P_{cc}^{N} states in a unitarized coupled-channel approach

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Starting from an effective Lagrangian with heavy quark spin symmetry embedded, the coupled-channel dynamics of the doubly charmed systems $D^{(*)}\Sigma_c^{(*)}$ is investigated. The potential underlying our investigation includes t-channel pseudoscalar and vector meson exchanges. A series of S-wave bound states with isospin I=1/2 is found by applying the first iterated solution of the N/D method one state with binding energy 23 MeV in the $5/2^ D^*\Sigma_c^*$ channel, three states with binding energy 26, 30 and 7 MeV (relative to the thresholds from low to high, respectively) in the $3/2^ D\Sigma_c^*$ - $D^*\Sigma_c^ D^*\Sigma_c^*$ system and three states with binding energy 32, 8 and 16 MeV in the $1/2^ D\Sigma_c$ - $D^*\Sigma_c$ - $D^*\Sigma_c^*$ system. Those P_{cc}^N states serve as the open-charm partners of the hidden charm pentaquarks P_{ψ}^N observed by the LHCb Collaboration.

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I. INTRODUCTION

Searches of exotic hadrons, whose valence quark composition is beyond the conventional picture where mesons and baryons are composed of a pair of quark-antiquark $(q\bar{q})$ and three quarks (qqq), respectively, have become an important project for most of the collider facilities especially after the experimental observations of tetraquark and pentaquark candidates [1-7]. In 2015, the LHCb Collaboration announced the first evidence of two hidden-charm pentaquark-like states $P_\psi^N(4380)$ and $P_\psi^N(4450)$ in the $J/\psi p$ invariant mass spectrum measured from the decay process $\Lambda_b^0 \to J/\psi K^- p$ [6]. In 2019, the mass spectrum of P_ψ^N pentaquarks was updated to three states, that is, $P_\psi^N(4312)$, $P_\psi^N(4440)$ and $P_{ab}^{N}(4457)$ by the LHCb Collaboration [8]. In 2020, a new hidden-charm pentaquark state with strangeness, namely $P_{\psi s}^{\Lambda}(4459)$, was observed in the $J/\psi\Lambda$ invariant mass distribution from the $\Xi_b^- \to J/\psi \Lambda K^-$ decay [9]. And very recently, another hidden-charm pentaquark with strangeness, $P_{\psi s}^{\Lambda}(4338)$, was announced by the LHCb Collaboration [10]. It is observed in the invariant mass spectrum of $J/\psi\Lambda$ in the decay $B^- \to J/\psi\Lambda\bar{p}$. In the last decade, the LHCb Collaboration has found many surprises in exotic hadron spectroscopy. One can expect that the richness of the exotic spectrum will continue to increase in the foreseeable future.

The theoretical investigations on the exotic spectroscopy date back to the birth of the quark model in 1964. The existence of pentaguark states was first

pointed out by Gell-Mann in his famous paper on the quark model [11]. From the point of view of modern physics, neither the multiquark states that made up of more than three valence quarks such as tentraquarks $(qq\bar{q}\bar{q})$ and pentaguarks $(qqqq\bar{q})$, nor the hybrid states that have both valence quarks and gluons or the glueballs that are composed of pure valence gluons are forbidden by Quantum Chromodynamics (QCD), which is the fundamental theory of the strong interactions. Before the first experimental evidence of P_{ψ}^{N} , these pentaguark states have been predicted by the theoretical work based on the phenomenological coupled-channel approach in 2010 [12]. In that work, one $\bar{D}\Sigma_c$ and one $\bar{D}^*\Sigma_c$ bound state were found around 4 GeV, which can be related to the observed $P_{\psi}^{N}(4312)$ and $P_{\psi}^{N}(4440)$ or $P_{\psi}^{N}(4457)$ states, respectively. In particular, the newly reported $P_{\psi s}^{\Lambda}(4459)$ and $P_{\psi s}^{\Lambda}(4338)$ states are also compatible with the predicted $\bar{D}^*\Xi_c$ and $\bar{D}\Xi_c$ bound states, respectively. Such impressive consistence between the experimental observations and the theoretical predictions on the pentaguark spectrum around 4 GeV is a strong indication for the molecular nature of those pentaguarklike states. The hadronic molecule picture has become are much discussed approach to explain the nature of the exotic candidates, as seen by the many theoretical studies of exotic hadron spectroscopy during the last decades, see the recent reviews in Refs. [13–17].

In Ref. [18], the authers claim that the near-threshold structures exist generally in two heavy hadron systems as long as the interaction between them is attractive. Hundreds of hadronic molecules are proposed in the heavy-heavy [19–21] and heavy-antiheavy sectors [22]. And very recently, several five-flavored bound states are predicted in the $B^{(*)}\Xi_c^{(\prime)}$ system [23]. Among these exotic baryons, the doubly-charmed pentaquarks are straightforward extensions of the P_ψ^N states, whose quark content can be

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written as $ccqq\bar{q}$ (q = u, d). In this work, we explore the mass spectrum of the doubly-charmed pentaguarklike states around 4 GeV. Some earlier investigations on the this system are given in Refs. [19-21]. Ref. [19] constructed the contact, one-pion-exchange, and two-pionexchange potentials for the coupled-channel $D^{(*)}\Sigma_c^{(*)}$ system within the framework of a chiral effective field theory and found the S-wave bound states by solving the nonrelativistic Schrödinger equation. Ref. [21] updated the configuration in Ref. [19] by introducing the S-D mixing effect within the one-boson-exchange (OBE) model assisted with heavy quark spin symmetry (HQSS) and Ref. [20] only studied the single channel case. All those three works give a similar mass spectrum for the doublycharmed pentaguark candidates, that is, one $1/2^ D\Sigma_c$, one $3/2^ D\Sigma_c^*$, two $D^*\Sigma_c$ with spin-parity $1/2^-$ and $3/2^-$, and three $D^*\Sigma_c^*$ bound states with spin-parity $1/2^-$, $3/2^-$ and $5/2^-$. All those doubly-charmed bound states have isospin I = 1/2 and the binding energies var from several MeV to tens of MeV.

In the present work, both pseudoscalar (π, η) and vector (ρ, ω) meson exchanges are considered by means of an effective Lagrangians that is constrained by the HQSS together with the chiral symmetry for the pseudoscalar meson part and the hidden local symmetry for the vector meson part. The bound states and resonances are found as poles of the coupled-channel scattering amplitudes given by a unitarized Bethe-Salpeter equation (BSE) in the on-shell factorization approach. Further, the first iterated solution of the N/D method is employed to avoid the unphysical left-hand-cut problem in the on-shell factorized BSE [24, 25].

This work is organized as follows. In Sec. II, we present the theoretical framework of our calculation. The numerical results for the $D\Sigma_c$, $D\Sigma_c^*$, $D^*\Sigma_c$ and $D^*\Sigma_c^*$ coupled-channel dynamics and relevant discussions are presented in Sec. III. Finally, a brief summary is given in Sec. IV. Some technicalities are relegated to the Appendices.

II. FORMALISM

Before going to the details of our theoretical calculations, we briefly count the number of channels in the coupled-channel $D^{(*)}\Sigma_c^{(*)}$ system included in our work. It should be noticed that we only consider S-wave scattering with isospin I=1/2 throughout the present work. The quantum numbers of the various $D^{(*)}\Sigma_c^{(*)}$ channels are listed in Tab. I. It shows that we have three channels $(D\Sigma_c, D^*\Sigma_c$ and $D^*\Sigma_c^*)$ for spin-parity $J^P=1/2^-$, three channels $(D\Sigma_c^*, D^*\Sigma_c$ and $D^*\Sigma_c^*)$ for $J^P=3/2^-$, and one channel $(D^*\Sigma_c^*)$ for $J^P=5/2^-$.

TABLE I: S-wave channels in the $D\Sigma_c$ - $D\Sigma_c^*$ - $D^*\Sigma_c$ - $D^*\Sigma_c$ coupled-channel system. The quantum numbers are given in the spectroscopic form $^{2S+1}L_J$.

Channel	$J^P = 1/2^-$	3/2-	5/2-
$D\Sigma_c$	$^{2}S_{1/2}$	-	-
$D\Sigma_c^*$	-	$^{4}S_{3/2}$	-
$D^*\Sigma_c$	$^{2}S_{1/2}$	$^{4}S_{3/2}$ $^{4}S_{3/2}$	-
$D^*\Sigma_c^*$	${}^{2}S_{1/2}$ ${}^{2}S_{1/2}$	$^{4}S_{3/2}$	$^{6}S_{5/2}$

A. Effective Lagrangians and the on-shell factorization approach of the Bethe-Salpeter equation

Chiral perturbation theory (ChPT) developed in the 1980s [26–28] has achieved great success in describing low-energy experiments of the strong interaction, especially the $\pi\pi$ and πN scattering [29]. A variety of ChPT variants were proposed to solve various specific stronginteraction systems. Among them, the heavy baryon and heavy meson chiral perturbation theory are designed to describe the interactions between two hadrons containing one or more heavy quarks [30–33], see Ref. [17] for a recent review. Similar to the ChPT language, the interactions between two heavy hadrons are constructed from light pseudoscalar exchange, the Goldstone bosons from the spontaneous breaking of chiral symmetry. The interactions between heavy hadrons and light vector mesons are built by using the hidden local symmetry approach [34–36]. All the relevant effective Lagrangians used here are given as [22, 37]

$$\mathcal{L}_{1} = ig\langle H_{b}^{(Q)} \mathcal{A}_{ba}^{\mu} \gamma_{\mu} \gamma_{5} \bar{H}_{a}^{(Q)} \rangle,
\mathcal{L}_{2} = -i\beta \langle H_{b}^{(Q)} v^{\mu} \rho_{\mu,ba} \bar{H}_{a}^{(Q)} \rangle
+ i\lambda \langle H_{b}^{(Q)} \sigma^{\mu\nu} (\partial_{\mu} \rho_{\nu} - \partial_{\nu} \rho_{\mu})_{ba} \bar{H}_{a}^{(Q)} \rangle,
\mathcal{L}_{S} = \frac{3}{2} g_{1} (iv_{\kappa}) \epsilon^{\mu\nu\lambda\kappa} \bar{S}_{\mu,ab}^{(Q)} \mathcal{A}_{\nu,bc} S_{\lambda,ca}^{(Q)} - i\beta_{S} \bar{S}_{\mu,ab}^{(Q)} v^{\alpha} \rho_{\alpha,bc}
\times S_{ca}^{(Q)\mu} + \lambda_{S} \bar{S}_{\mu,ab}^{(Q)} (\partial^{\mu} \rho^{\nu} - \partial^{\nu} \rho^{\mu})_{bc} S_{\nu,ca}^{(Q)}, \tag{1}$$

where the Lorentz indices are given by the greek letters μ, ν, \cdots , the SU(3) flavor indices are denoted by latin symbols a, b, \cdots and $\langle \cdots \rangle$ is the Dirac trace. The summation over repeated indices is implicit. The explicit formulae for all vertices in our t-channel potentials are then obtained by expanding the above Lagrangians. The expansion of all the underlying Lagrangian and the definitions of the involved field operators are displayed in App. A.

The scattering amplitude T is unitarized through the Bethe-Salpeter equation, namely

$$T = V + VGT, (2)$$

where V is the scattering kernel that is expressed in terms of the t-channel one-boson-exchange transitions between

two channels and G is the two-meson loop function. In the coupled-channel case, T is a $n \times n$ matrix (n denotes the number of the coupled channels) and G becomes a n-dimension diagonal matrix with all the elastic loop functions g_i as its elements (i is the channel index). Working with the on-shell factorization approach, the integral equation (2) is reduced to an algebraic equation and one can solve the amplitude T with

$$T = (1 - VG)^{-1}V. (3)$$

Unitarity and analyticity of T are guaranteed in the on-shell factorization prescription in most cases [38–40]. And it has been applied commonly into various hadron systems to study the low-energy strong interaction dynamics in them, see e.g. [14, 17, 41–48] (and references therein). It should also be mentioned that in some cases where the unphysical left-hand cuts (usually come from the partial-wave projection of V) are not far away from the energy regions of interest, the unitarity and analyticity of T become problematic due to the existence of such unphysical cuts, see e.g. Refs. [25, 49] for more details.

For the two-meson Green's function g_i , we adopt dimensional regularization to arrive at

$$g_{i} = i \int \frac{d^{4}q}{(2\pi)^{4}} \frac{1}{(p-q)^{2} - M_{B}^{i2} + i\varepsilon} \frac{1}{q^{2} - M_{P}^{i2} + i\varepsilon}$$

$$= \frac{1}{16\pi^{2}} \left\{ a(\mu) + \ln \frac{M_{B}^{i2}}{\mu^{2}} + \frac{M_{P}^{i2} - M_{B}^{i2} + s}{2s} \ln \frac{M_{P}^{i2}}{M_{B}^{i2}} + \frac{\bar{q}_{i}}{\sqrt{s}} \left[\ln(s - M_{B}^{i2} + M_{P}^{i2} + 2\bar{q}_{i}\sqrt{s}) + \ln(s + M_{B}^{i2} - M_{P}^{i2} + 2\bar{q}_{i}\sqrt{s}) - \ln(-s - M_{B}^{i2} + M_{P}^{i2} + 2\bar{q}_{i}\sqrt{s}) - \ln(-s + M_{B}^{i2} - M_{P}^{i2} + 2\bar{q}_{i}\sqrt{s}) \right] \right\}, \tag{4}$$

where $s = p^2$ and

$$\bar{q}_i = \sqrt{\frac{(s - (M_B^i + M_P^i)^2)(s - (M_B^i - M_P^i)^2)}{4s}}.$$

 M_B^i, M_P^i denote the baryon and meson masses in the channel i, and $a(\mu)$ is the scale-dependent subtract constant. Note that $\operatorname{Im}(\bar{q}) \geq 0$ indicates that Eq. (4) gives the loop function in the physical sheet, denoted as g_i^I . The loop function in the unphysical sheet, denoted as g_i^{II} , is then expressed as

$$g_i^{II}(s) = g_i^I + 2i\rho_i(s), \tag{5}$$

with $\rho_i(s) = \bar{q}_i/(8\pi\sqrt{s})$. Another strategy commonly adopted to calculate the loop function G is to introduce a phenomenological form factor, such as the Gaussian regulator, in the integral, that is,

$$g_i = \int \frac{q^2 dq}{4\pi^2} \frac{\omega_B^i + \omega_P^i}{\omega_B^i \omega_P^i} \frac{e^{-2q^2/\Lambda^2}}{s - (\omega_B^i + \omega_P^i)^2 + i\varepsilon}, \quad (6)$$

where ω_B^i and ω_P^i are the on-shell energies for the baryon and meson in *i*th channel, respectively. In this work, we take the convention of Eq. (4), where $a(\mu)$ is estimated by matching g_i of Eq. (4) to the one of Eq. (6) at *i*th threshold with the empirical values of cutoff Λ , i.e., around 1 GeV.

B. First iterated solution of the N/D method

As mentioned above, the existence of the left-hand cuts (LHC) could invalidate the on-shell factorization formula of Eq. (3) in some cases. Unfortunately, this indeed happens in the $D^{(*)}\Sigma_c^{(*)}$ systems of interest. In this subsection, we briefly introduce the N/D method that can be used to treat those unphysical LHC properly [24, 25, 40, 50, 51].

In the N/D method, the unitarized scattering amplitude T is constructed through the dispersion relations and it has the general form

$$T(s) = D(s)^{-1}N(s),$$
 (7)

where the numerator N(s) and denominator D(s) contain the analytic information of the left- and right-hand cuts, respectively. The general expressions of N(s) and D(s) for the S-wave are given by

$$D(s) = \sum_{m=0}^{n-1} a_m s^m + \frac{(s-s_0)^n}{\pi} \int_{s_{\text{thr}}}^{\infty} ds' \frac{\rho(s')N(s')}{(s'-s)(s'-s_0)^n},$$

$$N(s) = \sum_{m=0}^{n-1} b_m s^m + \frac{(s-s_0)^n}{\pi} \int_{-\infty}^{s_{\text{left}}} ds' \frac{\text{Im}T(s')D(s')}{(s'-s)(s'-s_0)^n},$$
(8)

where the polynomials $\sum_m a_m s^m$ and $\sum_m b_m s^m$ denote the subtraction terms with a_m and b_m the corresponding subtraction constants. s_0 is the subtraction point and n is the number of subtractions that is required to ensure the convergence of the dispersion integrals. Note that the so-called Castiliejo-Dalitz-Dyson (CDD) poles [52] are dropped here. The difficulty caused by the unphysical LHC will be overcome if one solves exactly the N/D integral equations Eq. (8). An approximative strategy called the first-iterated solution of the N/D method that was proposed in Refs. [24, 25] is utilized in our work. It states that we approximate the numerator N(s) as the tree-level potential V(s) and then the denominator D(s) can be expressed as

$$D_{ij}(s) = \gamma_{0ij} + \gamma_{1ij}(s - s_{\text{thr}}^{j}) + \frac{1}{2}\gamma_{2ij}(s - s_{\text{thr}}^{j})^{2} + \frac{(s - s_{\text{thr}}^{j})s^{2}}{\pi} \int_{s_{\text{thr}}^{j}}^{\infty} ds' \frac{V_{ij}(s')\rho_{j}(s')}{(s' - s_{\text{thr}}^{j})(s' - s)s'^{2}}, \quad (9)$$

where i and j are the channel indices. The subtraction point s_0 is set to be s_{thr}^j , the jth threshold, and n=3.

Three subtraction constants γ_{0ij} , γ_{1ij} and γ_{2ij} are determined by matching $D_{ij}(s)$ of Eq. (9) and $\delta_{ij} - V_{ij}(s)G_j$ around the threshold s^j_{thr} for each i and j, specifically. We fit $D_{ij}(s)$ of Eq. (9) to $\delta_{ij} - V_{ij}(s)G_j$ in the small energy region from the threshold s^j_{thr} to 100 MeV above it. It is worth mentioning that G_j in the matching procedure is the loop function in the physical sheet. Consequently, the scattering amplitude T of the physical sheet is calculated with the D function constructed in Eq. (9). Further, T in the unphysical sheets are defined as [25],

$$T^{II}(s) = \frac{1}{[T^I(s)]^{-1} - 2i\rho},\tag{10}$$

where ρ denotes the diagonal matrix diag $\{N_i\rho_i(s)\}$ with $N_i=0$ and 1 representing the physical and unphysical sheets for the *i*th channel. The uncertainty stemming from the ambiguity of the choice of the matching energy region will be discussed when we present our numerical results.

III. RESULTS AND DISCUSSIONS

First, we focus on the case of $J^P = 5/2^-$ and explain the necessity of treating the LHC in the partial-wave projected potentials by using the N/D method. The S-wave potential for the $5/2^-$ - $D^*\Sigma_c^*$ system is given in terms of the t-channel ρ -, ω -, π - and η -exchange diagrams. The numerical potential $V_{D^*\Sigma_c^*\to D^*\Sigma_c^*}$ is presented in Fig. 1. One can clearly see that the S-wave projection for each boson exchange produces one LHC located below the $D^*\Sigma_c^*$ threshold. Among them, the LHC corresponding to the π exchange is quite close to the threshold and that from the vector meson exchange is stronger.

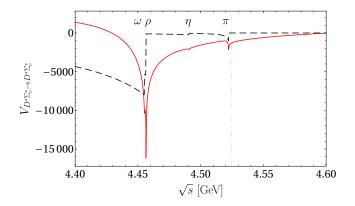


FIG. 1: S-wave potential for $D^*\Sigma_c^* \to D^*\Sigma_c^*$ with (I,J)=(1/2,5/2) (real part: solid curve, imaginary part: dashed curve). The corresponding threshold is represented by the orange dash-dotted line. The left-hand cuts from left to right are generated by the ω , ρ , η and π exchange, respectively.

The (1-VG) term in Eq. (3) and D in Eq. (9) of the physical Riemann sheet RS(+) with the cutoff $\Lambda=0.7\,\mathrm{GeV}$ for the S-wave $5/2^--D^*\Sigma_c^*$ system are shown in the left and right panel of Fig. 2, respectively. The pole position of scattering amplitude T in the single channel case is represented by the zeros of (1-VG) in the on-shell factorization BSE approach or equivalently the zeros of the D function in the N/D method.

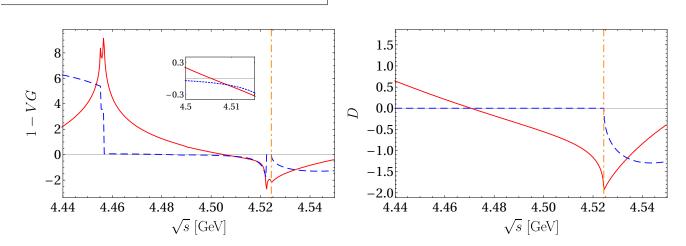


FIG. 2: (1 - VG) (left panel) and D (right panel) in the physical sheet for the S-wave $D^*\Sigma_c^*$ system with (I, J) = (1/2, 5/2) (real part: solid curve, imaginary part: dashed curve). The cutoff is $\Lambda = 0.7 \,\text{GeV}$. The corresponding threshold is represented by the orange dash-dotted line.

It can be seen from Fig. 2 that there no pole appears in the function 1-VG and one pole located below the

threshold at the real axis can be found in the D function which is related to a $D^*\Sigma_c^*$ bound state with binding

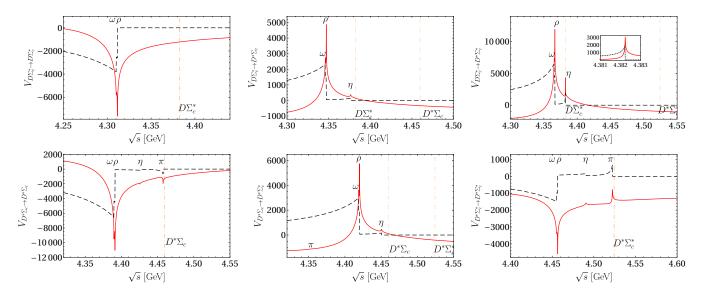


FIG. 3: V matrix for the S-wave $D\Sigma_c^* - D^*\Sigma_c^* - D^*\Sigma_c^*$ coupled-channel system with (I, J) = (1/2, 3/2) (real part: solid curve, imaginary part: dashed curve). The corresponding threshold is represented by the orange dash-dotted line.

energy $E_B=53\,\mathrm{MeV}$. However, all the previous works imply the existence of an S-wave $D^*\Sigma_c^*$ bound state with $J^P=5/2^-$ [19, 20]. To be more specific, Ref. [19] claimed an S-wave $D^*\Sigma_c^*$ state with binding energy being around 20 MeV in the non-relativistic one- and two-pion exchange potentials. Ref. [20] found a bound state from the non-relativistic ρ - and ω -exchange interactions with the binding energy varying from 2 to 40 MeV. The S-wave $D^*\Sigma_c^*$ state with $E_B=53\,\mathrm{MeV}$ obtained from the N/D method is consistent with the previous non-relativistic investigations.

Next, we turn to the $D\Sigma_c^*-D^*\Sigma_c$ - $D^*\Sigma_c^*$ coupled-channel system with (I,J)=(1/2,3/2). The V matrix is presented in Fig. 3. Similar to the $J^P=5/2^-$ single channel case, all the meson-exchange potentials contain the left-hand cuts after implementing the partial-wave projection except the one-pion-exchange parts in $V_{D\Sigma_c^*\to D^*\Sigma_c}$ and $V_{D\Sigma_c^*\to D^*\Sigma_c^*}$. In particular, the left-hand cuts from the η -exchange interaction in $V_{D\Sigma_c^*\to D^*\Sigma_c^*}$ and the π -exchanges in the elastic potentials $V_{D^*\Sigma_c\to D^*\Sigma_c^*}$

and $V_{D^*\Sigma_a^*\to D^*\Sigma_a^*}$ are quite close to and below the lower thresholds. For these cases, the dispersion relations used in N/D method are still valid since no LHC goes above the RHC. However, the effects left in the unitarized scattering amplitude T by those near-threshold LHC are difficult to remove completely. To do so as much as possible, we shift the energy region, where the subtraction constants in the D functions of Eq. (9) are fitted to functions (1 - VG) in Eq. (3) a bit above the corresponding thresholds. In practice, the D_{ij} functions are fitted to $\delta_{ij} - V_{ij}G_j$ in the energy range starting from 10 MeV above the jth threshold up to $100 \,\mathrm{MeV}$ above it. Then the pole positions are found to be zeros of determinants of D for the physical sheet and $([T^I]^{-1} - 2i\rho)$ for the unphysical sheets. The pole trajectory in the unphysical sheet RS(--+) for the $3/2^ D\Sigma_c^*$ - $D^*\Sigma_c$ - $D^*\Sigma_c^*$ coupledchannel system is shown in Fig. 4. When Λ is less than $0.5\,\mathrm{GeV},\,\mathrm{no}\;D^*\Sigma_c^*$ bound state with $J^P=3/2^-$ is found. The cutoff range from 0.6 to $0.8\,\mathrm{GeV}$ is then chosen to give the full mass spectrum in this work.

The calculations for the $D\Sigma_c$ - $D^*\Sigma_c$ - $D^*\Sigma_c^*$ coupled channels with (I,J)=(1/2,1/2) can be done along similar lines. The V matrix for (I,J)=(1/2,1/2) is given in Fig. 5.

Finally, all the obtained pole positions are collected in the Tab. II. Similar to other previous works [19–21], one $D\Sigma_c$ bound state with $J^P=1/2^-$, one $D\Sigma_c^*$ bound state with $J^P=3/2^-$ and one $D^*\Sigma_c^*$ bound state with $J^P=5/2^-$ are found. These three states are located at the physical real axis and below the threshold of the lowest channel in the corresponding systems, thus they

are bound states. Two resonances located at the complex plane of the unphysical sheets RS(-++) and RS(--+) are found in both $1/2^-$ and $3/2^-$ coupled channels, which are related to the $D^*\Sigma_c$ and $D^*\Sigma^*$ bound states with corresponding spin-parities obtained in the single-channel investigations by Ref. [20]. The sensitivities of various pole positions to the cutoff are different, the most sensitive case is the $5/2^ D^*\Sigma_c^*$ pole, whose binding energy varies from 23 MeV to 110 MeV as the cutoff changed from 0.6 GeV to 0.8 GeV, and the least sensitive one is the $1/2^ D^*\Sigma_c$ pole, whose binding energy varies from

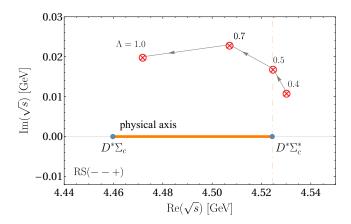


FIG. 4: The pole trajectory on the sheet $\mathrm{RS}(--+)$ of the S-wave $D\Sigma_c^*$ - $D^*\Sigma_c$ - $D^*\Sigma_c^*$ coupled-channel system with (I,J)=(1/2,3/2). The cutoff Λ is varied from 0.4 to 1.0 GeV. The highest threshold is represented by the orange dash-dotted line. The pole positions are denoted as the red crosses. The sheet $\mathrm{RS}(--+)$ intersects with the physical region at the thick orange line that connects the $D^*\Sigma_c$ and $D^*\Sigma_c^*$ thresholds.

 $8\,\mathrm{MeV}$ to $18\,\mathrm{MeV}.$ Noted that all the widths of resonances depend on the Λ only mildly and thus the averaged values are listed.

We plot the full mass spectrum together with the uncertainty caused by the cutoff Λ in Fig. 6. The theoretical error from the ambiguity of the matching energy region that needed for the determination of the subtraction constants in Eq. (9) is also estimated. The variation of this matching range from (10, 100) MeV above threshold to (30, 100) MeV induces an additional uncertainty of 10% to 20% on the binding energies of the various poles. It should be mentioned that the elimination of those nearthreshold left-hand cuts in this work by employing the first-iterated N/D method is rather qualitative and not a rigorous treatment. To investigate the mass spectrum in the $D^{(*)}\Sigma_c^{(*)}$ coupled-channel systems more quantitatively and precisely, a more rigorous and elegant treatment to the uphysical left-hand cuts is required. Nevertheless, there is no doubt that the open-charm partners of those LHCb pentaguark states do exist and are located close to the corresponding $D^{(*)}\Sigma_c^{(*)}$ thresholds. One can expect those P_{cc}^N states around 4 GeV will be observed experimentally in the near future.

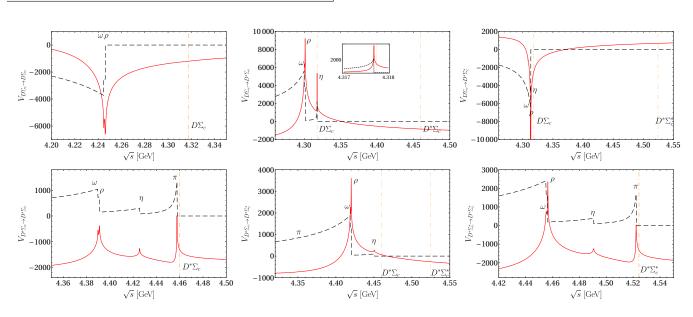


FIG. 5: V matrix for the S-wave $D\Sigma_c$ - $D^*\Sigma_c$ - $D^*\Sigma_c$ coupled-channel system with (I,J)=(1/2,1/2). For notations, see Fig. 3.

IV. SUMMARY

Recently, several pentaquark-like states, P_{ψ}^{N} close to $\bar{D}^{(*)}\Sigma_{c}^{(*)}$ thresholds and $P_{\psi s}^{\Lambda}$ close to the $\bar{D}^{(*)}\Xi_{c}^{(*)}$, were reported by the LHCb Collaboration. All these pentaquark states are suggested to be the hadronic molecule candidates composed of the corresponding mesons close

to the pertinent thresholds, see e.g. Refs. [13–17] (and references therein). The extensions of the mass spectrum of pentaquark states assisted by the heavy quark spin symmetry and flavor symmetry are expected straightforwardly. In this work, we explore the possible open-charm partners of P_{ψ}^{N} states by investigating the S-wave $D^{(*)}\Sigma_{c}^{(*)}$ coupled-channel dynamics. Unitarized scattering amplitudes are constructed by means of the first-

TABLE II: Pole positions obtained from the $5/2^ D^*\Sigma_c^*$, $3/2^ D\Sigma_c^*-D^*\Sigma_c-D^*\Sigma_c^*$ and $1/2^ D\Sigma_c-D^*\Sigma_c-D^*\Sigma_c^*$ systems. The cutoff Λ is varied in the range of (0.6, 0.8) GeV. All the pole positions are given in units of MeV and the corresponding thresholds are listed in the brackets in the first column.

Channel	$1/2^{-}$	3/2-	5/2-
$D\Sigma_c(4318)$	(4223, 4286)	-	-
$D\Sigma_c^*(4382)$	-	(4295, 4356)	-
$D^*\Sigma_c(4460)$	(4442, 4452) + 20i	(4367, 4440) + 1i	-
$\overline{D^*\Sigma_c^*(4524)}$	(4490, 4508) + 10i	(4496, 4518) + 20i	(4415, 4501)

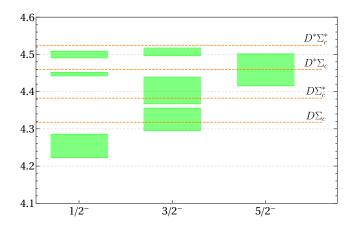


FIG. 6: Mass spectrum of the S-wave $D\Sigma_c^*-D^*\Sigma_c-D^*\Sigma_c^*$ coupled-channel dynamics with (I,J)=(1/2,3/2), S-wave $D\Sigma_c-D^*\Sigma_c-D^*\Sigma_c^*$ system with (I,J)=(1/2,1/2) and S-wave $D^*\Sigma_c^*$ channel with (I,J)=(1/2,5/2), respectively. The cutoff Λ is varied from 0.6 GeV to 0.8 GeV. For the resonances, only the real parts of their pole positions are presented.

iterated N/D method. The open-charm partners of P_{ψ}^{N} are predicted as follows: for $I(J^{P})=1/2(1/2^{-})$, one $D\Sigma_{c}$ bound state with $E_{B}=32\,\mathrm{MeV}$, one $D^{*}\Sigma_{c}$ state with $E_{B}=8\,\mathrm{MeV}$ and one $D^{*}\Sigma_{c}^{*}$ state with $E_{B}=16\,\mathrm{MeV}$; for $I(J^{P})=1/2(3/2^{-})$, one $D\Sigma_{c}^{*}$ bound state with $E_{B}=26\,\mathrm{MeV}$, one $D^{*}\Sigma_{c}$ state with $E_{B}=20\,\mathrm{MeV}$ and one $D^{*}\Sigma_{c}^{*}$ state with $E_{B}=7\,\mathrm{MeV}$; for $I(J^{P})=1/2(5/2^{-})$, one $D^{*}\Sigma_{c}^{*}$ state with $E_{B}=23\,\mathrm{MeV}$. Further investigations are required to solidify these results.

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Appendix A: Effective Lagrangians and exchange potentials

After expanding the effective Lagrangians in Eq. (1), one has

$$\mathcal{L}_{1} = \frac{-\sqrt{2}g}{F_{\pi}} \left[-i\epsilon_{\tau\alpha\mu\nu}v^{\alpha}P_{b}^{*(Q)\mu}\partial^{\nu}\Pi_{ba}P_{a}^{*(Q)\tau\dagger} + P_{b}^{*(Q)\mu}\partial_{\mu}\Pi_{ba}P_{a}^{(Q)\dagger} + P_{b}^{(Q)}\partial^{\nu}\Pi_{ba}P_{a\nu}^{*(Q)\dagger} \right],$$

$$\mathcal{L}_{2} = \sqrt{2}\beta g_{V} \left(P_{b}^{*(Q)\mu}v^{\nu}V_{\nu,ba}P_{\mu,a}^{*(Q)\dagger} - P_{b}^{(Q)}v^{\nu}V_{\nu,ba}P_{a}^{(Q)\dagger} \right) - i2\sqrt{2}\lambda g_{V} \left[P_{b}^{*(Q)\mu}(\partial_{\mu}V_{\nu,ba} - \partial_{\nu}V_{\mu,ba})P_{a}^{*(Q)\nu\dagger} + i\epsilon_{\tau\lambda\nu\alpha}v^{\alpha}P_{b}^{*(Q)}\partial^{\tau}V_{ba}^{\lambda}P_{a}^{*(Q)\nu\dagger} - i\epsilon_{\alpha\mu\tau\lambda}v^{\alpha}P_{b}^{*(Q)\mu}\partial^{\tau}V_{ba}^{\lambda}P_{a}^{(Q)\dagger} \right],$$

$$\mathcal{L}_{S} = \frac{-3g_{1}}{2\sqrt{2}F_{\pi}}v_{\kappa}\epsilon^{\mu\nu\lambda\kappa} \left[\bar{B}_{6\mu,ab}^{*(Q)} + \frac{1}{\sqrt{3}}\bar{B}_{6,ab}^{(Q)}\gamma_{5}\gamma_{\mu} \right] \partial_{\nu}\Pi_{bc} \left[B_{6\lambda,ca}^{*(Q)} - \frac{1}{\sqrt{3}}\gamma_{\lambda}\gamma_{5}B_{6,ca}^{(Q)} \right]$$

$$+ \frac{\beta_{S}g_{V}}{\sqrt{2}} \left[\bar{B}_{6\mu,ab}^{*(Q)} + \frac{1}{\sqrt{3}}\bar{B}_{6,ab}^{(Q)}\gamma_{5}(\gamma_{\mu} + \nu_{\mu}) \right] v^{\alpha}V_{\alpha,bc} \left[B_{6,ca}^{*(Q)\mu} - \frac{1}{\sqrt{3}}(\gamma^{\mu} + v^{\mu})\gamma_{5}B_{6,ca}^{(Q)} \right]$$

$$+ \frac{i\lambda_{S}g_{V}}{\sqrt{2}} \left[\bar{B}_{6\mu,ab}^{*(Q)} + \frac{1}{\sqrt{3}}\bar{B}_{6,ab}^{(Q)}\gamma_{5}(\gamma_{\mu} + v_{\mu}) \right] (\partial^{\mu}V^{\nu} - \partial^{\nu}V^{\mu})_{bc} \left[B_{6\nu,ca}^{*(Q)} - \frac{1}{\sqrt{3}}(\gamma_{\nu} + v_{\nu})\gamma_{5}B_{6,ca}^{(Q)} \right], \tag{A1}$$

in terms of the various field operators

$$P^{(*)(Q)} = \sqrt{M}(D^{(*)0}, D^{(*)+}, D_s^{(*)+}), \ \Pi = \begin{pmatrix} \frac{\pi^0}{\sqrt{2}} + \frac{\eta}{\sqrt{6}} & \pi^+ & K^+ \\ \pi^- & -\frac{\pi^0}{\sqrt{2}} + \frac{\eta}{\sqrt{6}} & K^0 \\ K^- & \bar{K}^0 & -\sqrt{\frac{2}{3}}\eta \end{pmatrix}, \ V = \begin{pmatrix} \frac{\rho^0}{\sqrt{2}} + \frac{\omega}{\sqrt{2}} & \rho^+ & K^{*+} \\ \rho^- & -\frac{\rho^0}{\sqrt{2}} + \frac{\omega}{\sqrt{2}} & K^{*0} \\ K^{*-} & \bar{K}^{*0} & \phi \end{pmatrix},$$

$$B_6^{(Q)} = \begin{pmatrix} \sum_c^{++} & \frac{1}{\sqrt{2}} \sum_c^{+} & \frac{1}{\sqrt{2}} \Xi_c^{\prime+} \\ \frac{1}{\sqrt{2}} \sum_c^{+} & \sum_c^{0} & \frac{1}{\sqrt{2}} \Xi_c^{\prime0} \\ \frac{1}{\sqrt{2}} \Xi_c^{\prime+} & \frac{1}{\sqrt{2}} \Xi_c^{\prime0} & \Omega_c^{0} \end{pmatrix}, \ B_6^{*(Q)} = \begin{pmatrix} \sum_c^{*++} & \frac{1}{\sqrt{2}} \sum_c^{*+} & \frac{1}{\sqrt{2}} \Xi_c^{*+} \\ \frac{1}{\sqrt{2}} \sum_c^{*+} & \sum_c^{*0} & \frac{1}{\sqrt{2}} \Xi_c^{*0} \\ \frac{1}{\sqrt{2}} \Xi_c^{*+} & \frac{1}{\sqrt{2}} \Xi_c^{*0} & \Omega_c^{*0} \end{pmatrix}.$$
(A2)

The values of the coupling constants adopted in the present work are given in Tab. III.

TABLE III: The coupling constants adopted in the calculation. Here, the values are taken from Refs. [20, 53, 54].

g_V	g_1	g	λ_S	λ
5.8	0.94	-0.59	$-3.31/\mathrm{GeV}$	$0.56/{ m GeV}$
β	β_2	β_S	F_{π}	
0.9	-0.9	-1.74	$0.092\mathrm{GeV}$	

For each process, besides the pseudo-potential \mathcal{V} , which can be derived from the effective Lagrangians above, there also exists a combined coupling factor (CF) and an isospin factor (IF). The combined coupling factor comes from the combination of different charged particles for the same process. We collect all the involved combined coupling factors in Tab. IV. The isospin factors are calculated using [55]

$$IF(I) = \sum_{m_1, m_2, m_3, m_4} \langle I_1 m_1 I_2 m_2 | II_z \rangle \langle I_3 m_3 I_4 m_4 | II_z \rangle \times \langle I_3 m_3 I_4 m_4 | M^{iso} | I_1 m_1 I_2 m_2 \rangle, \tag{A3}$$

with the baryon-first convention is applied. All the isospin factors for the considered processes are also given in Tab. IV. Then the transition amplitude with specified quantum numbers is expressed by

$$V = -(CF) \times (IF) \times \mathcal{V}. \tag{A4}$$

Here the minus sign comes from the definition that ensures a negative V corresponds to an attraction interaction.

Appendix B: Partial-wave analysis

Once the potential respecting heavy quark spin symmetry is calculated, we can get the partial-wave potential in the JLS basis using [24, 56]

$$\begin{split} T_{LS;L'S'}^{J}(s) &= \frac{Y_{L'}^{0}(\hat{\mathbf{z}})}{2J+1} \sum_{\substack{s_{1z}, s_{2z}, s_{3z}, \\ s_{4z}, m}} \int \mathrm{d}\hat{\mathbf{p}}' Y_{L}^{m}(\hat{\mathbf{p}}')^{*} \\ &\times (s_{1z}s_{2z}S_{z}|s_{1}s_{2}S)(mS_{z}S'_{z}|LSJ) \\ &\times (s_{3z}s_{4z}S'_{z}|s_{3}s_{4}S')(0S'_{z}S'_{z}|L'S'J)T(s) \\ &\stackrel{m=0}{=} \frac{\sqrt{(2L'+1)(2L+1)}}{2J+1} \sum_{\substack{s_{1z}, s_{2z}, \\ s_{2z}, s_{3z}, s_{3z},$$

$$\times P_L(\cos\theta)(s_{1z}s_{2z}S_z|s_1s_2S)(0S_zS_z'|LSJ) \times (s_{3z}s_{4z}S_z'|s_3s_4S')(0S_z'S_z'|L'S'J)T(s).$$
 (B1)

Since only the S-wave term is considered in the present work, e.g. $L=L^\prime=0,$ we then have the following expression

$$T_{0S;0S'}^{J}(s) = \frac{1}{2J+1} \sum_{\substack{s_{1z}, s_{2z}, \\ s_{3z}, s_{4z}}} \frac{1}{2} \int_{-1}^{1} d\cos\theta(s_{1z}s_{2z}S_z | s_1s_2S) \times (s_{3z}s_{4z}S_z | s_3s_4S')T(s) . \tag{B2}$$

TABLE IV: The combined coupling factors and isospin factors for all the involved processes. It implies the first seven reactions contains the π , η , ρ and ω exchange, while the last two only have the ρ and ω exchange.

Process	Exchanged Particle	CF	IF
$\begin{array}{c} \Sigma_c D \to \Sigma_c D^* \\ \Sigma_c D \to \Sigma_c^* D^* \end{array}$	π	$\frac{1}{2}$	2
$\Sigma_c^* D \to \Sigma_c D^*$ $\Sigma_c^* D \to \Sigma_c^* D^*$ $\Sigma_c D^* \to \Sigma_c D^*$ $\Sigma_c D^* \to \Sigma_c^* D^*$ $\Sigma_c D^* \to \Sigma_c^* D^*$ $\Sigma_c^* D^* \to \Sigma_c^* D^*$	η	$\frac{1}{6}$	1
	ρ	$\frac{1}{2}$	2
	ω	$\frac{1}{2}$	1
$\Sigma_c D \to \Sigma_c D$	ω	$\frac{1}{2}$	1
$\Sigma_c^* D \to \Sigma_c^* D$	ρ	$\frac{1}{2}$	2

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