support S. The corresponding drives (not shown) act from the northern side via the Machine Center. The actual azimuthal position is recorded directly in units of 0.005° by a shaft encoder which fits in a 6 cm bore (normally for the tuning line of the RF-system) on the top side. The radial position is taken from the driving belt with a resolution of < 0.1 mm. The encoders are of the incremental type with reference marks which are positioned at R.R. and A.R. resp. (Fig. 1). These positions are related to their corresponding position/safety switch and the azimuthal switch is adjusted to be 60° from the hill centerline. This centerline coincides within 0.1° with the centerline of the symmetric part of the hill sector (Radius < 0.5 m). The scaling of the azimuthal position is given by the encoder resolution. The absolute value is not relevant for isochronism.

Electronics and software: The control system for the measurements was developed as part of an engineers certificate⁵⁾. A Personal Computer (80286) is used for control and on line data handling. The PC reads the actual positions from the encoders (minirot 450) equipped with the corresponding counter electronics and actuates the drives until the programmed position is reached. The positioning procedure and the mechanics have been optimized to keep within the azimuthal resolution (0.005°) and minimize the time for one field value (7 sec/point). The magnetic field was taken with a Hall probe (BH 15, resolution 10^{-5} T, absolute accuracy 10^{-4} T) via an interface. The PC-Code performs the mechanical calibration with the aid of the operator which types the parameters (fixed coordinate, coil no.,...) and starts the measurement. One measurement is defined as only one set of data for one fixed coordinate. Having finished the measurement successfully the PC stores the data with a programmed filename in which the parameters are keyed.

Measurement and results: The main field was measured at a level of 0.65 T (NMR probe in the valley) mainly in the azimuthal sequence ($R = \text{const}$). The azimuthal grid was identical with the following trim coil measurements. The steps of $\Delta R = 4$ cm in the radius in between the sets of measurements for $R = \text{const}$ was adapted to the subsequent coil measurements ($\Delta R = 4, 8, 12$ cm). For max. positive excitation of a coil the field difference to the main field can then be calculated and displayed. It turned out that a second measurement with max. negative excitation of the trim coil enables results with much better accuracy and reliability as a pure factor of two. One reason is that the trim coil fields modulate the main field permanently. (A cancellation has not been tried.) The other reason is that the main and the coil field measurement are taken at different times (ca. 1 week difference) and the conditions did not keep constant: drifts, probe and mechanical accidents,...). The main error source was an unreliable connection between the axis of rotation and the shaft encoder. The measurements started with a reproducibility of 0.005° (Fig. 2) and ended with up to 0.04° . (The time schedule did not allow for repetition.)

Figure 3 shows an example of an azimuthal profile of the main field (B) together with that of coil No. 1 (ΔB) for $R = 0.8$ m at optimal conditions. The B - and ΔB -profiles look very similar but the scaling and the negative values of ΔB in the valleys (stray field) make the main difference!. There are two slopes of ΔB only for positive θ (beam exit side) which are caused by the fact that the coil here extends 12° into the valley gap.

Figure 4 shows an example for the azimuthal profile of the stray field in the hill (-4 mT) as well as in

the valley (-1 mT) for coil No. 7, $R = 1.08$ m and the approximations "sharp edge" (S) with 4 and Fourier series (F) with 13 parameters. Both approximations are suitable to check the consistence of data by looking on the radial dependance of the parameters.

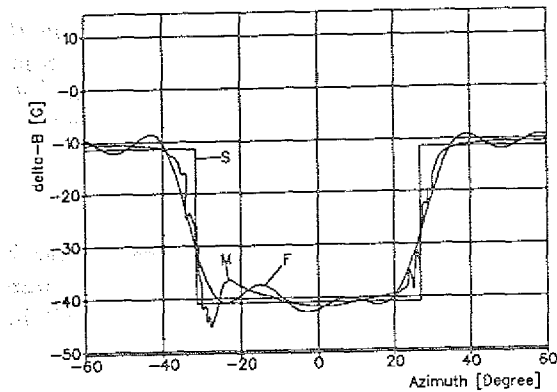


Figure 4:

Delta-B as in Fig. 3 but in the stray field region of coil 7 at $R = 1.08$ m. M: measured, S: sharp edge, F: Fourier approximation

Figure 5 shows the radial profiles (delta-B) of the fields of all 9 coils for $\theta = 0^\circ$ which are very similar to the mean value in the Fourier series or the hill value in the sharp edge approximation. The ("not wanted", negative) stray field in comparison with the ("wanted", positive) field inside a coil increases with the coil number. The phase shift caused by a coil is the integral of the mean of delta-B times the beam turn number¹⁾. For the smallest coil No. 9 the stray field contribution dominates which means that the larger coils (compare Fig. 1) have to compensate this stray field tail. The trim coil phase shifts are strongly coupled with each other which means practically that the beam phase can not be optimized without the aid of a computer.

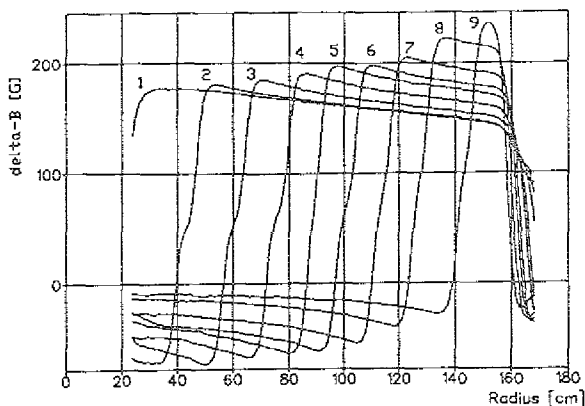


Figure 5:

Delta-B as in Fig. 3 but vs. R ("radial trim coil profiles") along the hill center line, $\theta = 0^\circ$, for all nine coils, see text

One set of trim coil currents has been calculated by applying (not corrected) Fourier parameters and tested with beam: Only one coil had to be changed by 18% of its max. excitation to accelerate the internal beam to extraction radius.

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The Injection Beam Line

P. Jahn, R. Maier, J. Reich, M. Rogge

The injection beam line has to transport the beam over a distance of 95 m from the cyclotron to the COSY ring. Almost all parts of the injection beam line have been installed by the Siemens company and set to work in their main functions. The (H^-)-branch is built up now. Some of the mechanical elements are needed for the (H_2^+)-branch, too, which is prepared and will be used first with beam.

The dipole and quadrupole magnets had to be readjusted to compensate floor settings of up to more than 1 mm due to shielding walls. The positions of the diagnostic devices remained unchanged, being readjusted later electronically. The vacuum system of the beam line is divided by valves into three sections. The pressure is below 10^{-7} mbar everywhere. Heating out and UHV-tests of the last part will be done during the next weeks. The control system has been installed and front-end tests are being performed.

The 57 main magnets connected to 19 power supplies were checked with regard to field strength and polarity. Both polarities are necessary because of particles with different charge signs. The steering magnets have been set to work as well. They are able to deflect a beam with the maximum rigidity of 2 Tm by ± 2 mrad in the horizontal and vertical planes.

The first two dipoles were investigated more thoroughly. The maximum field-strength B is limited by the power supply to 1.51 T. The remanence measured after 3 min is ± 0.6 mT. The

field values calculated by the control system on the basis of an empirical $B(I)$ table agree within 5 mT with the measured values. Comparison of the fields of the two identical magnets connected in series at the maximum current give an agreement better than 25 mT.

The thermal security of all magnets was proved by running them with full power and without water cooling. All power supplies involved were switched off after typically 20 minutes at about 70 degrees with one exception at 90 degrees.

Ion-optical calculations have been done in order to study beam geometries beyond the reference design. So, restrictions concerning the sign of the dispersion at the stripping target have been removed by applying an additional mode of operation in the 8-element FODO-structure.

An application program has been written, which couples the field-strengths of the quadrupoles in a linear way together, in order to vary focusing properties keeping other ion optical conditions constant. Such conditions can, for example, be: "focusing in the horizontal plane" and/or "focusing in the vertical plane", "magnification = 1" etc. A similar program is prepared for the setting of the steering dipoles¹⁾.

Reference:

- 1) P. Jahn, M. Rogge, Beam Steering in the Injection Transfer Line, Contribution to this Annual Report

Beam Steering in the Injection Transfer Line

P. Jahn, M. Rogge

The possibilities of beam steering in the injection beam line¹⁾ have been investigated by calculations using the computer program TRANSPORT.

Deviations of the beam and their corrections by steering are of special importance at the beginning and at the end of the beam line. The beam leaving the cyclotron has to be adjusted to the axis of the beamline and the beam entering the COSY ring has to be adapted in position and angle on the stripping target.

As an example the possibilities for steering in the horizontal plane of the last ion optical sections are demonstrated in Fig. 1 for one of the standard ion optical settings. The horizontal beam geometry involves a phase space of 3.18π mm mrad and a beta value of 31.8 m at the stripping target. The deviations of the central trajectory from its normal path induced by a $+1$ mrad kick from each of the steerers or deflection magnets are plotted (dashed lines). The effects of several steerers can be linearly superimposed.

In the representation chosen the deviations are compared with an effective aperture, which is the mechanical aperture diminished by halve the beam dimension on both sides. In this case the touching of these deviations with the drawn boundary (solid

line) is equivalent to the touching of the beam envelope with the real pipe aperture.

Critical points, where the free aperture is narrow, show up directly. A linear combination of kicks from selected steerers can be quantitatively estimated from the drawing, to allow a passage of the beam through the critical aperture.

The geometrical limitations are especially stringent in the last 4 m of the beam line. The calculations show, that for these ion optical conditions horizontal angles of more than about ± 3 mrad at the stripping target cannot be realized.

An application program is in progress, which allows to change the currents in the steerer magnets as a linear function of three free input parameters to satisfy up to three prescribed ion optical conditions. So it becomes for example possible to change the angle at the target by one "knob" keeping the position of the beam at the target constant and to change the position on the target by an other "knob" keeping the angle constant.

Reference:

- 1) J. Reich, M. Rogge, IKP Annual Report 1988, Jül-Spez-499 (1989), p. 129

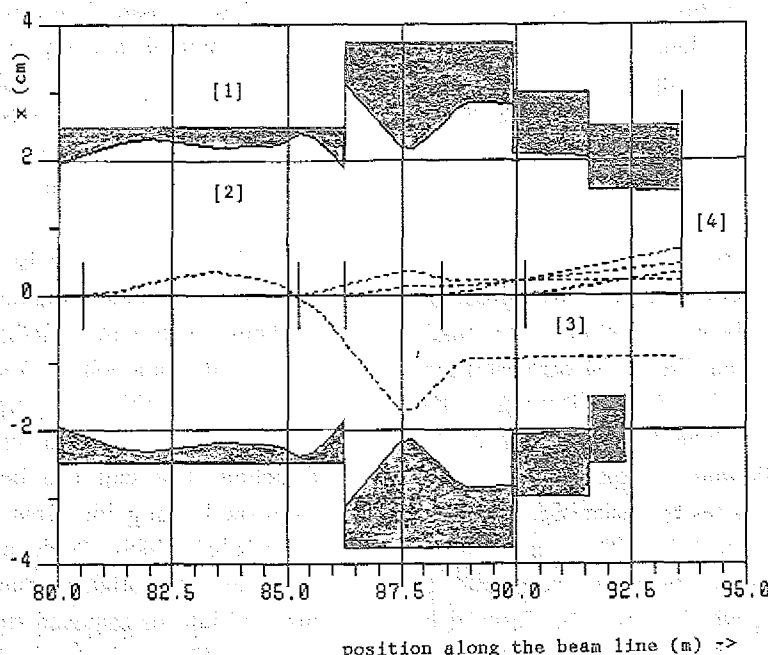


Figure 1:

Geometrical limitations, beam extension, and steering possibilities in the last section of the injection beam line.

- [1] mechanical limitations (beam tube etc.)
- [2] aperture for the central trajectory which equals [1] reduced by the beam extension (shaded area)
- [3] central trajectories deflected by one of five steerers indicated at the axis; the angle is $+1$ mrad in all cases
- [4] stripping target

Measurement of the COSY Injection Energy

H. Labus, J. Dietrich, N. Dolfus, R. Maier

All components to perform the measurement of the COSY injection energy are now available and put into operation. The measurement principle is based on time of flight measurement as applied in other labs, e.g. GSI Darmstadt. The signals of two capacitive probes at the beginning and end of a $L=28,9$ m long straight section are amplified by 48 db with 50 Ω RF-amplifiers and then transmitted to the control room by 125 m coaxcables of type RG 273/U. The signal propagation times are matched to less than 0,5 ns. The RF-amplifiers are connected with the probes via 2,5 m semirigid coaxcables to reduce radiation damage of the RF-amplifiers. To get the energy $E=m_0c^2\sqrt{1-v^2/c^2}$ it is necessary to measure the bunch velocity $v=L/T$, where L is the constant flight distance and T the variable flight time between the probes. $T=T_i+T_f$ consists of two parts: The integer part T_i is precisely defined by the number N of bunches, which are aligned along L and by the time difference $1/f$ of two bunches, where f is the cyclotron frequency. At the standard COSY-injection with 80 MeV H_2^+ at $f=27,8$ MHz we get $N=8$, where the precision of T_i is $<10^{-6}$ due to the

frequency stability. The fractional part $T_f < 1/f$ is the time difference between an arbitrary signal from the first probe and the next following signal of the second probe. T_f will be measured in a preliminary setup with a fast digital storage scope (TEK 2440) with a precision of $<0,5$ ns, so energy resolution of $\pm 0,3$ % is available. This meets the requirements to check if the injection beam fits to the COSY acceptance in momentum space.

At a final setup the scope will be replaced by a VXI module (HP E1428A) with larger storage depth. Appropriate fit procedures, applied to the signal streams of both probes, are investigated to find the phase difference, e.g. the fractional part T_f with resolution <0.1 ns. A main advantage of this method is that the energy measurement only depends on the precision of the cyclotron frequency and the probe distance, whereas the calibration of the scope has no influence, as well as signal slope degradation between the probes.

Beam Profile Measurement in the COSY-Extraction

H. Labus, J. Dietrich, R. Maier, W. Renfle

Beam profile and position measurements are performed at 14 positions along the extraction beam lines with multiwire proportional chambers (MWPC), as developed at GSI Darmstadt. Each station consists of a detector unit and a mechanical drive.

A detector unit has an active area of 64×64 mm² and is able to record two profiles in x- and y-direction, perpendicular to the beam direction. The space resolution is <1 mm according to 64 wires at 1 mm distance for each plane. The inner diameter of the beam pipe is 100 mm. The free diameter of the detector frame is >100 mm. The gas amplification is controlled within 5 decades by variation of the high voltage between 50 V and 5 KV. Thus, profiles may be recorded at current densities from 10^6 to 10^{11} 1/s cm², a range, which cannot be covered by any other measurement principle. The smallest time window to record one profile is 200 μ s, so time dependent oscillation of the beam center may be measured up to frequencies of some kHz. Larger integration times of 2 ms or 20 ms can be selected by software control. Thus, statistical current deviations

at the lowest beam densities can be smoothed. Each detector unit is a compact module consisting two MWPC's, individual preamplifiers for each wire with current readout. An onboard digital processor collects the data and serves the IEEE-488-Interface. A prototype is under test.

A mechanical drive consists of a 400 mm deep cup with 140 mm inner diameter and 0.5 mm wall thickness, whose outer ledge is connected with the beam guide via a bellows. A stepping motor control is able to move this cup perpendicular to the beam direction with ± 0.2 mm reproducibility in x- and y-direction. The cup can be pulled out completely from the beam guide cross section. The mechanical drive is bakeable to 200°C and was developed at the ZAT of KFA Jülich. The detector unit can be inserted into or removed from the cup very quickly. Thus, baking procedures or exchange of defect detector units can be performed without vacuum distortion. Even rupture of an MWPC window does not stop an experiment in progress. Two prototypes have been built and successfully tested.

H₂⁺-Source:

The low-emittance H₂⁺ source described last year¹⁾ has in the meantime injected particles into the cyclotron. Up to 150 μ A of H₂⁺ ions with an energy of 8.2 keV - needed for a cyclotron extraction energy of 80 MeV - have been transported to the hyperboloid inflector in the center of the cyclotron with almost 100% transmission in the source beamline and the vertical cyclotron injection path. The effect of the low beam emittance on the cyclotron transmission could not be determined so far since the beam has not yet been extracted from the cyclotron. The source has shown reliable and stable performance over several weeks of operation.

The beamline has been connected to the new COSY computer control system and so shall the source. Some improvements of the control system are still necessary to guarantee reliable handling of the source and the source beam line. The pulsing system¹⁾ driven by an amplifier built by the electronic workshop of the IKP can produce beam pulses with a minimum rise time of less than 150 ns. The pulse length can be as short as 1 μ s.

H⁻-Source:

The H₂⁺-source can be replaced by an H⁻, D⁻-source¹⁾ which has been bought from the IBA company in Louvain-la-Neuve (Belgium). Both sources have been mounted on waggons and can be easily exchanged without disconnecting any electrical cables or cooling hoses. This source exchange would later on not waste COSY beam time since a switch from positive to negative ions needs much more time for changes at the COSY injection.

The H⁻-source has been designed by the company to run at an extraction voltage of 25 to 30 kV. At an extraction voltage of only 4.5 kV needed here some optical problems occurred in the extraction area. The result was an incomplete matching of the source to the existing beam line, caused by the large divergence of the low energy beam. The matching has been improved remarkably by introducing two additional electrostatic electrodes behind the puller. The source with the full extraction system is schematically shown in Fig. 1. The first new electrode is on negative voltage up to several kV whereas the second one is grounded.

After a period of cyclotron operation with H₂⁺ particles, which includes the acceptance test of the new transport beam line between the cyclotron and

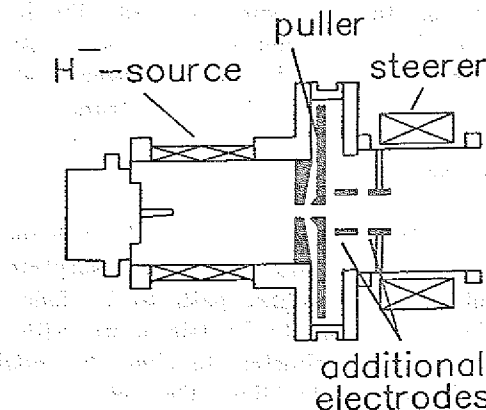


Figure 1:

Modified extraction system of the H⁻-source

the COSY ring, the H⁻-source shall be connected to the cyclotron to study the cyclotron performance for H⁻ particles. These tests have to be performed as soon as possible since they may result in changes in the cyclotron extraction region to achieve optimum performance for polarized H⁻ particles²⁾.

14 GHz-ECR-Source:

Ion implantation measurements with low energy highly charged ⁴⁰Ar and ⁸⁴Kr ions on silicon targets have been performed by R. Herrmann from the university of Frankfurt. The purpose of this work was to compare the implantation properties of singly and multiply charged ions. M. Loiselet from the university of Louvain-la-Neuve (Belgium), who is involved in a project to build an ECR source for radioactive elements, has started to determine the lower limit of gas consumption in the Jülich ECR source necessary for the efficient production of highly charged ions. A preliminary result is that the source can produce several μ A of ²⁰Ne⁸⁺ with a Ne gas leak of about $1 \cdot 10^{-5}$ mbar $\cdot$ l s⁻¹.

Further on it has been shown that the source can work with a biased plate or an Al₂O₃ layer on an Al backing in state of a first plasma stage²⁾. Studies to compare the source performance in the three different set ups - first stage, biased plate and Al₂O₃ layer - are in progress.

References:

- 1) H. Beuscher, W. Hoffmeister, A. Müller, H. Hadamek, U. Rindfleisch, IKP Annual Report 1990, Jül-2462, p. 175 and p. 197
- 2) H. Beuscher, Proc. of the 4th Int. Conf. on Ion Sources, Rev. of Scient. Instr. to be published

Magnetic Spectrometer BIG KARL

J. Engel, H. Machner, R. Maier, P. von Rossen, P. Schmidt, R. Tölle

The year gone by has seen a steady and substantial progress in all fields pursued to ready the large spectrometer for the upcoming experiments. After extensive preparative work, that had been stretching for over a year, the 200 ton instrument was anchored to its new position with a precision in the millimeter range.

Exact and in depth studies that preceded this move assured a smooth execution of all necessary steps. Besides others the complex path to be followed during that move had to be laid down with an accuracy of some centimeters to obey the restrictions given by the building. One of the most expensive preparations was the reinforcement of the floor along the taken path, which had to withstand near point-like loads of 90 tons. In addition the

central concrete shielding weighing 150 tons had to be completely dismantled beforehand. The same was holding for the former power supply gallery that was then erected adjacent the south wall of the Nord Halle after undergoing appropriate modifications.

The moving of the heavy instrument was done using its air bearings and the hall crane, which was lifting the central area around which the spectrometer usually rotates. The actual course was steered with two hydraulic driven pistons, that held it precisely on the marked path to its new location.

The concept for the experimental cave structure that was based on investigations to ensure radiation protection outside has been laid down. Figure 1 depicts the area.

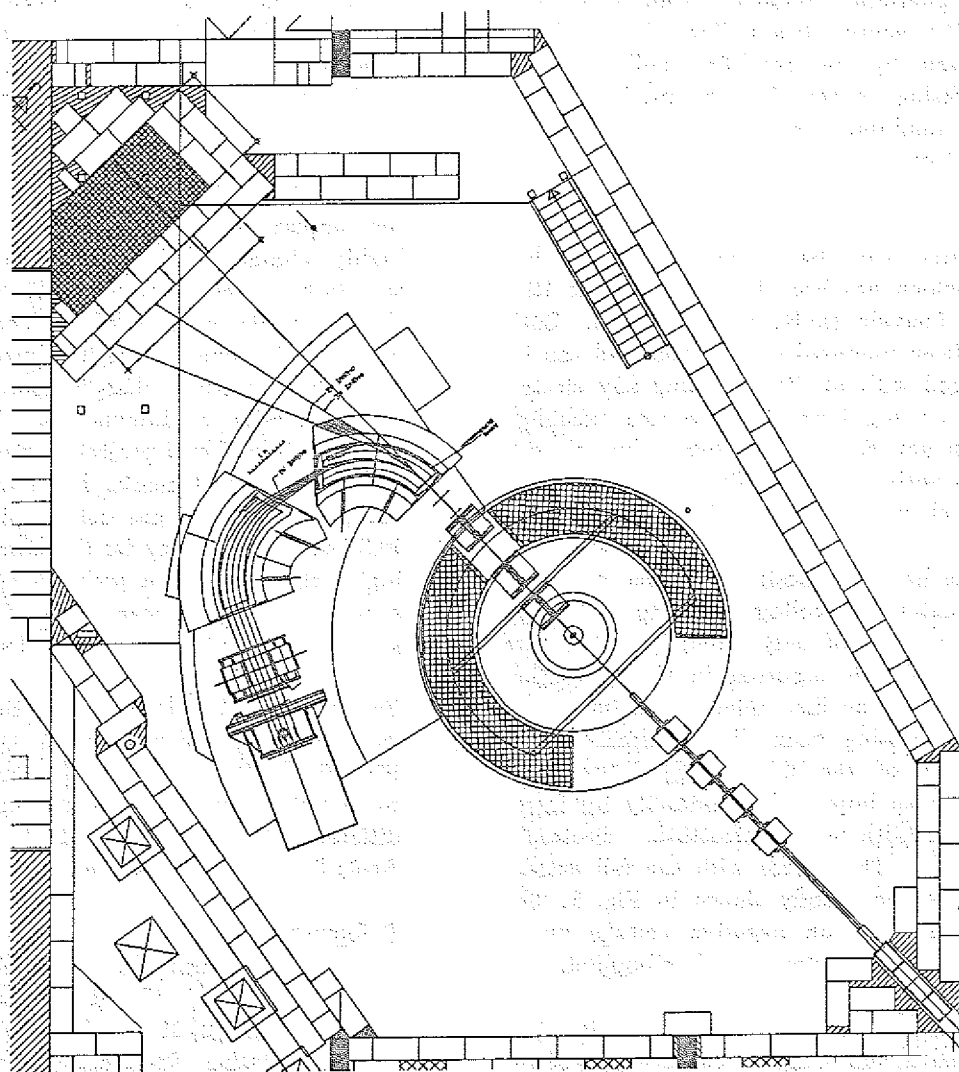


Figure 1: Shielded experimental area for the spectrometer BIG KARL. The area is being entered through the labyrinth shown at the top. For more information see text.

The extracted beam is entering from the lower right and is focused with four quadrupoles acting as a symmetric triplet onto target. The spectrometer is shown under zero degree, a position that will be used for most of the planned experiments. The direct beam enters the spectrometer and leaves it via a thin vacuum window and a hole in the outer side yoke of the first dipole. A massive beam dump finally absorbs the high energy particles. The immediate target area is surrounded by an additional circular concrete shielding and can be accessed from the rear. According to these plans a formal request for permission of realization has been filed to the authorities.

The ionoptical modernization of the entrance region which was into discussion for quite a time will be now realized with three high power, high field gradient quadrupoles that were given to us on loan from DESY. The systems are now on site and have already been successfully tested with respect to their electric, cooling, and vacuum properties. Taking their specifications, extensive calculations were done to optimize acceptance and the highly desired reconstruction of particle trajectories by measuring position and angle in the focal plane of the spectrometer.

The higher power demand of those new quadrupoles will be met in the following way: For the first

quadrupole the characteristics of the new power supply were specified and the order was placed, the second quadrupole will be powered by paralleling the existing power supplies of the old entrance region, the third quadrupole is fed with a 1000 A power supply that was freed by swapping.

The engineering study to exit the direct beam out of dipole 1 has been completed. It comprised the modification of the vacuum chamber and the side yoke of the first dipole as well as the steps to dismount and to reassemble all parts. Commission was given to an external company to modify and to manufacture the parts according to the engineering specifications.

Development and set-up of the components for the BIG KARL control system have been carried on. Racks containing the various modules have been assembled and were installed at the spectrometer support and on the power supply gallery. Tests of VME and G64 systems were conducted.

Work on the focal plane detectors has been progressed in parallel by our colleagues in Krakow. The first detector planes have been fabricated and single wire readout was successfully performed. Some of the topics cited here are covered in more depth by special articles in this report.

Studying the Ion-Optical Bandwidth of the QQQDDQ-System

M. Kistryn, P. von Rossen, R. Tölle

When investigating ejectiles with BIG KARL close to zero degree it is in many instances a foremost concern to obtain besides the momentum also the precise scattering angle by measuring position and angle in the focal plane.

As the optimization of one property often adversely affects the others, one has to evaluate the bandwidth of qualities that can be covered by varying the entrance region of the QQQDDQ-System. Relevant qualities are: momentum resolution, momentum range, maximum momentum, horizontal and vertical acceptance, entrance angle reconstruction, behavior of the direct beam (size at the exit of dipole 1), envelope size for a start trigger e.g. after Q2.

The restrictions that had to be observed were: the dipoles with a gap of 60 mm and the dispersion quadrupole being fixed at their original positions, the maximum width of the usable field for each dipole, the butterfly-shaped-apertures of the entrance quadrupoles and the dispersion quadrupole with its flat vacuum chamber (gap ca. 70 mm), the dipole correction windings allowing only limited radial field deformations (quadratic term in field expansion), not every position being available for the focal plane detectors, and there is an upper limit for the pole-tip-field of the quadrupoles.

Our study investigates how to optimize specific properties of the QQQDDQ-System. The following degrees of freedom were evaluated in looking for the classes of sets that give the best results:

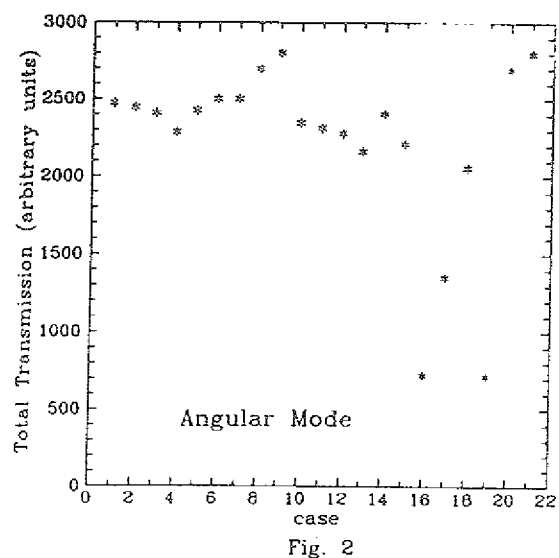
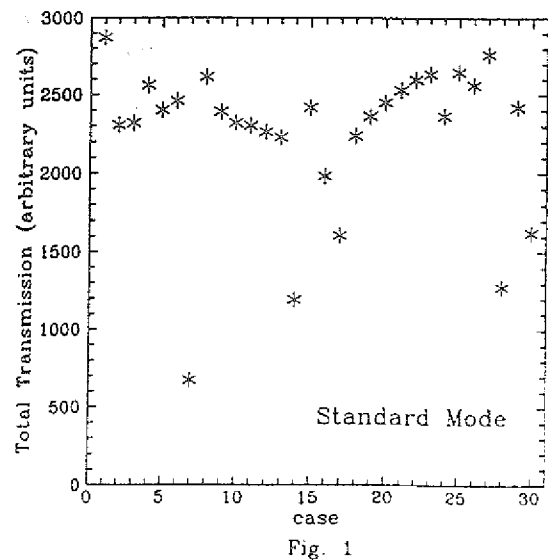
- I. position of the first quadrupole Q1,
- II. distance between the first and the second quadrupole,
- III. position of the focal plane,
- IV. polarity and value of the quadrupole Q2A (only for the standard mode, because for the angular mode this quantity is fitted),
- V. polarity and value of the quadrupole Q3.

The calculations for angular and standard mode were made using the programs TRANSPORT and TURTLE. The evaluation proceeds along the following steps:

- a) decide on design parameters
- b) set up TRANSPORT input file
- c) fit the settings for the lenses

- d) put values into TURTLE input set
- e) find overall acceptance between -4% to 4%
- f) show acceptance in x , x' and y , y' .

The results are summarized in figures and in tables. Figures 1 and 2 show the total transmission for both modes separately, for all calculated cases (shown in Table 1 and Table 2). For all cases the calculations were performed with 5000 particles in a phase space of $x' \pm 30$ mrad and $y' \pm 120$ mrad with a central momentum of 1.1 GeV/c. Figure 3 depicts the total transmission versus the distance of the quadrupoles Q1 and Q2.



case	total trans- mission	Q1 [kGs]	Q2A [kGs]	Q3 [kGs]	distance Q1-Q2 [cm]	distance target-Q1 [m]	distance Q3-detector [m]
1	2871	16.59	0	-6	22	0.4	1.71401
2	2308	14.07	0	-6	22	0.62	0.71401
3	2324	14.03	0	-6	22	0.62	0.76401
4	2564	13.58	0	-6	22	0.62	1.71401
5	2402	13.57	0	-6	22	0.62	1.01401
6	2461	13.77	0	-6	22	0.62	1.21401
7	672	13.16	0	-6	22	0.62	3.71401
8	2587	13.33	0	-6	22	0.62	2.71401
9	2387	12.75	0	-6	40	0.58	2.14501
10	2318	12.35	0	-6	40	0.62	2.14501
11	2301	12.22	0	-4	40	0.62	2.14501
12	2264	12.04	0	-2	40	0.62	2.14501
13	2232	11.83	0	0	40	0.62	2.14501
14	1187	11.3	-4	-6	40	0.62	2.14501
15	2424	12.14	-2	-6	40	0.62	2.14501
16	1988	12.45	2	-6	40	0.62	2.14501
17	1612	12.5	4	-6	40	0.62	2.14501
18	2242	12.17	0	-6	44	0.62	2.14501
19	2370	12.56	0	-6	36	0.62	2.14501
20	2457	12.78	0	-6	32	0.62	2.14501
21	2535	13.03	0	-6	28	0.62	2.14501
22	2603	13.31	0	-6	24	0.62	2.14501
23	2635	13.46	0	-6	22	0.62	2.14501
24	2365	13.95	0	-6	22	0.62	0.88901
25	2649	13.99	0	-6	22	0.58	1.7401
26	2564	13.58	0	-6	22	0.62	1.7401
27	2766	14.97	0	-6	22	0.50	1.7401
28*	1274	11.3	-4	-6	40	0.62	2.14501
29*	2424	12.14	-2	-6	40	0.62	2.14501
30*	1626	12.5	4	-6	40	0.62	2.14501

Table 1:

List for all calculated cases for standard mode.

All the data (except * cases) in the table is calculated with 1 m distance between Q2A and first dipole, for * cases with 1.1 m.

case	total trans- mission	Q1 [kGs]	Q3 [kGs]	distance Q1-Q2 [cm]	distance target-Q1 [m]
1	2473	14.3	0	22	0.62
2	2446	14.1	0	26	0.62
3	2410	13.8	0	30	0.62
4	2287	13.2	0	40	0.62
5	2427	14.3	2	22	0.62
6	2504	14.3	-2	22	0.62
7	2504	14.3	-4	22	0.62
8	2700	15.7	0	22	0.50
9	2802	17.3	0	22	0.40

case	total trans- mission	Q1 [kGs]	Q3 [kGs]	distance Q1-Q2 [cm]	distance target-Q1 [m]
10	2350	14.5	0	22	0.62
11	2315	14.2	0	26	0.62
12	2281	13.9	0	30	0.62
13	2167	13.34	0	40	0.62
14	2410	14.50	2	22	0.62
15	2214	14.66	-2	22	0.62
16	723	13.03	-4	22	0.62
17	1352	13.03	-3.5	22	0.62
18	2061	*14.74	-3	22	0.62
19	719	13.02	-6	22	0.62
20	2699	15.92	0	22	0.50
21	2805	17.53	0	22	0.40

Table 2:

List for all calculated cases for angular mode.

The data above was calculated with 1 m distance between Q2A and first dipole; the data below with 1.1 m.

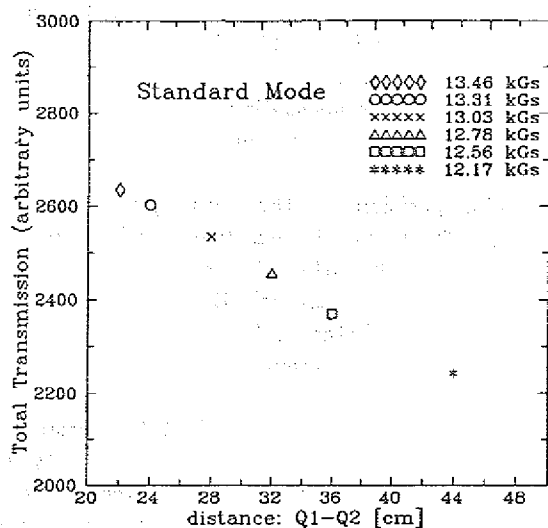


Fig. 3

The transmission was also calculated for increasing momentum deviation. Defining the transmission of the 0% momentum deviation i.e. the central momentum for a specific case to be 100% and dividing the calculated numbers for the other momentum bins by the one for the 0% momentum deviation bin one obtains the relative transmission shown in Fig.4.

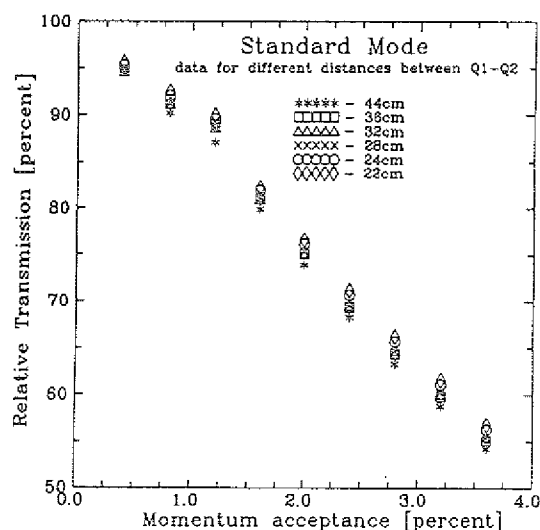


Fig. 4

To choose the best set-up for a given experiment one has to look at both kinds of plots, the relative transmission and the total one. If for a given set-up the value of the first quadrupole, given in the tables, is higher than 12.8 kG one can use these parameters but with the restriction that the momentum of the observed particles has to be correspondingly lower than 1.1 GeV/c.

Conclusions:

I. The distance between target and first quadrupole Q1 was changed within the range 40 - 62 cm for both standard and angular mode. One should choose the smallest possible distance: "target-Q1" to get highest transmission for both modes.

II. The distance between the first and the second quadrupole was varied within the range 22 - 44 cm for standard mode and 22 - 40 cm for angular mode. Again, the smallest distance is the best if one looks for the best transmission.

III. The position of the focal plane was also modified. The value of total transmission does not change very much within the range 0.9 m - 2.8 m (middle of Q3 - - detector). It is somewhat better to have a larger distance, but at 3.7 m the transmission has fallen off steeply.

IV. The value of the quadrupole Q2A usually was 0 kG for the standard mode. This quantity was always fitted for the angular mode. Hence, for the standard mode calculations were performed to see how the value of the Q2A quadrupole influences the

transmission. Computations used values of -4, -2, 0, 2, 4 kG for the Q2A. High transmission is obtained for -2 kG and 0 kG. Compared to the -2 kG transmission, the 2 kG transmission is 15% lower, the 4 kG transmission falls down by 30%, and with Q2A at -4 kG it is halved.

V. The value of the quadrupole Q3 was changed within the ranges -6 to 0 kG and -4 to 2 kG for standard and angular mode, respectively. This quality does not significantly affect total transmission.

In general it was found, that there is smooth behavior of the ion optical parameters investigated. This gives some freedom to shape one parameter without too heavy tradeoffs at the others.

From the experience the kind and the quantity of the radiation protection duties in the IKP are largely defined by the prescriptions named in the radiation protection ordinance (among other things: determinations of radioactivity; contamination investigations; radiation protection informations; dose rate measurements; control of the radiation exposure, of the release of the air born radioactivity and of the dose rates in radiation protection areas; organisation of the radioactive transports; execution of the conditions in connection with efforts of foreign crews in controlled areas and their care and supervision; order of the radiation protection devices; regular function tests of the movable and fixed radiation protection and personnel safety system devices; collection and declaration of the radioactive waste; planning of the working order and surveillance in case of handling with radioactive materials; markings; dispatch of the results of the air monitoring; regard for conditions of authorizations and internal prescriptions; furnishing with diverse personal dosimeters; erection of temporary and permanent shielding; decontaminations; different notes and reports accordingly to the radiation protection ordinance; storing and safeguarding of the radioactive material; execution of the measures for the starting and operation of x-ray devices; closings; control of restrictions and inhibitions in connection with the handling with radioactive material; directives for stay restrictions; informing about news and advances on the radiation protection field; informing about the status of the lawful regulations, instructions and recommendations on the radiation protection field), and the working hours being at disposal are filled up.

Yet for the setting in operation of the isochronous cyclotron as COSY-injector and of the beam line a series of labours had to be carried out additionally in 1991: the reconstruction of the cyclotron necessitated an essential modification of the personnel safety system (PSA) of the cyclotron. By that the total "hardware" (bells, horns, rotary lamps, quit switches, emergency switches) was reinstalled and the software of the cyclotron-PSA was created and implemented in the Simatic S5 so that the PSA could be set in operation before the start of the cyclotron operation.

For the acceptance of the injection beam line from the cyclotron to COSY a series of radiation protection duties had to be carried out: with the object of the licence of erecting comprehensive statements, particularly about shielding, air and cooling water activation, release of radioactivity, security measures and safety considerations, were submitted. The licence of erecting was given on September 5, 1991. For the test operation of the injection beam line a motion to a further licence had to be proposed. Thereto further statements, particularly about the organisation in the IKP and the PSA of the injection beam line, were given to the competent authority. The limited test operation licence was given on October 9, 1991. Beside the administrative duties above all the realization of the measures necessary for the operation of the PSA is to be mentioned: the "hardware" inclusive the interrupt devices was installed and wired. The software was created by taking into account special demands: it must be possible to operate the cyclotron and the beam line separately and to include the operation of COSY, extraction beam line and external experimental areas in future. The operation of the cyclotron referring to the PSA is controlled by the new software in the meantime.

Because during operation of the beam line higher beam losses can occur as assumed in the shielding calculations it is necessary to control the dose rates on accessible places continuously. For this reason along the injection tunnel on the outside of the shielding walls a series of neutron monitors well suitable for the dose rate measurements were installed and wired. The software needed because of the especially kind of the radiation monitoring (the mean permissible dose rate is very low; as intervall for average one hour can be choosen; for an impending excess of the permissible dose rate a "prewarning" should occur) was created and tested. The radiation monitoring system was set in operation.

Moreover a system for the monitoring of the radioactivity released with air was selected (a very large sensitivity was required) and ordered. In the meantime the system was delivered.

Computer Aided Shield Layout at COSY

D. Filges, H.R. Langohr, H.J. Probst, B. Wolfertz *

*Bergische Universität GHS Wuppertal

High-energy radiation transport codes, namely the HERMES system [1] are being applied to investigate some fundamental shielding problems associated with COSY. It is not practical to run such a large Monte Carlo particle transport system for detailed calculations to address all of the numerous radiation safety questions and engineering type design problems associated with experimental areas and accelerator shielding.

Therefore, we have written an initial version of a simple, fast running code, named CASL for Computer Aided Shield Layout, which has rather general source, material and geometry capabilities, but which requires minimal computer and setup time. While the CASL code uses approximate, semiempirical physics treatment, it can provide sufficient accuracy for many practical problems.

The basic calculational approach of the CASL code is to assume a simple, exponential attenuation relation for the shield attenuation of high energy radiation. That means

$$D(\vec{r}) = \frac{(\tilde{D} \cdot \dot{S}_n(\vec{\Omega} \cdot \vec{\Omega}_0; E_0, A_t) B(E_0, t) \exp[-t/\lambda(E_0, t)])}{\bar{F}_h |\vec{r} - \vec{r}_0|^2}$$

$D(\vec{r})$ = dose-equivalent rate (location $\vec{r}$, direction $\vec{\Omega}$), from a proton beam (energy E_0 , direction $\vec{\Omega}_0$), target material (A_t , location $\vec{r}_0$).

The high energy neutron flux at material thickness t is

$$\phi_h(t) = \frac{B \cdot \exp(-t/\lambda)}{|\vec{r} - \vec{r}_0|^2}$$

B = "buildup factor", λ = attenuation length.

The dose rate due only to the high-energy neutrons is then $\phi_h(t)/\bar{F}_h$, where $\bar{F}_h$ is the average flux-to-dose factor for neutrons > 100 MeV.

The dose is determined predominately by the low energy neutrons. Therefore, the dose-ratio $\tilde{D}$ (total dose/dose due to neutrons > 100 MeV) is applied. $\tilde{D}$ depends on the neutron spectra of the particular shield material used obtained in detailed transport calculations.

To illustrate the application of the code the shield design of a typical beam stop arrangement for COSY

is applied. Values for the following parameters are assumed:

proton beam energy: $E = 1$ GeV
average beam current: 1.6×10^{-7} mA = 10^9 p/s
target material: iron.

The dimensions and materials for the beam stop are given together with some iso-dose line plots in the x-y plane in Fig. 1. The proton beam hits the iron target inside the re-entry hole at $U = -125$ cm. Fig. 2 shows the attenuation of the fast neutron flux and the dose dependence ($\tilde{D}$) through the beam stop in zero degree direction (proton beam direction).

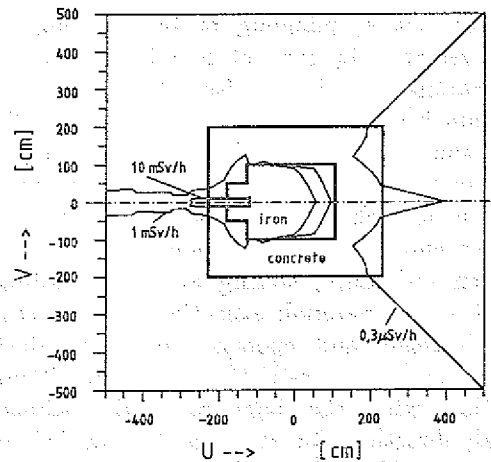


Fig. 1: Beam stop example with calculated iso-dose lines (proton beam hits at $U = -125$ cm)

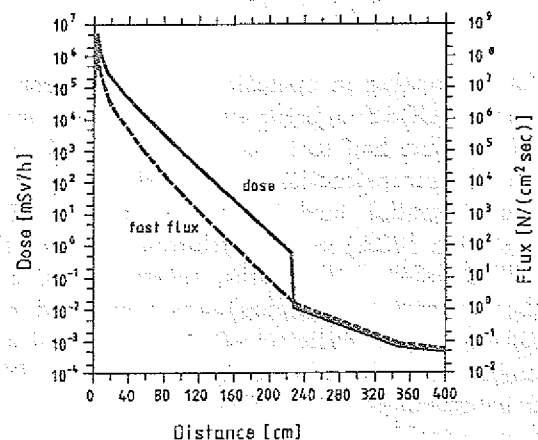


Fig. 2: Dose and fast neutron flux in zero degree direction through beam stop (zero at $U = -125$ cm in Fig. 1)

Reference:

/1/ P. Cloth et al., HERMES, Jül-2203, May 1988

IV. Technical Developments

Computer Development

IKP-VAX System

M. Karnadi, R. Nellen, K.H. Watzlawik

On the IKP Local Area VAX Cluster (IKP-LAVC) [1] and the both stand alone systems BK-VAX780 and DP-VAX8200 the operating system VMS Rel. 5.4 and the compilers FORTRAN Rel. 5.7, PASCAL Rel. 4.1 and VAX-C Rel. 3.2 have been implemented. The user systems on these computers have been updated accordingly. All in our institute existing VAX stations have been involved in the VAX Cluster. The VMS/ULTRIX Connection software (UCX) has been installed on the server and all satellites. Therefore, many kinds of UNIX functions are realized under the operating system VMS, especially TCP/IP protocol, the Internet functions Telnet and FTP as well as NFS. Within the Cluster a transition to a heterogenous Computer Network was established.

To increase the CPU and I/O performance for off-line processing and on-line applications, especially

the COSY experiment preparation, the IKP VAX-Cluster will be extended. The Server DP-VAX-11/780 will be substituted by a Server VAX4000-300. Additionally, the IKP-LAVC will be extended by four VAX-Stations-4000/60, equipped with 48MB memory, local disk and EXABYTE 8500 tape drive each. Multiuser operation environment will be distributed on individual workstations with up to eight users per workstation. The server act as boot server for all VMS stations and disk server for VMS, UNIX and OS9 systems as well as batch processing for experiment data evaluation. Therefore, the server is provided with 4.6 GB disk storage capacity and tape drives (EXABYTE 8200 and 8500, 1600/6250Bpi drives). The planning for the new IKP-LAVC is completed, the hardware components are ordered.

COSY Experiment Network, IKPnet

M. Karnadi, R. Nellen, K.H. Watzlawik

The institute-wide Ethernet based Local Area Network for COSY experiments, the IKPnet [1,2] has been enlarged according to the TCP/IP functionality. Therefore, the IKPnet was involved as a local Internet subnet (ikpnet) into the campus-wide Local Area Network, the KFAnet. All to the IKPnet connected Terminal Servers (Peripheral Port Servers, PPS) has been extended regarding the TCP/IP Telnet function. PCs connected to the IKPnet are equipped with 3COM (3C503, 3C523) Ethernet interfaces and the PC/FTP Network Software for DOS. Thereby, network-wide internet communications including Telnet, FTP, SMTP and NFS are achieved for PCs. With SmarTerm software normal graphic terminal emulation and with X-Vision software X-Windows will be realized on the PCs side. In summary, the IKPnet was extended to the heterogenous Computer Network. DECnet, LAT, LAVC as well as TCP/IP with all Internet functions and NFS are supported for various computers with the operating system such as VMS, ULTRIX, OS9, AIX, SUN-OS and MS-DOS.

At this time for electronic mail on the VAX side BITNET and VMS-Mail/Internet are used. The

protocol conversion between VMS-Mail and Internet SMTP take place on the central mail router in ZAM.

The extension of IKPnet for COSY experiment requirements is in progress [3]. Therefore, the IKPnet, involved into a one Internet subnet, will be separated by a local bridge into an on-line and off-line area, with the aim to isolate the experiment network traffic from the rest of the network. The off-line area comprises the existing network. For the on-line area a mini network, located in the experiment measurement room, will be built up by means of a DELNI and thin-wire multiport repeaters. Within the experiment area thin-wire Ethernet will be used. Dedicated thin-wire cords are assigned for local detached fields, such as experiments located at COSY ring and at the external beam, instrumentation racks located inside the experiment measurement room and user rooms. For the communication with COSY-Control, which is involved in a different subnet, a gateway will be installed in the near future. A certain contingent of Internet IP and DECnet addresses are reserved for external experimentators.

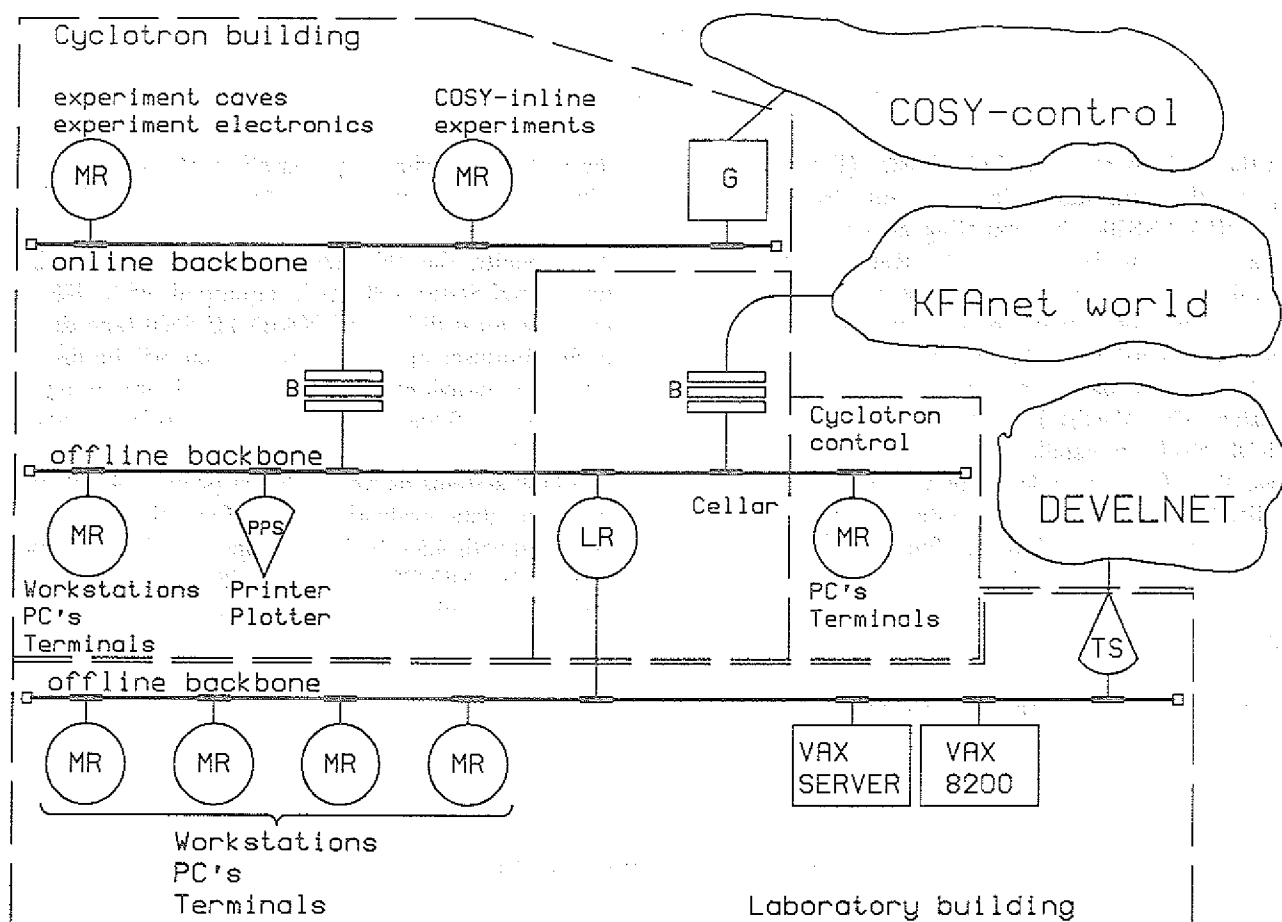


Fig. 1: COSY Experiment Network Topology, IKPnet
 MR - thin-wire Multiport Repeater
 LR - Local Repeater
 PPS - Peripheral Port Server
 B - LAN Bridge

COSY Experiment Data Processing

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In collaboration with institutes inside the KFA, neighbouring universities, other research centers and companies in Germany and other countries the activities due to the preparation of COSY experiments have been proceeded.

For the off-line evaluation of COSY experiment data the computing requirements have been prepared [4] in collaboration with scientists from our institute. These requirements consist of a summary of following data for the design of the off-line computing concept:

- estimation of average raw data produced by COSY experiments, recording methods and the volume of expected tape material,
- evaluation methods and steps,
- estimation of the required computer performance as well as the expected storage capacity for the reduced and evaluated experiment data. At this time, a solution for the off-line computing will be striven in collaboration with ZAM.

For COSY experiment instrumentation the design of 19" racks take place in collaboration with the LNPI-

Petersburg. In particular, rack mechanics, at CERN developed hydrostatic leakless cooling system as well as remote control of the cooling system and control of built-in crates are involved in this intention.

The built-up of various test systems based on CAMAC, FASTBUS and VME are in progress for detector developments. Therefore, hardware and software tools are in preparation and the implementation support is provided for controller, VME-Workstations and VME-Interfaces, operating systems VMS and OS9, driver and network software as well as several application software.

For the preparation of the experiments "Production of Lambda Hypernuclei at COSY" and "Experimental Investigation of Low-Lying States of Pionic Atoms" the data acquisition system FATIMA [5,6] is used. Therefore, the front-end hardware and the program system of FATIMA was extended. The data acquisition software was enlarged for the operation at the Local Area VAX Cluster, IKP-LAVC. This means, that various different configured FATIMA CAMAC systems, connected to several satellites across the Cluster, have access to the same program system, resided on the server. For each FATIMA CAMAC system exists its own hardware configuration file, with informations, such as crate number and number, type and position of ADCs and/or TDCs. The attachment between the CAMAC system and program take place by logical names during the login procedure.

To allow continuous, error free processing due to single or multi event buffering mode as well as mixed mode operation of various types of ADCs and TDCs, a new CAMAC module, the "Fera System Controller" (FSC)[7] has been developed in collaboration with the University of Dresden. The previous used "Memory Switching Unit", the MSU [6] is supplemented by this module.

At now, test and implementation of mixed mode operation of LeCroy and Silena ADCs and/or TDCs are still in progress. The remote handling of this operation mode will be adapted in the program system.

As programming tool for users using OS9 systems, the Unibridge cross-software was ported by ZEL on a ULTRIX DEC-Station-5000 and alternatively at the VAX4000 supported by NFS. The cross-software has been successfully tested on the Motorola, Eltec and CES VME CPUs.

The design of the digitizing systems for several experiments at COSY have been continued. Therefore, a device list for the instrumentation of COSY experiments was prepared. This device list comprises by the CANU commission "Datenaufnahme COSY" proposed instrumentation and digitizing systems

such as crates, CAMAC and FASTBUS modules, especially ADCs and TDCs, including their specifications and global functionality [8]. The functionality of FASTBUS masters, at CERN developed FVSBI and CHI in conjunction with Phillips Scientific ADCs and TDCs have been testing in cooperation with other groups.

For the Zero-Degree-Facility experiments at COSY the design of a FASTBUS based digitizing system for several detector arrangements was started.

The hardware and software concepts of the COSY experiment data acquisition have been defined in the "Summary of Recommendations" [9] by the CANU commission "Datenaufnahme COSY" in numerous meetings during the last two years. The aim, the optimal coherence regarding a modular structure and employment of standards in hardware and software for data acquisition systems at COSY of all participating groups has been achieved.

From our side, the modular data acquisition system for COSY experiments based of standardized intercrate connection, the VICbus and the Eventbuilder utilized by VME microcomputers has been defined and documented accordingly [10,11,12].

The development of the modular data acquisition system for COSY experiments was concentrated around the implementation of the VICbus as a new standard intercrate connection for instrumentation systems CAMAC and FASTBUS, the VME Eventbuilder and Workstations.

In cooperation of ZEL with CES, Geneva a CAMAC Crate Controller, type VCC2117 with standard VICbus interface has been developed and put into operation [13]. For this Controller an ESONE CAMAC library was written and implemented in a VME-Processor (FIC8232), running the operating system OS9.

For detector development, a test system was built-up, to be composed of LeCroy FERA-System at CAMAC, a VME-Workstation with OS9 and the ULTRIX DEC-Station-5000, to demonstrate the controller functionality. On the DEC-Station, at CERN developed program PAW (Physics Analysis Workstation) was ported and extended by Ethernet TCP/IP data transfer from the VME-Station to the DEC-Station.

Accordingly, the development of the VICbus interface for the FASTBUS master, the CHI was started in collaboration with the Struck company.

To connect the instrumentation systems and workstations to the VME Eventbuilder CPUs, the development of a low cost VICbus VSB interface, as well as the development of the VICbus TURBOchannel interface for Workstations are intent by different groups at ZEL.

As extension of the LeCroy FERA-System a bit-pattern CAMAC module including control and data transfer according the FERA control and ECL-databus specifications is under development at ZEL. To utilize CAMAC based digitizing systems by the FERA-system the FSC (Fera System Controller), the

mixed mode configuration of LeCroy and Silena ADCs and/or TDCs as well as the bit-pattern module will be used by COSY experiments in one or more crates per FERA system.

FERA System Controller, FSC

R. Nellen, R. Krause*, K.H. Watzlawik
(* TU Dresden)

The LeCroy FERA digitizing system becomes a standard in our institute for the CAMAC instrumentation of medium scaled detector arrangements and is used currently for detector tests and in the future at several COSY experiments.

To achieve continuous, error free processing for single or multi event buffering mode, especially to realize mixed mode configuration of various types of digitizing modules with different timing characteristics, such as LeCroy or Silena, a new CAMAC module, the "Fera System Controller" (FSC) has been developed in collaboration with the University of Dresden.

Therefore, the Fera System Controller (FSC) has been equipped with the following functionality. Hangups of FERA, occurring at now sporadically are avoided by a programmable timeout in steps of 100ns. Assuming a hangup, unless data ready signal is not present during the timeout sequence, a FERA clear will be executed.

Of course, various digitizing modules possess different conversion times and require different gates regarding the width and the start time. Therefore, a flexible timing has been established by increasing the number of gates and providing a programmable readout delay in steps of 100ns.

Two gate inputs with different start time of gates are implemented. In the time sequence, the gate inputs will be enabled by the common Master Gate and are blocked after gates come true during the conversion and readout time or the timeout sequence, including 2µs delay, after the clear signal was generated. Hangups, previously occurred by gates during the clear sequence, are prevented. The readout delay effects the data readout across the ECLbus when the last module in the configuration has finished the conversion. By this function data loss is prevented, caused by different timing characteristics.

Single and multi event double buffering modes are realized in conjunction with two LeCroy memory modules (Mod.4302) for data transfer to the micro-computer. The FSC performs the control of the memories. In multi event mode the storage of

multiple events in the memory at eventboundary of the last one will be achieved by switching to the passive memory upon overflow of the active memory (settable pointer) after the present event readout is finished. For eventmarking the first channel of the first module is used. In single event mode consecutive events are stored in different memory modules. LAM generation after the memory switching can be enabled or disabled in both modes. In case of high data rates all gates will be blocked, if both memories are filled with data.

The following count signals for scalers will be generated to perform statistical informations and timing control:

- total event count corresponds to the Master Gate,
- accepted events, the count of valid acquired events,
- rejected events, the counts of inhibited gates by the FSC,
- timeouts detected, the count of clear sequences caused by timeouts.

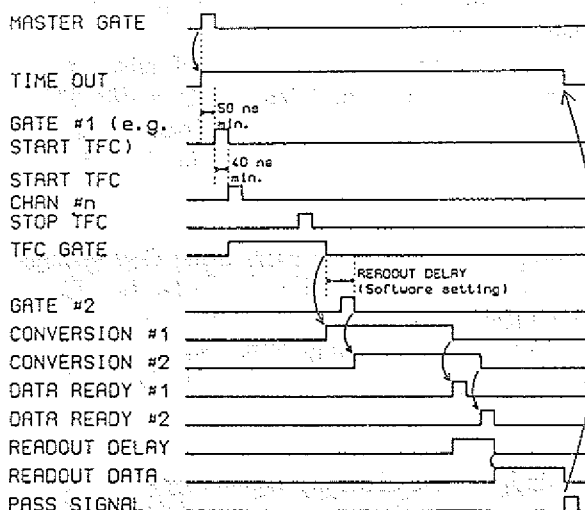


Fig. 2: Timing diagram of FERA in conjunction with the FSC

Testing EXABYTE Tape Drives

K.H. Watzlawik, J. Stock, M. Karnadi, R. Nellen

In contrast to the previous experiment operations at Cyclotron Julic with quasi continuous data stream, follows by COSY experiments periodical after each experiment cycle of 0.1 to 100s an interruption time of 1 to 100s for deceleration, injection, acceleration and cooling of the beam. Depending on the operating mode of COSY, the maximum interruption interval can be in the order of 3.6s for nocooling beam, 6.2s for electron cooling and 105s in case of stochastic cooling [14].

By reason of good experience with Exabyte 8200 tape drives for recording of experiment data and backup [15], Exabyte 8200 and the more advanced and faster drive Exabyte 8500 have been tested according to the ability for recording of data on COSY experiments, especially with regard to their start/stop behaviour.

The test configuration consists of Exabyte 8200 drive with Viking QTO controller or Exabyte 8500 with Emulex UC08 controller connected to the Qbus of the VAX-Station-3200. Alternatively Exabyte 8500 drive connected to the SCSIbus of the VAX-Station-3100/76 was tested. Standard 8mm data cartridges MP-90 are used as storage medium. On the VAX side VMS rel. 5.4 and standard TMSCP DEC drivers are used. Both drives are equipped with a speed-matching buffer (1MB Exabyte 8500 and 250KB 8200 type) organized as dual-ported circular memory.

In comparison with the Exabyte 8200, the 8500 type is sturdier and more reliable. No unconditional error or hangups have been occurred, requiring the drive or controller reset, during the test period of the type 8500. Recording specifications like response time at BOT, initialize, mount and rewind speed are similar for both types. File search is seven times and write tape marks in continuous operation five times faster (require 2s) than the 8200 type. The average transfer rate depends on the block size and amounts of a maximum of 240 KBps at the Exabyte 8200 and 500 KBps at the 8500 (fig. 3).

Regarding response time, the start behaviour of both drive types are depended on the tape position and the data stream interruption time. At the beginning of tape, the drives require a response time of 34s during which the tape is accelerated and a logical begin of tape (LBOT) pattern written. In other tape positions, measured by maximum recording speed and default drive settings, exist for the data stream interruption time (dt) two thresholds (dt1 and dt2), each

having a constant dead time (T_{0i}) for the begin of data recording within the ranges as shown in table 1. Both, the values of the thresholds and the dead time are slightly depended on the blocksize (see fig. 4,5).

EXABYTE 8200		
$T_{01} = 0s$	for $0 \leq dt < dt_1 = 15s$	
$T_{02} = 2.5s$	for $dt_1 \leq dt < dt_2 = 76s$	
$T_{03} = 7.2s$	for $dt \geq dt_2$	

EXABYTE 8500		
$T_{01} = 0s$	for $0 \leq dt < dt_1 = 16s$	
$T_{02} = 2.3s$	for $dt_1 \leq dt < dt_2 = 78s$	
$T_{03} = 9.6s$	for $dt \geq dt_2$	

Table 1: Average values of dead time for begin of data recording (T_{0i}) corresponding to the thresholds (dti) for data stream interruption time (dt).

These dead times are characteristic for the start/stop emulation on helical scan streamer drives in coherence with the dual-ported circular data buffers and are interpreted as follows [16]:

- by data stream interruption of less than dt_1 , the tape motion is continued and when the buffer is empty gap tracks are written on the tape,
- in case of data stream interruption of $dt_1 \leq dt < dt_2$ the tape path will be untensioned. At start time, data are written in the buffer and when the buffer is filled, the subsystem is blocked for T_{02} before data recording is continued,
- by data stream interruption of longer than dt_2 , the tape drum will stop rotating, and after the buffer is filled, the tape is blocked for the time T_{03} , to start tape motion and synchronisation.

The response time T_{02} can be depended on individual drive settings such as thresholds for motion, reconnection and gaps [17].

Generally, during a data stream interruption of less than dt_1 no dead time occurs and data recording takes place continuously by a blocksize greater than 4KB. That means, after the start, the data recording will be done continuously with a data rate of between $4KB/dt_1 = 0.25KBps$ and maximum speed in average. In this range of data rate, interruption times

less than $dt1$ included, the subsystems prove a satisfactory start/stop behaviour in connection with the speed-matching buffers.

In this operation mode Exabyte tape drives can be used directly for data recording of COSY experiments, excluding experiments with stochastic cooled beam, requiring periodical interruption intervals of longer than $dt1$ between the experiment cycles.

In case of longer interruption intervals than $dt1$ during the recording of COSY experiment data, the loss of experiment data during the tape dead time by start of the experiment cycle must be prevented by intermediate storage of data in the eventbuilder memory. Therefore, in case of Exabyte 8500, provided at worst-case a tape dead time of 10s and maximum speed of 500KBps, a memory of 5MB is required, organized as a queue. Alternatively, the data can be stored intermediate on the workstation disk.

Currently, all adequate tape drives such as DAT or IBM3480 [18], excluding extremely expensive and susceptible start/stop drives, are using the streamer technology and have comparable characteristics regarding their start/stop behaviour.

In summary, the tape drive Exabyte 8500 is the best approach regarding data transfer speed, tape capacity, data density and technical accomplishment for low cost bulk data storage at COSY experiments [9]. The drive is wide-spread, has a standardized SCSI interface and is used and approved on several experiments at different laboratories as well as in our institute.

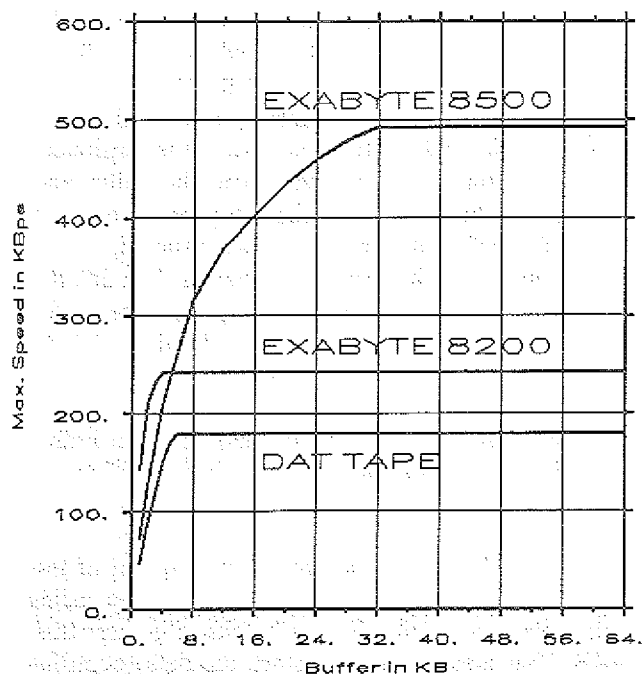


Fig. 3: Function of transfer rate in relation to the blocksize for several tape drives

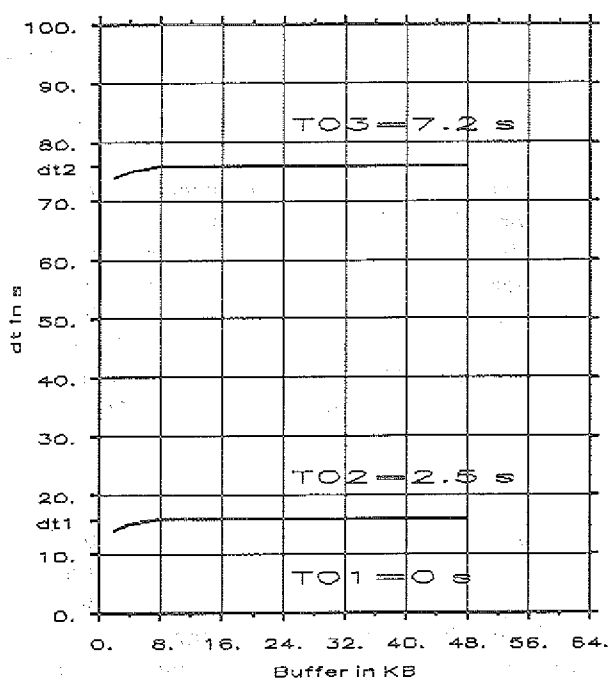


Fig. 4: Quantity of thresholds (dti) for the data stream interruption time (dt) and tape deadtime (TOi) in relation to the buffer-size for the Exabyte 8200 drive

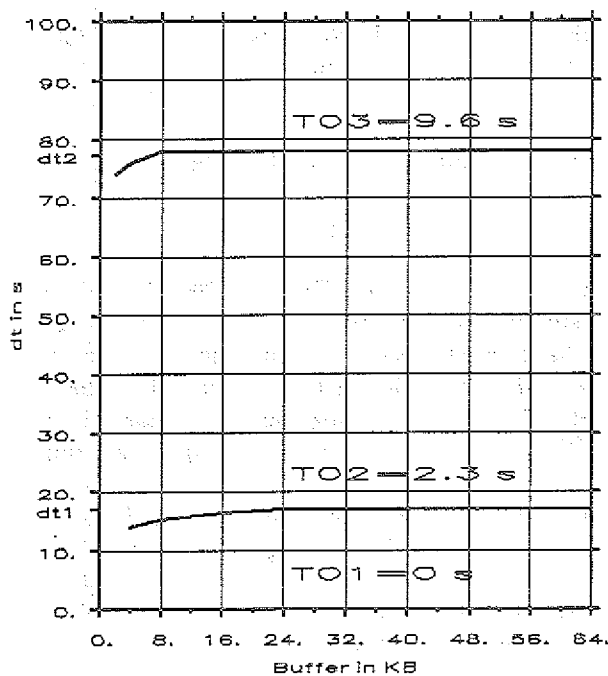


Fig. 5: Quantity of thresholds (dti) for the data stream interruption time (dt) and tape deadtime (TOi) in relation to the buffer-size for the Exabyte 8500 drive

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Feasibility Test to Use RISC Workstations for COSY Offline Data Analysis and Reconstruction

G. Sterzenbach, H.R. Langohr, B. Stugu*, M. Weber**

Offline data analysis requires processing of large datasets and a considerable amount of CPU-power. The expected requirements for the analysis of COSY experiments suggest to share work between local workstations and the mainframes at KFA. The workstation can probably do the first data reduction steps, whereas the higher level (CPU-consuming) reconstruction can be submitted to the mainframes. Therefore, the software should represent a standard running on either environment. We decided to use standard FORTRAN programs and only few library software (CERNLIB, NAG), which is available on both systems.

A first test using the reconstruction software for the CERN experiment PS-202 (Jetset) should demonstrate the feasibility of doing the first data reduction on a workstation and the interchange of datasets between the workstation and mainframes.

Therefore, the CERNLIB software was installed on the IBM RS6000 (and on a SUN Sparc station). Both the libraries (HBOOK etc.) and the utility programs (PAW) are running without problems.

Communication between the workstations (RS6000 at IKP and at ZAM and a SUN Sparc station at IKP) was established to allow window based remote login, remote printing, system-dumps etc. Communication between the workstations and the mainframe is done by TCP/IP services (see Fig. 1).

The offline data evaluation program used for the experiment PS-202 (Jetset) at CERN LEAR was implemented on the RS6000. As a first test, data were transferred via TCP-IP (FTP) from a disk of the IBM mainframe to the hard-disk of the workstation. The transfer rate was around 55 kbyte/s during daytime. When large quantities of data have to be processed, this e.g. implies a transfer time of data from a full 150 Mbyte IBM cassette of 45 minutes.

The performance of the reconstruction program for the PS 202 experiment was investigated processing experimental data. Preliminary CPU time

measurements showed the workstation to be about 50 % slower than the running of the reconstruction program on the IBM mainframe. In Table 1 the CPU time consumption for the reconstruction of 4800 events of the second level reduction is summarized.

Table 1: CPU time comparison

	CPU time for 4800 events/s
IBM mainframe	178
IBM RS6000	264

We learn from the first experience, that further tests also on other data reduction levels are necessary.

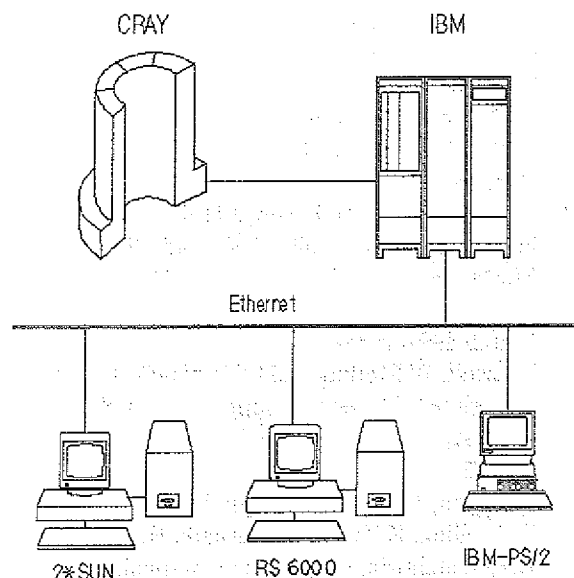


Fig. 1: Network access of workstations

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Semiconductor Detectors

G. Fiori, A. Hamacher, H. Machner, H. Metz, D. Protic, G. Riepe

The continuation of developing thin xy-position sensitive transmission type detectors from high purity germanium (HPGe) formed - like in the last year - the central subject of the activities. This detector type will be employed as the first dE-detector of special detector stacks, as one is planned to be used for COSY-experiments (see contributions of IKP1). For this purpose a position structure is projected, which consists of 200 spirals running from the circumference of each detector side to a central hole - in opposite sense on both sides.

Such a detector will be realized following established procedures of the laboratory, i.e. contacting HPGe wafers by ion implantation on both sides and subsequent creating of the structures by photolithography and plasma etching.

In view of the high count rates during operation individual amplifier channels will be connected to the single structure elements.

A first prototype of such a detector with 8 strips on either side (orthogonal to each other) had already been fabricated and tested in 1990.

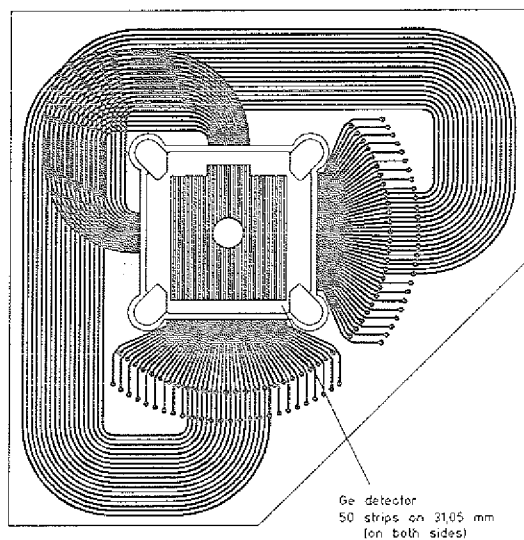


Fig. 1: The 50-strip detector with its two connecting boards.

A more advanced, second prototype with 50 strips on both sides and a central hole (7,00 mm diameter)

was now completed, fixed in a mount providing insulation on both sides and connected to 100 leads (with reduced thermal conduction) inside a special cryostat equipped with two 50-pin feedthroughs.

The thickness amounts to only 1,7 mm. The strips have a width of 535 μm . The grooves in between are 85 μm wide. (As a spare part a second diode with structure was fabricated.) Attached to the vacuum feedthroughs are boxes containing commercial pre-amplifiers in SMD-technology, which are connected through twisted-pair cables to the appropriate linear amplifiers and discriminators. These are mounted in NIM-modules, each containing 16 of them. All 100 strips were checked for proper operation in connection with their 100 amplification channels and the gains were set to the same value.

The resolution obtained in this setup amounts to 3.6 keV for 661 keV (^{137}Cs). For a preliminary data-reduction only the discriminator-signals were used, and sent through a bit-encoding and gating circuitry to a 2-parameter data-taking system. This way a 2-parameter image of the detector pixels, each resulting from two crossing strips, was feasible. In meantime the electronic circuits for the later needed 400 channels were designed, the layouts for the prints were completed, and the assembling will be started.

For an experiment previously planned for Saturne/Saclay and meant to be performed as a preparation for the experimental setup at COSY a detector-system consisting of four single HPGe-diodes had been envisaged. In the course of the plans the complete detector cryostat with many of the components and details was designed: like the mounts for the position-sensitive detectors (the first with 400 contacts, the following with 32 each), the mechanism for adjusting the detectors, the internal wiring, the joints to a D_2 -target. As a very convenient connection between the cooled detectors and the vacuum feedthroughs at room temperature proved to be a Kapton-foil equipped with printed leads.

In the scope of a collaboration with the university of Stockholm a telescope type detector-system made from HPGe was completed, to be used in experiments at the CELSIUS ring, Uppsala. The first 1.5

mm thick xy-position sensitive ΔE -detector received a pattern of 66 vertical strips on the front and 18 horizontal strips on the rear side, thus providing a sensitive area of $66 \times 36 \text{ mm}^2$. The second is a 15 mm thick E-detector of appropriate size, the fabrication of which created some problems because of difficulties with the quality of large diameter base material. (To be available as a spare for the ΔE -detector another diode, also with structure, was fabricated.) The read-out of the position-sensitive detector is made through a resistive chain bonded to the strips. Since this system will be used in an ultra high vacuum environment, the material of all components of the mounts had to be chosen such that a bake-out for 24 hours of at least 200°C can be accomplished. For the resistor chain a ceramic board with printed leads had to be prepared, which after several attempts finally could be achieved.

To be used in an experiment at PSI/Villigen, planned by a group of the university of Bonn, a detector-

system consisting of a commercial Si-detector and two transmission type HPGe-detectors (fabricated in the laboratory) with a total thickness of 28.5 mm was prepared. All detectors were mounted inside a special cryostat and tested for proper operation.

A series of commercial semiconductor detectors (Si and Ge) was tested or regenerated.

Groups of IFF were supported with some special technical procedures (plasma etching and bonding).

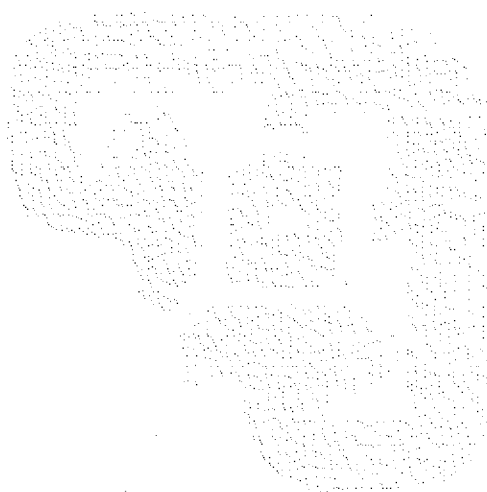
Collaboration existed with: Universities of Dortmund and Saarbrücken (position-sensitive detectors for synchrotron radiation), Tandem Accelerator Lab. Uppsala and University of Stockholm (special detectors for charged particles); furthermore inside KFA with ISI (photolithographic masks) and ZEL (nuclear electronics).

Targets

J. Pfeiffer, G. Riepe

Aside from the central subject, the development and fabrication of targets for the COSY injection system, some special targets were prepared for experiments in external laboratories. A series of Kapton-foils was coated by evaporation with gold.

For a collaboration with groups of TU München and of Heavy Ion Laboratory of University Warschau various targets were made from thorium and uranium, to be used in experiments at LEAR/CERN.



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Prof. C. Leluc-Lechanoine	University of Genf
Prof. Dr. U. Lynen	GSi Darmstadt
Prof. Dr. V. Metag	University of Gießen
Dr. W. Oelert	IKP, KFA Jülich
Prof. Dr. E. Rössle (Chairman)	University of Freiburg
Prof. Dr. B. Schoch	University of Bonn
Prof. Dr. O. Schult	IKP, KFA Jülich
Prof. Dr. J. Speth	IKP, KFA Jülich
Prof. Dr. Y. Terrien	SATURNE Saclay

VII. PERSONNEL

Institute for Nuclear Physics:

Dr. K. Kilian

Director of the Institute for Nuclear Physics
(Experiment I)
Professor of Physics at the University of Bonn

Dr. R.A. Maier

Head of the Accelerator Division

Dr. O. Schult

Director of the Institute for Nuclear Physics
(Experiment II)
Professor of Physics at the University of Köln
Managing director of the Institute in 1991

Dr. J. Speth

Director of the Institute for Nuclear Physics
(Theory)
Professor of Physics at the University of Bonn

Scientific Staff:

DP D. Anagnostopoulos (E2)

DP N. Baron (Th)

Dr. G. Baur (Th)
(Priv. Dozent at the Univ. of Basel)

Dr. U. Bechstedt (LI)

Dr. H. Beuscher (LI)

Dr. J. Bojowald (Ec)

Dr. K. Bongardt (LI)

DI N. Bongers (LI)

Prof. Dr. G. Borchert (E2)
(Apl. Prof. at the Univ. of Cologne)

DI W. Bräutigam (LI)

Dr. R. Broda (E2)
until Aug. 23, 1991

Dipl. R. Bröders (E1)
until Nov. 30, 1991

DP M. Buballa (Th)

DP M. Büscher (E2)

Dr. P. Cloth (Sd)

DP R. Collatz (E2)

DP G. Dagge (Sd)
until Aug. 31, 1991

DP M. Dahmen (E1)
since Nov. 1, 1991

DP G. Decker (E1)

Dr. J. Dietrich (LI)
since March 1, 1991

Dr. A. Djaloeis (E1)

Dr. S. Drożdż (Th)
(Inst. of Nucl. Physics
Krakov, Poland)
until Nov. 5, 1991

Dr. V. Drüke (Sd)

DP M. Elble (E2)

Dipl. R. Eßer (E2)
since Sept. 2, 1991

Dipl. J. Fabian (E2)
from Jan. 2 to Dec. 31, 1991

Dr. D. Filges (Sd)

Dr. W. Gast (E1)

DI H. Gorke (E2)

Dr. D. Gotta (E2)

Dr. F. Grümmer (Th)
(on leave to Univ. of Bochum since Febr. 2, 1989)

Dr. D. Grzonka (E1)

DI U. Hacker (LI)

Dr. J. Haidenbauer
until Jan. 31, 1991

Dr. A. Hardt (LI)

DI K. Henn (LI)

DP V. Herrmann (Th)

Dipl. M. Hoffmann (Th)
since Nov. 4, 1991

Prof. K. Holinde (Th)
(Apl. Prof. at the Univ. of Bonn)

Dipl. V. Jaeckle (E1)

Dr. P. Jahn (E1)

Dipl. G. Janßen (Th)
since Sept. 1, 1991

DP R. Jessenberger (LI)

DP M. Kalassa (Th)
since Oct. 17, 1991
(DAAD-Fellow)

Prof. Dr. K. Kilian (E1)
(Prof. at the Univ. of Bonn)

DP H. Kim (Th)

Dr. V. Klemt (Th)

Dr. H.R. Koch (E2)

DP B. Körfgen (Th)

Dipl. S. Kopp (Th)
since Dec. 1, 1991

DP D. Kopyto (E2)

Prof. S. Krewald (Th)
(Apl. Prof. at the Univ. of Bonn)

Dipl. A. Kurtenbach (E1)
since July 1, 1991

DP D. Kutchin (E1)

Dr. H. Labus (Ec)

H.R. Langohr (Sd)

Dr. H. Lawin (LI)

DP M. Liang (E2)

Prof. Dr. R.M. Lieder (E1)
(Prof. of the Univ. of Bonn)

DP Ch. Lippert (E1)

Dr. H. Machner (E1)

Dr. R.A. Maier (LI)

Dr. S. Martin (LI)

DP R. Menegazzo (E2)
(Univ. of Padua, Italy)
until July 31, 1991

DP E. Migli (Th)
(Univ. of Pisa, Italy)
until Sept. 30, 1991

Dr. P. Morsch (E1)

DP V. Mull (Th)

Prof. Dr. K. Nakayama (Th)

Dr. R.D. Neef (Sd)

Dr. S. Nishizaki (Th)
(A.v.H.-Fellow)
since Oct. 1, 1991

Dr. W. Oelert (E1)
(Priv. Dozent at the Univ. of Bonn)

Dr. H. Ohm (E2)

DP P. Oltmanns (Th)

Prof. F. Osterfeld (Th)
(Apl. Prof. at the Univ. of Bonn)

DI N. Paul (Sd)

Dr. B.C. Pearce (Th)

Dr. D. Prasuhn (LI)

DP H.J. Probst (Rp)

DP D. Protic (Dt)

DP J. Reich (LI)

Dipl. A. Reuber (Th)

Ch. Reul (Sd)

Dr. G. Riepe (Dt)

Dr. E. Roderburg (E1)

Dr. K. Roehrich (E1)

Dr. M. Rogge (E1)

Dr. P. von Rossen (LI)

A. Saad (E2)
(Scholar from Egypt)

DI M. Sauer (LI)
since Nov. 4, 1991

Dr. H. Schaal (Sd)

DP H. Schnare (E1)
until Febr. 28, 1991

DI A. Schnase (LI)

Dipl. Ch. Schütz (Th)
since Sept. 1, 1991

DI G. Schug (LI)

Prof. Dr. O. Schult (E2)
(Prof. at the Univ. of Cologne)

W. Schwab (E1)

Prof. J. Sedlacek (E2)
from Sept. 5 to Dec. 20, 1991

Dr. T. Sefzick (E1)
since Dec. 1, 1991

DI S. Selbach (E2)
until Sept. 30, 1991

DI T. Servene (E2)
since Oct. 15, 1991

Dr. H. Seyfarth (E2)

Dr. K. Sistemich (E2)
(Prof. at the Univ. of Cologne)

DP B. Spang (E1)
until April 30, 1991

Prof. Dr. J. Speth (Th)
(Prof. at the Univ. of Bonn)

Dipl. K.M. Spohr (E1)
until Aug. 31, 1991

DI R. Stassen (LI)
since July 1, 1991

Dr. H.J. Stein (E2)

DP O. Steinkamp (E1)

Dr. M. Stephan (LI)

G. Sterzenbach (Sd)

Dr. H. Stockhorst (LI)

DP K. Strähle (E1)

Prof. I. Strakovsky (E2)
from Sept. 2 to Dec. 31, 1991

DP R. Stratmann (E1)
since Jan. 1, 1991

Dr. A. Szczurek (Th)
(Inst. of Nucl. Physics
Krakov, Poland)

Dr. R. Tölle (LI)

Dr. P. Turek (E1)

Dr. W. Unkelbach (Th)
until Sept. 30, 1991

DI V. Vau (LI)

Dr. R. Wagner (LI)

Dr. T.S. Walhout (Th)
(Humboldt-Fellow)
until Sept. 30, 1991

Dipl. S. Walsh (E1)
since Aug. 1, 1991

Prof. Dr. J. Wambach (Th)
(Prof. at the Univ. of Bonn)

DP M. Waters (E1)
since Febr. 1, 1991

Dr. K.-H. Watzlawik (Da)

A. Weinert (LI)
since July 1, 1991

Dr. P. Wucherer (LI)

DP P. Wyborny (Th)

DI Ch. Zhang (E2)

DP M. Ziolkowski (E1)
(Univ. of Krakov, Poland)
until Nov. 10, 1991

Research Visitors (for one week to six months):

Dr. T. Ainsworth (Th)
from June 11 to July 26, 1991
(Texas A&M Univ., USA)

Dr. T. Bocheva (E1)
from Febr. 1 to April 30, 1991
(Bulg. Academy of Sciences, Sofia, Bulgaria)

Dr. M. Bogdanovic (E2)
from Sept. 5 to Oct. 31, 1991
(Boris Kidric Inst., Belgrad, YU)

Prof. V. Brown (Th)
from Nov. 16, 1990 to Jan. 3, 1991
(Lawrence Livermore Lab., Livermore, CA, USA)

Dr. A. Büchner (E1)
from Jan. 21 to April 20, 1991
from June 3 to Aug. 31, 1991
(ZfK Rossendorf, Dresden, Germany)

Prof. D. Cha (Th)
from Jan. 14 to Febr. 23, 1991
(Inha Univ., Inchon, Korea)

Dr. D. Dinev (LI)
until Nov. 13, 1991
(Bulgarian Academy of Sciences, Sofia)

Prof. J. Durso (Th)
from June 6 to Aug. 5, 1991
(Mount Holyoke Coll., South Hadley, MA, USA)

A.K. Georgiev (E1)
since June 10, 1991
(Bulgarian Academy of Sciences, Sofia)

L.M. Glennane (E1)
until Febr. 28, 1991
(Dublin Inst. of Technology, Irland)

Prof. H. Güven (E2)
from Aug. 11 to Oct. 4, 1991
(Univ. Istanbul, Turkey)

Prof. W-Y. P. Hwang (Th)
from Jan. 15 to March 29, 1991
(Univ. of Taipei, China)

Prof. L. Jarczyk (E1) + (E2)
from Jan. 25 to Febr. 16, 1991
from June 17 to July 15, 1991
from Nov. 17 to Dec. 20, 1991
(Univ. Krakov, Poland)

Dr. M.B. Johnson (Th)
from May 26 to June 2, 1991
(Los Alamos Nat. Lab., Los Alamos, NM, USA)

Prof. R.R. Jayanti (E1)
since Febr. 1, 1991
(Banaras Hindu Univ., India)

Prof. S. Kamedzhiev (Th)
from March 3 to May 1, 1991
from July 31 to Aug. 31, 1991
(Obninsk, GUS)

M. Kistryn (E1)
from Nov. 19 to Dec. 19, 1991
(Univ. Krakov, Poland)

Dr. St. Kistryn (E2)
from Jan. 25 to Febr. 16, 1991
from July 7 to Sept. 27, 1991
(Univ. Krakov, Poland)

Prof. V. Komarov (E2)
from Jan. 14 to Febr. 1, 1991
from April 21 to 30, 1991
(JINR Dubna, GUS)

Dr. V. Koptev (E2)
from Jan. 14 to 26, 1991,
from April 21 to 30, 1991
from June 16 to 23, 1991,
from July 28 to Aug. 11, 1991
from Sept. 29 to Oct. 6, 1991
(LNPI Gatchina, GUS)

Dr. A. Kozela (E1)
from June 12 to July 12, 1991
(Jagellonian Univ. Krakov, Poland)

R. Krause (E1)
from Sept. 16 to Dec. 13, 1991
(T.U. Dresden, Germany)

Prof. E. Lomon (Th)
from June 9 to July 9, 1991
(MIT, Cambridge, Mass., USA)

Prof. M. Macfarlane (Th)
from June 3 to July 4, 1991
(Univ. of Indiana at Bloomington, USA)

J. Majewski (E1)
from June 3 to Aug. 9, 1991
from Oct. 9 to Dec. 18, 1991
(Jagellonian Univ., Krakov, Poland)

Prof. A. Marinov (E1)
from Aug. 5 to Oct. 11, 1991
(Univ. Jerusalem, Israel)

G.V. Marti (E1)
(Dept. de Fisica, CNEA, Argentina)

A. Misiak (E1)
from Oct. 9 to Dec. 18, 1991
(Jagellonian Univ. Krakov, Poland)

Dr. T. Morek (E1)
from July 7 to Sept. 30, 1991
(Univ. Warsaw, Poland)

Dr. H. Müller (E2)
from Jan. 13 to 26, 1991
from March 20 to 28, 1991
from April 23 to May 1, 1991
from July 30 to Aug. 8, 1991
from Sept. 9 to 20, 1991
(ZfK Rossendorf, Dresden, Germany)

Prof. F. Myhrer (Th)
from May 1 to 24, 1991
(Univ. of South Carolina, Columbia, SC, USA)

Prof. H. Nann (E1)
from May 13 to 24, 1991
(Indiana Univ., USA)

Dr. B. Naumann (E1)
from Jan. 21 to Febr. 15, 1991
from Nov. 11 to Dec. 6, 1991
(ZfK Rossendorf, Dresden, Germany)

Dr. J. Okolowicz (Th)
from June 11 to July 20, 1991
from Oct. 7 to 28, 1991
(Inst. of Nuclear Physics, Krakov, Poland)

Dr. T.-St. Pawlat (LI)
until Nov. 27, 1991
(Inst. of Nucl. Physics, Krakov, Poland)

M.G. Pia (E1)
from March 18 to 23, 1991
(INFN, Genova, Italy)

Prof. H. Plendl (E1)
from May 1 to Aug. 31, 1991
(Florida State Univ., Tallahassee, Florida, USA)

Dr. P.B. Rao (Sd)
until May 15, 1991
(Indira Gandhi Centre, Kalpakkam, India)

Dr. K. Rashid (E2)
from July 11 to 17, 1991
(Pakistan Inst. of Nucl. Science, Islamabad, Pakistan)

Dr. T. Rzaca-Urban (E1)
until Febr. 28, 1991
from March 1 to Nov. 15, 1991
(Inst. of Exp. Physics, Warsaw, Poland)

Prof. M. Schepkin (E1)
from Sept. 13 to 23, 1991
(ITEP Moscow, GUS)

Dr. G. Smyrski (E1)
from Febr. 1 to 21, 1991
from June 20 to July 6, 1991
from Dec. 9 to 20, 1991
(Jagellonian Univ. Krakov, Poland)

M. Sokolowski (E1)
from Dec. 6 to 20, 1991
(Jagellonian Univ. Krakov, Poland)

Prof. A. Strzalkowski (E1) + (E2)
from Jan. 25 to Febr. 16, 1991
from June 17 to July 15, 1991
from Nov. 17 to Dec. 20, 1991
(Jagellonian Univ., Krakov, Poland)

Dr. B. Stugu (E1)
(Univ. Oslo, Norwegen))

Prof. R. Tegen (Th)
from Dec. 1 1991 to Febr. 15, 1992
(Univ. of the Witwatersrand, Johannesburg, SA)

Dr. G. Tertychny (Th)
from March 3 to May 1, 1991
from July 31 to Aug. 31, 1991
(Obninsk, GUS)

A. Turowiecki (E1)
from May 13 to 31, 1991
(Inst. of Nucl. Studies, Warsaw, Poland)

Prof. T. Udagawa (Th)
from May 20 to June 27, 1991
(Univ. of Texas at Austin, USA)

Dr. T. Venkova (E1)
from April 15 to July 9, 1991
(Bulg. Acad. of Sciences, Sofia, Bulgaria)

Dr. E. Zaplatine (LI)
since Oct. 4, 1991
(JINR Dubna, GUS)

Prof. D. Zawischa (Th)
from Febr. 25 to March 24, 1991
(Univ. of Hannover, Germany)

DP M. Ziolkowski (E1)
from Nov. 12 to Dec. 20, 1991
(Jagellonian Univ. Krakov, Poland)

Technical and Administrative Staff:

U. Banken (Ad)
H.G. Böge (LI)
W. Borgs (E2)
H. Borsch (LI)
R. Brings (LI)
G. Brittner (Ec)
P. Brittner (LI)
E. Brökel (Th)
U. Burck (Ad)
J. But (Ws)
L. Conin (LI)
B. Dahmen (LI)
W. Derissen (Cd)
H. Diesburg (Ec)
N. Dolfus (Ec)
F. Drees (LI)
R. Enge (LI)
R. Engelbrecht (Ws), until July 15, 1991
J. Engel (LI)
P. Engels (LI)
W.R. Ermer (E2)
W. Ernst (Ec)
K. Esser (Ad)
H.P. Faber (LI), since May 13, 1991
G. Fiori (Dt)
H.-W. Firmenich (Ws)
G. Gad (LI)
S. Geisler (Cd)
G. Göbbels (Ws)
A. Golik (Ws)
M. Gringel (Cd), since Jan. 25, 1991
H. Hadamek (Ws)
A. Hamacher (Dt)
J. Hanke (Sd), since Sept. 1, 1991
M. Hedrich (Da), until Dec. 31, 1991
U. Hilgers (LI)
H. Hintzen (Rp)
A. Hövelmeyer (Sd)
W. Hoffmeister (LI)
K.D. Jach (LI), since March 15, 1991
H.M. Jäger (E1)
I. Jannakos (LI)
H.J. Jansen (Ws)
R. Janssen (Ad)
A. Jeglorz (E2)

M. Karnadi (Da)
 R. Klein (E1)
 I. König (Sd)
 K. Krafft (Rp)
 M. Kremer (Ws)
 G. Krol (LI)
 K.P. Kruck (LI)
 T. Künster (Dt)
 K.G. Langenberg (LI), since July 1, 1991
 H.-J. Lindlar (LI), since Nov. 4, 1991
 W. Lorenz (Ad)
 G. Lürken (Ec)
 H. Metz (Dt)
 H. Meuser (Ws)
 A. Müller (LI)
 M. Müskes (LI)
 K. Nau (LI), until Febr. 28, 1991
 R. Nellen (Da)
 S. Peischard (LI)
 J. Pfeiffer (Dt)
 H. Pütz (LI), since April 8, 1991
 A. Retz (Cd)
 A. Richert (LI)
 U. Rindfleisch (Cd)
 DI D. Rosin (LI)
 N. Rotert (LI)
 Th. Sagefka (LI)
 A. Schaal (LI), since Nov. 4, 1991
 F. Scheiba (LI), since March 15, 1991
 P. Schmidt (LI)
 Jos. Schmitz (Ws)
 Jürg. Schmitz (LI)
 H. Schröder (Da), since Sept. 1, 1991
 F. Schultheiß (Ws)
 R. Seidemann (E2)
 H. Singer (LI)
 K. Sobotta (LI)
 W. Spieß (LI)
 E. Stegemann (Da)
 J. Stock (Da)
 C. Strauch (Ad)
 J. Strehl (Ws)
 K.P. Wieder (E2)
 W. Wilms (LI)
 K. Winkler (Ec)
 DI J.-D. Witt (LI)
 E. Wittig (Ad)
 I. Wollgarten (Cd),
 from July 4 to Sept. 30, 1991
 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

VIII. Publications

IKP-100191

Amian, W., Cloth, P., Dragovitsch, P., Drüke, V., Filges, D., Meier, M.M.

Neutron Spectra Measurements and Monte-Carlo Simulations of (p,xn) Reactions at 597 MeV and 800 MeV

Progress Report NEANDC(E)-322-U Vol. V, July 1991

20.86.0

IKP-100291

D.F. Anagnostopoulos, G.L. Borchert, D. Gotta, and K. Rashid:

K X-ray satellites of uranium bombarded with 500 MeV neon

Z. Phys. D-Atoms, Molecules and Clusters 18, 139-145 (1991)

20.50.0

IKP-100391

D. Anagnostopoulos, G.L. Borchert, D. Gotta, and K. Rashid:

Das Satellitenspektrum von K-Röntgenübergängen in Energiereichen Schwerionenstößen

12. Arbeitsbericht, Int. Winterschule Riezlern

20.50.0

IKP-100491

D. Anagnostopoulos, G.L. Borchert, D. Gotta and K. Rashid:

The K X-Ray Spectra of Ta and U following the Collisions with Energetic Ions

Proc.IOP Publishing Ltd., A. Hilger, London, in press

20.50.0

IKP-100591

Ardouin, D., Bellini, V., Bimbot, L., Bolore, M., Cerutti, C., Dabrowski, H., Erasmus, B., Gosset, J., Hisleur, J.M., Johansson, T., Julien, J., Kilian, K., Lebrun, C., Lebrun, D., Leray, S., Mougeot, A., Ngo, C., Papparlado, G.S., Saghai, B., Saintignon, P. De., Sannouillet, G., Sperduto, L.,

Search of an enhanced production of low energy pions in $\text{Cu}(p, \pi^+)X$, $\text{C}(p, \pi^+)X$ and $\text{Cu}(p, \pi^0)X$ reactions for $300 \text{ MeV} \leq T_p \leq 400 \text{ MeV}$; Z. Phys. A.

20.50.0

IKP-100691

R. Bacher, A. Badertscher, J. Egger, E. Morenzoni, L.M. Simons, J. Eades, K. Elsener, D. Gotta, P. Blüm, and K. Heitlinger:

High precision measurement of antiprotonic Hydrogen and Deuterium X-rays

Proc. First Biennial Conference on Low Energy Antiproton Physics, Stockholm, 2.-6. July 1990, ed.

P. Carlson, A. Kerek, and S. Szilagyi, World Scientific Singapore, 1991

20.50.0

IKP-100791

Balabanski, D.L., Lieder, R.M., Gast, W., Hebbinghaus, G., Krämer-Flecken, A., Kutsarova, T., Rzaca-Urban, T., Schnare, H. and Urban, W.

Study of Band Structures and Crossings in ^{179}Os

Nucl. Phys. A.

20.10.0

IKP-100891

Barnes, P.D., Birien, P., Breunlich, W.H., Decker, G., Diebold, G., Dutty, W., Eisenstein, R.A., Eyrich, W., Fischer, H., Franklin, G., Franz, J., Geyer, R., Hamann, N., Hertzog, D., Johansson, T., Kilian, K., Kirsch, W., Kraft, R., Lippert, C., Maher, C., Oelert, W., Ohlsson, S., Quinn, B., Roessle, E., Schmitt, H., Schumacher, R., Sefzick, T., Stinzing, F., Tayloe, R., Todenhagen, R., Urban, H.-J., Ziolkowski, M.

Nuclear photon scattering below and above the giant resonances.

Nucl. Phys. A530 (1991) 267

20.80.0

IKP-100991

N. Baron, G. Baur, M. Schumacher

Nuclear photon scattering below and above the giant resonances.

Nucl. Phys. A530 (1991) 267

20.80.0

IKP-101091

Baur, G.

Semiclassical and Glauber theories of electromagnetic excitation with high energy heavy ions.

Nucl. Phys. A531 (1991) 685

20.80.0

IKP-101191

Baur, G., Ferreira Filho, L.G.

Charge form-factor effects in γ - γ production of non-strongly interacting particles in relativistic heavy ion collisions.

Phys. Lett. B254 (1991) 30

20.80.0

IKP-101291

Bertulani, C.A., Baur, G., Hussein, M.S.

Two-neutron removal cross-sections of ^{11}Li -projectiles.

Nucl. Phys. A526 (1991) 751

20.80.0

IKP-101391

Beuscher, H.:

Status of ECR Source Development at Jülich,

Proc. of the 10th Intern. Workshop on ECR Ion Source, Oak Ridge (1990), p. 121

20.30.0

IKP-101491

G.L. Borchert, D. Gotta, O.W.B. Schult, L.M. Simons, K. Elsener, K. Rashid, J.J. Reidy, R.D. Deslattes, E.G. Kassler, and T. Mooney:

High resolution spectroscopy of X-rays from antiprotonic atoms using a low-energy crystal spectrometer

Proc. of the Fifth course of the Int. School of Physics of Exotic Atoms on Electromagnetic Cascade and Chemistry of Exotic Atoms, 14.-20. May 1989, Erice, Italy, Plenum Press, New York, 1990

20.50.0

IKP-101591

W. Borgs, M. Büscher, D. Gotta, H.R. Koch, W. Oelert, H. Ohm, O.W.B. Schult, H. Seyfarth, K. Sistemich, J. Ernst, F. Hinterberger, V. Koptev, S.V. Dshemuchadze, V.I. Komarov, M.G. Sapozhnikov, B.Zh. Zalyhanov, N.I. Zhuravlev, H. Müller, and P. Birien:

A 0^0 Facility in the COSY ring in Particle Production near Threshold

AIP conference proceedings no. 221, p. 299 (Nann H. and Stephenson E.J., eds.) (1991)

20.45.0

IKP-101691

Bowman, D., Daniel, H., David, P., von Egidy, T., Haninger, T., Hartmann, F.J., Hofmann, P., Jastrzebski, J., Kim, Y.S., Kurcewicz, V., Lieb, J., Lynch, W., Machner, H., Plendl, H., Riepe, G., Wright, B., Ziock, K.

Antiproton Induced Fission and Fragmentation

Proc. for Workshop on Nucleon-Antinucleon Interactions, ITEP NAN-91, ITEP, Moscow, 8.-11.7.1991

20.45.0

IKP-101791

Brinkmüller, B., Decowski, P., Marianski, B., Morsch, H.P., Paul, D. Rogge, M., Siebert, R., Spang, W., Turek, P., Zemlo, L.

Fission Decay of Giant Resonances in ^{238}U Excited in α -Scattering at Small Momentum Transfer

KFA-IPK(1)-1991-10

Phys. Rev. C.

20.45.0

IKP-101891

Brückner, J., Körfer, M., Wänke, H., Schroeder, A.N.F., Filges, D., Dragovitsch, P., Englert, P.A.J., Starr, R., Trombka, J.L., Taylor, I., Drake, D.M., Shunk, E.R.

Proton-Induced Radiation Damage in Germanium Detectors

IEEE Trans. Nucl. Sci., Vol. 38, No. 2, April 1991

20.86.0

IKP-101991

Buballa, M., Drożdż, S., Krewald, S., Speth, J. Nuclear electromagnetic response within continuum RPA theory.

Ann. Phys. 208 (1991) 346

20.80.0

IKP-102091

Buballa, M., Drożdż, S., Krewald, S., Szczurek, A. Final state interaction effects in the $(e, e'p)$ reaction.

Phys. Rev. C44 (1991) 810

20.80.0

IKP-102191

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IKP-500191
Dinev, D.:
Simulation and Correction of the Closed Orbit in the
Cooler Synchrotron COSY, Juli-2406, Nov. 1990
20.84.0

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Dinev, D.:
First Turn Simulations in the Cooler Synchrotron
COSY, Juli-2499, Juli 1991
20.84.0

IKP-500391
Filges, D., Paul, N.
Module für Detektor ZEUS gebaut — Komplexe
Detektoreinheiten für HERA bei DESY Hamburg
KFA-Intern 3/1991
20.86.0

IKP-500491
Zwoll, K., Meiling, W., Watzlawik, K.H.
COSY Experiment Datenerfassung, Systemübersicht
V1.1
KFA-ZEL-IB-500891
20.45.0

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IKP-600191

Beuscher, H.:

Status Ionenquellen und Quellenstrahlführung

Seminar Beschleunigerphysik COSY, IKP-KFA

Jülich, 26. Juni 1991

20.30.0

IKP-600291

Bongers, N.:

Inbetriebnahme AEG- und AMK-Netzgeräte am

COSY-Ring

Seminar Beschleunigerphysik COSY, IKP-KFA

Jülich, 19. Juni 1991

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IKP-600391

Jessenberger, R.:

Spinbewegung in zyklischen Beschleunigern

Seminar Beschleunigerphysik COSY, IKP - KFA

Jülich, 20. Feb. 1991

20.84.0

IKP-600491

Jessenberger, R.:

"Partial-Snake" für COSY

Seminar Beschleunigerphysik COSY, IKP-KFA

Jülich, 11. Dez. 1991

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IKP-600591

Lawin, H.:

Auslegung des Wall Current Monitors

Seminar Beschleunigerphysik COSY, IKP-KFA

Jülich, 23. Jan. 1991

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IKP-600691

Maier, R.

COSY Jülich Projekt

IFF Sommerschule, KFA Jülich, 6. März 1991

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Maier, R.:

COSY und Medizinforschung

2. Seminar der Führungskräfte der KFA Jülich

8.-9. Juni 1991

20.40.0

IKP-600891

Martin, S.:

Erste Auswertungen der Dipolfeldmessungen

Seminar Beschleunigerphysik COSY, IKP-KFA

Jülich, 9. Jan. 1991

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IKP-600991

Martin, S.:

Transmutation von radioaktivem Abfall

(Bericht von einer Arbeitstudie in Stockholm)

Seminar Beschleunigerphysik COSY, IKP-KFA

Jülich, 17. Juli 1991

20.84.0

IKP-601091

Martin, S.:

Langsame Extraktion aus COSY

Seminar Beschleunigerphysik COSY, IKP-KFA

Jülich, 30. Okt. 1991

20.84.0

IKP-601191

Prasuhn, D.:

Der COSY-Beschleunigungszyklus

Seminar Beschleunigerphysik COSY, IKP-KFA

Jülich, 2. Okt. 1991

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Stockhorst, H.:

Aspekte der Teilchenbewegung im longitudinalen

Phasenraum von COSY

Seminar Beschleunigerphysik COSY, IKP-KFA

Jülich, 5. Juni 1991

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COSY Experiment Data Processing

First Collaboration Meeting COSY Experiment 10,

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Experiment Control and on-line Data Analysis at

COSY Experiments

Collaboration Meeting COSY Experiment 11,

KFA-IPK 1.7.1991

20.45.0

XIII. INDEX OF AUTHORS

Abaev, V.	21	Derissen, W.	223
Alber, D.	92	Derrick, M.	72
de Angelis, G.	88	Dienel, S.	21
Amian, W.	75	Dietrich, J.	232,236,238,247,264
Anagnostopoulos, D.F.	66,101	Djalois, A.	46,48
Atac, A.	88,94	Dolfus, G.	221,264
Au, M.	136	Dombrowski, H.	21,41
Bach, C.	90	Dragovitsch, P.	75
Baktash, C.	67	Drake, D.M.	77
Balabanski, D.	90,92	Drochner, M.	276
Banerjee, P.	214	Drochner, W.	46,48
Baron, N.	158,160		
Baur, G.	157,158,160,213,214, 215,216,217,	Drożdż, S.	171,173,211
Bechstedt, U.	221	Drüke, V.	4,18,20,32,46,48,50, 72,75,77
Beck, E.M.	92	Dshemuchadze, S.V.	21
Beene, J.R.	67	Durso, J.	147
Berst, D.	107	Eberth, J.	90,103,107
Berst, M.	103	Egelhoff, P.	217
Bertrand, F.E.	67	Engel, J.	266
Beuscher, H.	265	Engels, P.	236
Binderberg, M.	90	Englert, P.A.J.	77
Birien, P.	21	Ermer, W.	221
Blomqvist, J.	94	Ernst, J.	21,48,121
Bochev, B.	90	Erven, W.	46,48,276
Böge, H.G.	259	Eschenauer, M.	90
Boivin, M.	69	Fabian, J.	39
Bojowald, J.	232,236,238,247	Fabian, U.	77
Bongardt, K.	232,236,238,247,252	Filges, D.	4,18,32,46,48,50,72, 74,75,77,272
Bongers, N.	250		
Borchert, G.L.	66,101	Fiori, G.	283
Borgs, W.	21,32	Forbes, S.	88
Borsch, H.	229	Frascaria, R.	69
Bräutigam, W.	256,259	Freiesleben, H.	3
Brand, H.	3	Frekers, D.	46
Brand, S.	3	Freund, S.	90
von Brentano, P.	32,90,92,103,107	Fritsche, A.	3
Brings, R.	256,259	Gast, W.	90,92,98,100,108,112
Brittner, P.	226	Gaul, G.	121
Broda, R.	85,96	Gehlen, F.	259
Bröders, R.	8,10,12,14,15	Georgiev, A.	90,98,100,108,112
Brückner, J.	77	Gjorup, N.	88
Buballa, M.	152	Glashausser, C.	162
Budzanowski, A.	41,46,48	Gleisberg, F.	276
Büchner, A.	108,112	Gotta, D.	21,32,66,101,121
Büscher, M.	21,24,26,29,81	Grawe, H.	85,90,96
Buhr, G.	122	Green, A.	162
Butzek, D.A.	4	Greenberg, W.R.	139
Cassing, W.	21,30	Grosse, E.	121
Cloth, P.	4,18,46,48,50,74,75, 77	Grzonka, D.	21,41,46,48,121
Collatz, R.	96	Gutknecht, D.	103,107
Conin, L.	222	Gysin, L.	106
Dabrowski, D.	46,48	Hacker, H.U.	226,250
Dagge, G.	77	Hadamek, H.	259
Dahmen, M.	8,10,12,14,15	Hagedorn, J.	244
Decker, G.	8,10,12,14,15	Hagemann, G.B.	88
Degenkolbe, G.	4	Haidenbauer, J.	133
		Halbert, M.L.	67
		Hamacher, A.	32,283
		Hamm, Ch.	241

Hardt, A.	221	Krug, J.	3
Hebbinghaus, G.	92	Kuhlmann, E.	3,121
Hebenstreit, J.	118	Kurtenbach, A.	8,10,12,14,15,55
Heese, J.	90	Kusakari, H.	88
Henck, R.	103,107	Kutchin, D.	90,100,103
Henn, K.	250	Labus, H.	232,264
Henneck, R.	217	Lach, M.	85,96
Herrmann, V.	180,183,186,189	LaFrance, P.	136
Herskind, B.	88	Lange, S.	3
Hessen, J.	100	Langohr, H.R.	272,282
Hinterberger, F.	21	Lawin, H.	236
Hintzen, H.J.	271	Lee, Y.	67
Hoffmeister, W.	265	Leege, K.W.	21
Hofmann, S.	85,96	Lenske, H.	119
Holinde, K.	127,129,131,133,145, 147	Lhersonneau, G.	81
Holzer, J.	46,48	Liang, M.	81,83
Horen, D.J.	67	Lieder, R.M.	88,90,92,98,100,103, 104,106,107,108,112
Hübel, H.	92	Lieb, B.J.	46,48
Ingebrechtsen, F.	88	Linz, U.	50
Jacobs, W.	69	Lippert, C.	8,10,12,14,15
Jaekle, V.	7	Loevenich, H.	276
Jäger, H.	103,107	Lomon, E.L.	136,139
Jahn, P.	262,263	Lürken, G.	232
Jannakos, I.	229	Lyachenko, V.	21
Jarczyk, L.	21,30,32,34,36,41,43, 46,48,118,119	Machner, H.	7,46,48,121,122,123, 266,276,283
Jaskola, M.	217	Maeckelburg, D.	61
Jensen, H.	88	Magiera, A.	41,118,119
Jerrestam, D.	88	Maier, K.H.	85,90,92,96,98,100, 103,107
Jessenberger, R.	252	Maier, R.	50,223,226,232,236, 238,247,252,262,264, 266
Kämmerling, H.	103,107	Majewski, J.	16
Kamerdzhev, S.	203,207,209	Manz, T.	121
Kamys, B.	21,30,41,118,119	Marti, G.	88,90,98,100,106
Karnadi, M.	275,276,279	Martin, S.	255
Khoukaz, A.	41	Matthäy, H.	3
Kilian, K.	4,5,7,8,10,12,18,41, 46,48,52,53,54,55,56, 58,60,63,64,79,121, 147	Meier, M.M.	75
Kim, H.	21,32,34,36,41,43,46, 48,268	Meiling, W.	46,48
Kistryn, St.	4,7	Menegazzo, R.	85,96
Klein, R.	85,94,96	Metz, H.	283
Kleinheinz, P.	90	Meuth, H.	236,241,244
Kluge, H.	21,32,39	Michel, P.	5
Koch, H.R.	41	Miethig, B.	5
Köhler, M.	166,168,169	Migli, E.	211
Körfggen, B.	21	Misiak, A.	16
Komarov, V.I.	46,48	Möller, K.	5
Konijn, J.	46,48,276	Morsch, H.P.	46,48,67,69
Kopp, H.	21	Mühry, H.	217
Koptev, V.	39	Müller, A.	265
Kopyto, D.	21,41	Müller, D.	94
Kozela, B.	271	Müller, H.	21,28,29
Krafft, K.	278	Mull, V.	131,133,142,145
Krause, R.	151,152,171,173	Mullins, S.	88
Krewald, S.	276	Musgrave, B.	72
Krivshitch, A.	221	Myhrer, F.	133,149
Krol, G.	229	Nakayama, K.	180,183,186,189,192
Kruck, K.		Naumann, B.	5

Naumann, L.	5	Santo, R.	21,41,121
Neef, R.D.	77,240	Santonocito, D.	88
Nellen, R.	275,276,278,279	Sarantites, D.G.	67
Nichitiu, F.	21	Sauer, M.	223,229
Nolte, M.	123	Schaal, H.	20,121
Nyberg, J.	88,94	Schalt, W.	229
Oelert, W.	7,8,10,12,14,15,16, 21,41,43,45,52,53, 54,55,56,58,60,61,63, 64,79,118,119	Schmidt, G.	5
		Schmidt, P.	266
Ohm, H.	21,32,81,83	Schmidt, R.	41
Oltmanns, P.	164,175,177	Schmitz, J.	229
Osterfeld, F.	164,166,168,169,175, 177	Schmitz, W.	92
		Schnare, H.	88,90,92,98
Ott, E.	90	Schnase, A.	241,244
Parlier, B.	77	Schneider, H.	229
Patnaik, A.	77	Schneidereit, Chr.	21
Paul, N.	4,72	Schramm, M.	85,90,96
Pawlat, T.	221	Schubart, R.	85,90,96
Pearce, B.C.	127,129	Schug, G.	226
Pfeiffer, J.	32,254,284	Schult, O.W.B.	21,24,30,32,34,36
Pfister, U.	223,226	Schwab, W.	223
Piiparinen, M.	88,94	Schwarz, U.	223,229
Plendl, H.S.	46,48	Sedlacek, M.	223
Plouin, F.	69	Sefzick, Th.	8,10,12,14,15,16,52, 53,54,58,61,64
Prasuhn, D.	223,226	Selbach, S.	66
Probst, H.J.	236,271,272	Seyfarth, H.	21,32,118,119
Protic, D.	46,48,283	Sick, I.	217
Pysz, K.	34	Siebert, R.	69
Ragnarsson, I.	83	Singer, H.	226
Rama Rao, J.	48,122	Sistemich, K.	21,24,28,30,32,81,83
Ramsoy, T.	94	Siudak, R.	119
Rebel, H.	215	Shyam, R.	214
Reedy, R.C.	77	Skwirczynska, I.	46,48
Reich, J.	262	Sletten, G.	88,94
Renftle, W.	264	Smyrski, J.	21,32,41,46,48
Reul, Ch.	20,240	Sobotta, K.	250
Richert, A.	250	Sokolowski, M.	41,43
Riepe, G.	32,46,48,254,283,284	Spang, W.	67,69
Rindfleisch, U.	4,259	Speth, J.	127,129,131,142,147, 151,154,195,197,200, 201,207,209,211
Ringe, P.	3		226
Roderburg, E.	8,10,12,14,15,41,46, 48,55	Spieß, W.	226
		Spohr, K.	90,98
Röhrich, K.	8,10,12,14,15,52,53, 55,56,58,60,61,63,64	Stassen, R.	226
		Stein, H.J.	223
Röser, A.	3	Stein, J.	108,112
Rogge, M.	262,263	Steinkamp, O.	8,10,12,56,58,60,61, 63
Romanowski, T.	72		3
Rook, J.	221	Steinke, M.	250
Rook, M.	41	Stephan, M.	18,282
Rosin, D.	221	Sterzenbach, G.	279
von Rossen, P.	46,48,118,119,266, 268	Stock, J.	226,232,238,247
		Stockhorst, H.	276
Rosso, O.	209	Stoff, H.	67
Rotert, N.	256,259	Stracener, D.W.	88,90,92,98,100,107
Rudy, Z.	118	Strähle, K.	39
Rzaca—Urban, T.	90,92,98	Strakovsky, I.	8,10,12,14,15
Sachsenhausen, H.	4	Stratmann, R.	
Sagefka, Th.	221		
Saghai, B.	69		

Strzalkowski, A.	21,30,32,34,36,41,43, 46,48,118,119	Ziolkowski, M.	8,10,12,14,15,16,41, 45,52,53
Stugu, B.	41,56,58,60,63,64, 282	Zolnierczuk, P.A.	36
Styrczen, B.	32,118	Zumloh, A.	223,226
Sugawara, M.	88	Zupranski, P.	69
Szczurek, A.	151,171,173,195,197, 200,201	Zwoll, K.	46,48,276
Tegen, R.	154		
Tegner, P.E.	69		
Teichert, W.	103,107		
Tertychny, G.	203,209		
Theuerkauf, J.	90		
Thoennessen, M.	67		
Thomas, H.G.	103,107		
Tjom, P.O.	88		
Tölle, R.	266,268		
Trautmann D.	213,217		
Travkin, V.	21		
Turek, P.	4,7,41,46,48,79,121		
Udagawa, T.	164,166,168,169,175, 177		
Unkelbach, W.	162,203,207,209		
Urban, W.	92		
Utzelmann, S.	90		
Varner, R.L.	67		
Vashegyi, T.	250		
Vau, V.	222		
Veiders, E.	229		
Vetterli, D.	217		
Virtanen, A.	88,94		
Wadsworth, R.	88		
Wänke, H.	77		
Walhout, T.S.	135		
Walsh, S.	8,10,12,14,15		
Wambach, J.	135,142,162,211		
Warde, E.	69		
Waters, M.	79,121		
Watzlawik, K.H.	21,32,275,276,278, 279		
Weber, M.	282		
Weinert, A.	250		
Werdin, R.	3		
v.d. Werth, A.	103,107		
Wiedner, C.A.	255		
Wimmer, J.	223		
Witt, J.—D.	223		
Wolfertz, B.	272		
Wolter, H.H.	119		
Wolters, H.	90,92		
Wucherer, P.	259		
Wuestner, P.	276		
Yonnet, J.	69		
Zalyhanov, V.	21		
Zawischa, D.	207		
Zeijma, J.	43		
Zell, K.O.	90,92,98		
Zhang, C.T.	85,96		
Zhuravlev, N.I.	21		

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