Multiquark contributions to charm baryon spectroscopy

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We study possible multiquark contributions to the charm baryon spectrum by considering higher order Fock space components. For this purpose we perform two different calculations. In a first approach we do a coupled-channel calculation of the ND system looking for molecular states. In a second step we allow for the coupling to a heavy baryon–light meson two-hadron system looking for compact exotic five-quark structures. Both calculations have been done within the framework of a chiral constituent quark model. The model, tuned in the description of the baryon and meson spectra as well as the NN interaction, provides parameter-free predictions for charm $+1$ molecular or compact two-hadron systems. Unlike the $N\bar{D}$ system, no sharp quark-Pauli effects are found. However, the existence of different two-hadron thresholds for the five-quark system will make the coupled-channel dynamics relevant. Only a few channels are candidates to lodge molecular or compact hadrons with a five-quark structure, being specially relevant the $(T)J^P = (0)1/2^-$ and $(T)J^P = (2)5/2^-$ channels. The identification of molecular states and/or compact hadrons with multiquark components either with or without exotic quantum numbers is a challenge of different collaborations like PANDA, LHCb, ExHIC or J-PARC.

I. INTRODUCTION

One of the most basic problems of QCD is to identify all the clusters of quarks, antiquarks and gluons that are sufficiently bound by QCD interactions so that they are either stable particles or appropriately long lived to be observed as resonances [1]. To this respect, charm hadron physics has become a cornerstone due to the experimental findings during the past decade. On the one hand, one encounters the outstanding discovery in charmonium spectroscopy of the flagship of the so-called XYZ states, the $X(3872)$ [2]. Before this discovery, and based on Gell-Mann conjecture [3], the hadronic experimental data were classified either as $q\bar{q}$ or qqq states according to $SU(3)$ irreducible representations. However, since 2003 more than 20 meson resonances reported by different experimental collaborations, most of them close to a two-meson threshold, had properties that made a simple quark-antiquark structure unlikely [4]. Although this observation could be coincidental due to the large number of thresholds in the energy region where the XYZ mesons have been reported, it could also point to a close relation between some particular thresholds and resonances contributing to the standard quark-antiquark heavy meson spectroscopy. On the other hand, a similar situation has arisen in the charm baryon spectrum during the past years with the advent of a large set of experimental data (see Refs. [5,6] for a comprehensive update of the experimental and theoretical situation of the heavy baryon spectra). The properties of some excited states show an elusive nature as three-quark systems. Like the charmonium spectrum, some of them are rather close to a baryon-meson threshold suggesting a possible molecular or compact structure [7–13]. It has been

already highlighted within a simple toy model the key role that S wave meson-baryon thresholds may play in matching poor light-baryon mass predictions from quark models with data [14]. Thus, the analysis of possible multiquark contributions close to meson-baryon thresholds with a full-fledged quark dynamical model could help in the understanding of heavy baryon spectroscopy.

The existence of molecular contributions in the charm baryon spectrum stems primarily on the interaction between charm mesons and nucleons, which on the other hand has turned into an interesting subject in several contexts [15]. It is particularly interesting for the study of chiral symmetry restoration in a hot and/or dense medium [16]. It will also help in the understanding of the suppression of the J/Ψ production in heavy ion collisions [17]. Besides, it may shed light on the possible existence of exotic nuclei with heavy flavors [18,19]. Experimentally, it will become possible to analyze the interaction of charm mesons with nucleons inside nuclear matter with the operation of the FAIR facility at the GSI laboratory in Germany [15]. There are proposals for experiments by the PANDA collaboration to produce D mesons by annihilating antiprotons on the deuteron. This could be achieved with an antiproton beam, by tuning the antiproton energy to one of the higher-mass charmonium states that decays into open charm mesons. These experimental ideas may become plausible based on recent estimations of the cross section for the production of $D\bar{D}$ pairs in proton-antiproton collisions [20]. There are also different theoretical estimations about the production rate at PANDA of Λ_c baryons through the direct process $p\bar{p} \to \Lambda_c \bar{\Lambda}_c$ [21,22]. The identification of hadronic molecular states and/or hadrons with multiquark components either

with or without exotic quantum numbers is also a challenge in relativistic heavy ion collisions offering a promising resolution to this problem [23,24]. Besides the LHCb [25] and CDF [26] collaborations are providing a huge data set of new measurements of heavy flavor spectroscopy. In the coming future, J-PARC also intends to contribute to the experimental measurement of exotic baryons. Thus, a good knowledge of the interaction of charm mesons with ordinary hadrons, like nucleons or $\Delta's$, is a prerequisite.

Before one can infer in a sensitive manner changes of the interaction in the medium [27,28], a reasonable understanding of the interaction in free space is required. However, here one has to manage with an important difficulty, namely the complete lack of experimental data at low energies for the free-space interaction. Thus, the generalization of models describing the two-hadron interaction in the light flavor sector could offer insight about the unknown interaction of hadrons with heavy flavors. This is the main purpose of this paper, to make use of a chiral constituent quark model describing the NN interaction [29] as well as the meson spectrum in all flavor sectors [30] to obtain parameter-free predictions that may be testable in future experiments. Such a project was already undertaken for the interaction between two charm mesons [31] and also for the interaction between anticharm mesons and nucleons [32] which encourages us in the present challenge.

The paper is organized as follows. In Sec. II we will first present a brief description of the quark-model wave function for the baryon-meson system. We will later revisit the interacting potential and finally we will summarize the solution of the two-body problem by means of the Fredholm determinant. In Sec. III we present and discuss our results. We will first briefly discuss the baryon-meson interaction in comparison to hadronic models. We will analyze the character of the different isospin-spin channels, looking for the attractive ones that may lodge resonances either as a molecule or as a compact five-quark state, to be measured by experiment. We will also compare with existing results in the literature. Finally, in Sec. IV we summarize our main conclusions.

II. THE BARYON-MESON SYSTEM

A. The baryon-meson wave function

In order to describe the baryon-meson system we shall use a constituent quark cluster model, i.e., hadrons are described as clusters of quarks and antiquarks. Assuming a two-center shell model the wave function of an arbitrary baryon-meson system, a baryon B_i and a meson M_j, can be written as

$$\Psi_{B_i M_j}^{\text{LST}}(\vec{R}) = \mathcal{A}\left[B_i\left(123; -\frac{\vec{R}}{2}\right)M_j\left(4\bar{5}; +\frac{\vec{R}}{2}\right)\right]^{\text{LST}}, \quad (1)$$

where $\mathcal{A}$ is the antisymmetrization operator accounting for the possible existence of identical quarks inside the

hadrons. In the case we are interested in, baryon-meson systems made of N or Δ baryons and D or D^* mesons, no identical quarks can be exchanged between the baryon and the meson and thus no sharp quark-Pauli effects are expected.

If we assume Gaussian $0s$ wave functions for the quarks inside the hadrons, the normalization of the baryon-meson wave function $\Psi_{B_i M_j}^{\text{LST}}(\vec{R})$ of Eq. (1) can be expressed as

$$\mathcal{N}_{B_i M_j}^{\text{LST}}(R)$$
$$= 4\pi \exp\left\{-\frac{R^2}{8}\left(\frac{4}{b^2} + \frac{1}{b_c^2}\right)\right\} i_{L+1/2}\left[\frac{R^2}{8}\left(\frac{4}{b^2} + \frac{1}{b_c^2}\right)\right], \tag{2}$$

where, for the sake of generality, we have assumed different Gaussian parameters for the wave function of the light quarks (b) and the heavy quark (b_c). In the limit where the two hadrons overlap ($R \to 0$), the Pauli principle does not impose any antisymmetry requirement. This can be easily checked for the $L = 0$ partial waves, where such effects would be prominent. Using the asymptotic form of the Bessel functions, $i_{L+1/2}$, we obtain the S wave normalization kernel in the overlapping region that behaves like a constant for $R = 0$,

$$\mathcal{N}_{B_i M_j}^{L=0 ST} \xrightarrow[R\to0]{} 4\pi\left\{1 - \frac{R^2}{8}\left(\frac{4}{b^2} + \frac{1}{b_c^2}\right)\right\}$$
$$\times \left\{1 + \frac{1}{6}\left(\frac{R^2}{8b_c^2}\right)^2\left(1 + \frac{4b_c^2}{b^2}\right)^2 + \cdots\right\}. \tag{3}$$

B. The two-body interactions

The two-body interactions involved in the study of the baryon-meson system are obtained from the chiral constituent quark model [29]. This model was proposed in the early 1990's in an attempt to obtain a simultaneous description of the nucleon-nucleon interaction and the baryon spectra. It was later generalized to all flavor sectors [30]. In this model hadrons are described as clusters of three interacting massive (constituent) quarks, the mass coming from the spontaneous breaking of the original $SU(2)_L \otimes SU(2)_R$ chiral symmetry of the QCD Lagrangian. QCD perturbative effects are taken into account through the one-gluon-exchange (OGE) potential [33]. It reads

$$V_{\text{OGE}}(\vec{r}_{ij}) = \frac{\alpha_s}{4}\vec{\lambda}_i^c \cdot \vec{\lambda}_j^c\left\{\frac{1}{r_{ij}} - \frac{1}{4}\left(\frac{1}{2m_i^2} + \frac{1}{2m_j^2} + \frac{2\vec{\sigma}_i \cdot \vec{\sigma}_j}{3m_i m_j}\right)\right.$$
$$\left.\times \frac{e^{-r_{ij}/r_0}}{r_0^2 r_{ij}} - \frac{3S_{ij}}{4m_q^2 r_{ij}^3}\right\}, \tag{4}$$

where λ^c are the $SU(3)$ color matrices, $r_0 = \hat{r}_0/\mu$ is a flavor-dependent regularization scaling with the reduced mass of the interacting pair, and α_s is the scale-dependent strong coupling constant given by [30]

$$\alpha_s(\mu) = \frac{\alpha_0}{\ln[(\mu^2 + \mu_0^2)/\gamma_0^2]}, \qquad (5)$$

where $\alpha_0 = 2.118$, $\mu_0 = 36.976$ MeV and $\gamma_0 = 0.113$ fm^{-1}. This equation gives rise to $\alpha_s \sim 0.54$ for the light-quark sector and $\alpha_s \sim 0.43$ for uc pairs.

Nonperturbative effects are due to the spontaneous breaking of the original chiral symmetry at some momentum scale. In this domain of momenta, light quarks interact through Goldstone boson exchange potentials,

$$V_\chi(\vec{r}_{ij}) = V_{\mathrm{OSE}}(\vec{r}_{ij}) + V_{\mathrm{OPE}}(\vec{r}_{ij}), \qquad (6)$$

where the one-pion exchange (OPE) and one-sigma exchange (OSE) potentials are given by

$$V_{\mathrm{OSE}}(\vec{r}_{ij}) = -\frac{g_{\mathrm{ch}}^2}{4\pi} \frac{\Lambda^2}{\Lambda^2 - m_\sigma^2} m_\sigma \left[Y(m_\sigma r_{ij}) - \frac{\Lambda}{m_\sigma} Y(\Lambda r_{ij}) \right],$$

$$V_{\mathrm{OPE}}(\vec{r}_{ij}) = \frac{g_{\mathrm{ch}}^2}{4\pi} \frac{m_\pi^2}{12 m_i m_j} \frac{\Lambda^2}{\Lambda^2 - m_\pi^2} m_\pi \left\{ \left[Y(m_\pi r_{ij}) - \frac{\Lambda^3}{m_\pi^3} Y(\Lambda r_{ij}) \right] \vec{\sigma}_i \cdot \vec{\sigma}_j \right.$$

$$\left. + \left[H(m_\pi r_{ij}) - \frac{\Lambda^3}{m_\pi^3} H(\Lambda r_{ij}) \right] S_{ij} \right\} (\vec{\tau}_i \cdot \vec{\tau}_j). \qquad (7)$$

$g_{\mathrm{ch}}^2/4\pi$ is the chiral coupling constant, $Y(x)$ is the standard Yukawa function defined by $Y(x) = e^{-x}/x$, $S_{ij} = 3(\vec{\sigma}_i \cdot \hat{r}_{ij})(\vec{\sigma}_j \cdot \hat{r}_{ij}) - \vec{\sigma}_i \cdot \vec{\sigma}_j$ is the quark tensor operator, and $H(x) = (1 + 3/x + 3/x^2)Y(x)$.

Finally, any model imitating QCD should incorporate confinement. Being a basic term from the spectroscopic point of view it is negligible for the hadron-hadron interaction. Lattice calculations suggest a screening effect on the potential when increasing the interquark distance [34],

$$V_{\mathrm{CON}}(\vec{r}_{ij}) = \{-a_c(1 - e^{-\mu_c r_{ij}})\}(\vec{\lambda}^c_i \cdot \vec{\lambda}^c_j). \qquad (8)$$

Once perturbative (one-gluon exchange) and nonperturbative (confinement and chiral symmetry breaking) aspects of QCD have been considered, one ends up with a quark-quark interaction of the form

$$V_{q_i q_j}(\vec{r}_{ij})$$
$$= \begin{cases} [q_i q_j = nn] \Rightarrow V_{\mathrm{CON}}(\vec{r}_{ij}) + V_{\mathrm{OGE}}(\vec{r}_{ij}) + V_\chi(\vec{r}_{ij}) \\ [q_i q_j = cn] \Rightarrow V_{\mathrm{CON}}(\vec{r}_{ij}) + V_{\mathrm{OGE}}(\vec{r}_{ij}), \end{cases}$$
$$(9)$$

where n stands for the light quarks u and d. Notice that for the particular case of heavy quarks (c or b) chiral symmetry is explicitly broken and therefore boson exchanges do not contribute. The parameters of the model are those of Ref. [32]. The model guarantees a nice description of the baryon (N and Δ) [35] and the meson (D and D^*) spectra [30].

In order to derive the local $B_n M_m \to B_k M_l$ interaction from the basic qq interaction defined above, we use a Born-Oppenheimer approximation. Explicitly, the potential is calculated as follows:

$$V_{B_n M_m(LST) \to B_k M_l(L'S'T)}(R) = \xi_{LST}^{L'S'T}(R) - \xi_{LST}^{L'S'T}(\infty), \qquad (10)$$

where

$$\xi_{LST}^{L'S'T}(R)$$
$$= \frac{\langle \Psi_{B_k M_l}^{L'S'T}(\vec{R})| \sum_{i<j=1}^5 V_{q_i q_j}(\vec{r}_{ij}) |\Psi_{B_n M_m}^{LST}(\vec{R})\rangle}{\sqrt{\langle \Psi_{B_k M_l}^{L'S'T}(\vec{R})|\Psi_{B_k M_l}^{L'S'T}(\vec{R})\rangle} \sqrt{\langle \Psi_{B_n M_m}^{LST}(\vec{R})|\Psi_{B_n M_m}^{LST}(\vec{R})\rangle}}.$$
$$(11)$$

In the last expression the quark coordinates are integrated out keeping R fixed, the resulting interaction being a function of the baryon-meson relative distance. The wave function $\Psi_{B_n M_m}^{LST}(\vec{R})$ for the baryon-meson system has been discussed in Sec. II A.

We show in Fig. 1 the different diagrams contributing to the baryon-meson interaction. As compared to the $N\bar{D}$ case [32] and due to the absence of quark-exchange
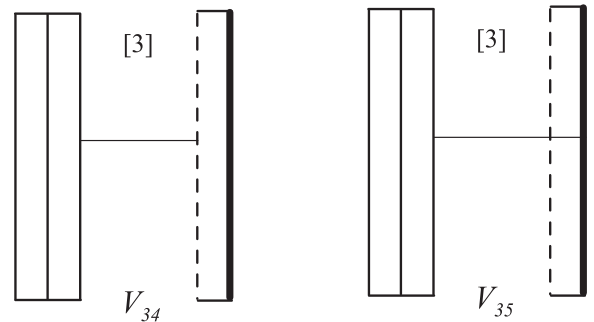

FIG. 1. Different diagrams contributing to the baryon-meson interaction. The vertical thin solid lines represent light quarks, the vertical thick solid line represents a heavy quark, the vertical dashed line stands for the light antiquark, and the horizontal solid line represents the exchanged particle. The number between square brackets stands for the number of diagrams topologically equivalent.

contributions, the number of diagrams is greatly reduced, getting just purely hadronic interactions.

C. Integral equations for the two-body systems

To study the possible existence of molecular states made of a light baryon, N or Δ, and a charmed meson, D or D^*, we have solved the Lippmann-Schwinger equation for negative energies looking at the Fredholm determinant $D_F(E)$ at zero energy [36]. If there are no interactions then $D_F(0) = 1$, if the system is attractive then $D_F(0) < 1$, and if a bound state exists then $D_F(0) < 0$. This method permitted us to obtain robust predictions even for zero-energy bound states, and gave information about attractive channels that may lodge a resonance in similar systems [31]. We consider a baryon-meson system $B_i M_j$ ($B_i = N$ or Δ and $M_j = D$ or D^*) in a relative S state interacting through a potential V that contains a tensor force. Then, in general, there is a coupling to the $B_i M_j$ D wave. Moreover, the baryon-meson system can couple to other baryon-meson states. We show in the first row of each spin cell of Table I the lowest light baryon—charm meson coupled channels in the isospin-spin (T, J) basis. They would contribute to our first approach, a coupled-channel calculation of the ND system looking for molecular states. As we have done in Ref. [31], we will later allow for the rearrangement of quarks at short distances giving rise to a coupling to a charm baryon—light meson two-hadron system, through the diagram represented in Fig. 2. For this case we show in the second row of each spin cell of Table I, between square brackets, the additional channels contributing to each (T, J) state. They would contribute to the second calculation, looking for compact five-quark states.

Thus, if we denote the different baryon-meson systems as channel A_i, the Lippmann-Schwinger equation for the baryon-meson scattering becomes

$$
t_{\alpha\beta;TJ}^{\ell_\alpha s_\alpha, \ell_\beta s_\beta}(p_\alpha, p_\beta; E) = V_{\alpha\beta;TJ}^{\ell_\alpha s_\alpha, \ell_\beta s_\beta}(p_\alpha, p_\beta) + \sum_{\gamma=A_1,A_2,\cdots} \sum_{\ell_\gamma=0,2}
$$
$$
\times \int_0^\infty p_\gamma^2 dp_\gamma V_{\alpha\gamma;TJ}^{\ell_\alpha s_\alpha, \ell_\gamma s_\gamma}(p_\alpha, p_\gamma)
$$
$$
\times G_\gamma(E; p_\gamma) t_{\gamma\beta;TJ}^{\ell_\gamma s_\gamma, \ell_\beta s_\beta}(p_\gamma, p_\beta; E);
$$
$$
\alpha, \beta = A_1, A_2, \ldots, \tag{12}
$$

where t is the two-body scattering amplitude, T, J, and E are the isospin, total angular momentum and energy of the system, $\ell_\alpha s_\alpha$, $\ell_\gamma s_\gamma$, and $\ell_\beta s_\beta$ are the initial, intermediate, and final orbital angular momentum and spin, respectively, and p_γ is the relative momentum of the two-body system γ. The propagators $G_\gamma(E; p_\gamma)$ are given by

$$
G_\gamma(E; p_\gamma) = \frac{2\mu_\gamma}{k_\gamma^2 - p_\gamma^2 + i\epsilon}, \tag{13}
$$

TABLE I.   Interacting baryon-meson channels in the isospin-spin (T, J) basis. See text for details.

	$T = 0$	$T = 1$	$T = 2$
$J = 1/2$	$ND - ND^*$	$ND - ND^* - \Delta D^*$	ΔD^*
	$[\Sigma_c \pi - \Lambda_c^+ \eta]$	$[\Lambda_c^+ \pi - \Sigma_c \pi]$	$[\Sigma_c \pi]$
$J = 3/2$	ND^*	$ND^* - \Delta D - \Delta D^*$	$\Delta D - \Delta D^*$
	$[\Sigma_c^* \pi - \Lambda_c^+ \omega]$	$[\Sigma_c^* \pi - \Lambda_c^+ \rho]$	$[\Sigma_c^* \pi]$
$J = 5/2$	$\cdots$	ΔD^*	ΔD^*
	$[\Sigma_c^* \rho]$	$[\Sigma_c^* \rho - \Sigma_c^* \omega]$	$[\Sigma_c^* \rho]$

with

$$
E = \frac{k_\gamma^2}{2\mu_\gamma}, \tag{14}
$$

where μ_γ is the reduced mass of the two-body system γ. For bound-state problems $E < 0$ so that the singularity of the propagator is never touched and we can forget the $i\epsilon$ in the denominator. If we make the change of variables

$$
p_\gamma = d \frac{1 + x_\gamma}{1 - x_\gamma}, \tag{15}
$$

where d is a scale parameter, and the same for p_α and p_β, we can write Eq. (12) as

$$
t_{\alpha\beta;TJ}^{\ell_\alpha s_\alpha, \ell_\beta s_\beta}(x_\alpha, x_\beta; E) = V_{\alpha\beta;TJ}^{\ell_\alpha s_\alpha, \ell_\beta s_\beta}(x_\alpha, x_\beta) + \sum_{\gamma=A_1,A_2,\cdots} \sum_{\ell_\gamma=0,2}
$$
$$
\times \int_{-1}^1 d^2 \left(\frac{1 + x_\gamma}{1 - x_\gamma}\right)^2 \frac{2d}{(1 - x_\gamma)^2} dx_\gamma
$$
$$
\times V_{\alpha\gamma;TJ}^{\ell_\alpha s_\alpha, \ell_\gamma s_\gamma}(x_\alpha, x_\gamma) G_\gamma(E; p_\gamma)
$$
$$
\times t_{\gamma\beta;TJ}^{\ell_\gamma s_\gamma, \ell_\beta s_\beta}(x_\gamma, x_\beta; E). \tag{16}
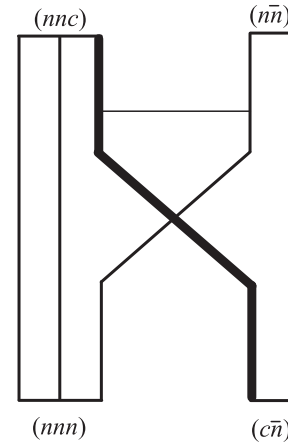$$

FIG. 2.   Diagram representing the coupling between a light baryon—charm meson channel and a charm baryon—light meson two-hadron system.

We solve this equation by replacing the integral from -1 to 1 by a Gauss-Legendre quadrature which results in the set of linear equations

$$\sum_{\gamma=A_1,A_2,\cdots}\sum_{\ell_\gamma=0,2}\sum_{m=1}^{N}M_{\alpha\gamma;TJ}^{n\ell_\alpha s_\alpha,m\ell_\gamma s_\gamma}(E)t_{\gamma\beta;TJ}^{\ell_\gamma s_\gamma,\ell_\beta s_\beta}(x_m,x_k;E)$$

$$= V_{\alpha\beta;TJ}^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}(x_n,x_k),\tag{17}$$

with

$$M_{\alpha\gamma;TJ}^{n\ell_\alpha s_\alpha,m\ell_\gamma s_\gamma}(E)=\delta_{nm}\delta_{\ell_\alpha\ell_\gamma}\delta_{s_\alpha s_\gamma}-w_m d^2\left(\frac{1+x_m}{1-x_m}\right)^2\frac{2d}{(1-x_m)^2}$$

$$\times V_{\alpha\gamma;TJ}^{\ell_\alpha s_\alpha,\ell_\gamma s_\gamma}(x_n,x_m)G_\gamma(E;p_{\gamma m}),\tag{18}$$

and where w_m and x_m are the weights and abscissas of the Gauss-Legendre quadrature while $p_{\gamma m}$ is obtained by putting $x_\gamma = x_m$ in Eq. (15). If a bound state exists at an energy E_B, the determinant of the matrix $M_{\alpha\gamma;TJ}^{n\ell_\alpha s_\alpha,m\ell_\gamma s_\gamma}(E_B)$ vanishes, i.e., $|M_{\alpha\gamma;TJ}(E_B)| = 0$.

III. RESULTS AND DISCUSSION

Regarding the ND interaction there are general trends that can be briefly summarized. It is worth noting the absence of quark-exchange diagrams, which also prohibits the OGE contribution, and thus quark-exchange effects are not present. Thus, the interaction comes determined by the OPE and OSE. For very-long distances ($R > 4$ fm) the dominant term is the OPE potential, since it corresponds to the longest-range piece. The OPE is also responsible altogether with the OSE for the long-range part behavior (1.5 fm $< R < 4$ fm), due to the combined effect of shorter range and a bigger strength for the OSE as compared to the OPE. The OSE gives the dominant contribution in the intermediate range (0.8 fm $< R < 1.5$ fm), determining the attractive character of the potential in this region. The short-range ($R < 0.8$ fm) potential is either repulsive or attractive depending on the balance between the OSE and OPE. Due to the nonexistence of quark-Pauli correlations from the norm as well as from the interacting potential, one gets a genuine baryonic interaction. Thus, dynamical quark-exchange effects do not play a relevant role in the ND interaction unlike the $N\bar{D}$ case.

Using the interactions described above, we have solved the coupled-channel problem of the baryon-meson systems made of a baryon, N or Δ, and a meson, D or D^*, as explained in Sec. II C. The existence of bound states or resonances will generate baryonic states with charm $+1$ that could be identified as some of the excited states measured in the charm baryon spectrum. In Table II we summarize the character of the interaction in the different (T,J) channels. It can be observed that due to the absence of quark-Pauli correlations and the contribution of the OSE the interaction is in general attractive, giving rise to states

TABLE II. Character of the interaction in the different $ND\,(T,J)$ channels.

	$T=0$	$T=1$	$T=2$
$J=1/2$	Attractive	Weak	Weakly attractive
$J=3/2$	Weak	Attractive	Attractive
$J=5/2$	$\cdots$	Weakly attractive	Attractive

that appear close to different thresholds. We have represented in Fig. 3 the masses and quantum numbers of the possible molecular ND states. The strongest interaction is obtained in the $ND^*\,(T)J^P = (0)1/2^-$ channel that it is coupled to the $ND\,(T)J^P = (0)1/2^-$ partial wave (see Table I), generating the best candidate to lodge a molecule. The expectation value of the isospin operator, -3 for isosinglet and $+1$ for isotriplet states, would reduce the attraction of isotriplet channels as compared to attractive isosinglet channels with the same spin J and vice versa, as can be easily checked in the first two columns of Table II.

Our results may be compared to those of Ref. [10] where the ND system has been analyzed by means of a hadronic model using Lagrangians satisfying heavy quark symmetry and chiral symmetry. They arrive at the same conclusion that the $(T)J^P = (0)1/2^-$ channel is the most attractive one. In Ref. [10] this channel presents a bound state of around 14.4 MeV for the model including only pion exchanges. The main difference of our results with those of Ref. [10] stems from the contribution of the scalar interaction and the consideration of explicit Δ degrees of freedom in our calculation. In a hadronic theory without explicit Δ degrees of freedom only a few channels survive and the coupled-channel dynamics would become simpler (see Table I). As a consequence, for example, one could not get $T = 2$ channels. In the chiral constituent quark model the importance of the Δ degrees of freedom is known since
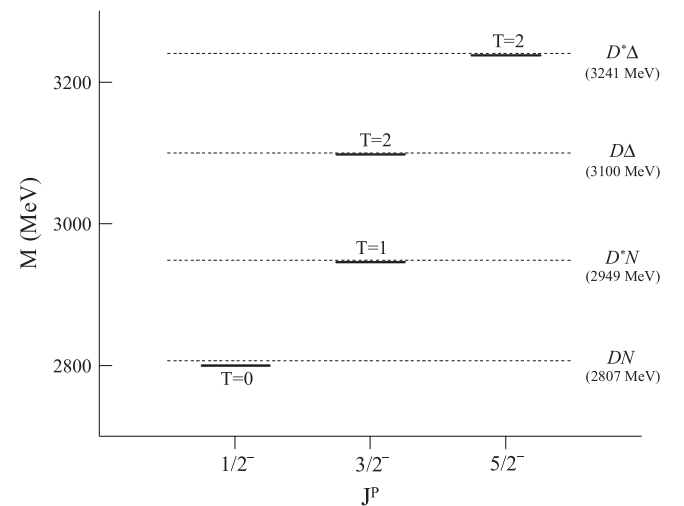

FIG. 3. Masses and quantum numbers of molecular DN states. The dashed lines stand for the different two-hadron thresholds.
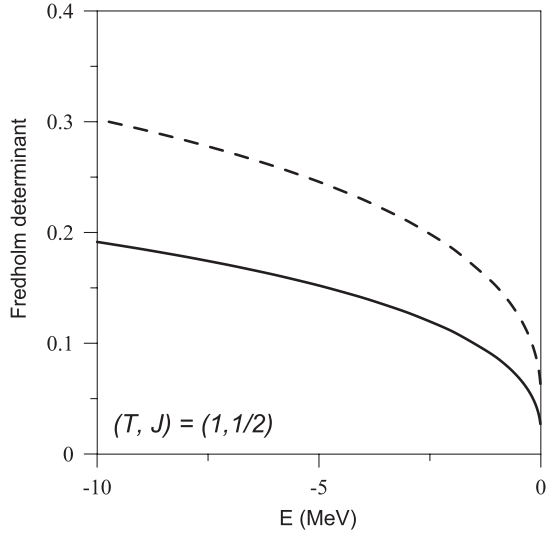
FIG. 4. Fredholm determinant of the $(T)J^P = (1)1/2^-$ channel considering all ND contributions of Table I (solid line) and neglecting the Δ degrees of freedom (dashed line).

long time ago [37,38]. It provides us with an isospin dependent mechanism that allows to correctly describe the low-energy NN S wave phase shifts through a coupled-channel effect, giving an important attractive contribution for the $^1S_0$ NN partial wave. To emphasize the importance of coupled-channel dynamics [39], we have repeated the calculation explained in Sec. II C for the $(T)J^P = (1)1/2^-$ channel but suppressing the states containing $\Delta's$. As can

be seen in Fig. 4, neglecting the Δ degrees of freedom the Fredholm determinant gets larger, indicating a loss of attraction. For the single channel calculation the $(T)J^P = (1)1/2^-$ bound state does not appear, in agreement with the conclusions of Ref. [10].

The DN molecular states appearing in Fig. 3 could be an important ingredient of the charm baryon spectrum. It has been recently suggested [12] the possibility of the $\Sigma_c(2800)$ being an S wave DN molecular state with $J^P = 1/2^-$ and the $\Lambda_c(2940)^+$ an S wave D^*N state with $J^P = 3/2^-$, which would agree rather well with the picture shown in Fig. 3. One may also find an experimental candidate for the $(T)J^P = (2)5/2^-$ resonance in one of the states recently reported by the *BABAR* collaboration [40], an unexplained structure with a mass of 3250 MeV/c^2 in the $\Sigma_c^{++}\pi^-\pi^-$ invariant mass. This state has also been recently suggested as a possible pentaquark [11], something that would be relevant in the second part of our discussion. In spite of this agreement, one should note that the assignment of quantum numbers to baryon resonances on the charm baryon spectrum [41] and the identification of their internal structure [7–13,39] is still an open issue that needs of further experimental analysis and also theoretical efforts. Such uncertainty has been recently revitalized by empha-sizing the potential importance of the relativistic kinematics of the light quark pair [42] casting doubts even on the assignment of quantum numbers to experimental states just based on the nonrelativistic quark model.

One could also find contributions to the charm baryon spectrum with a more involved structure such as compact
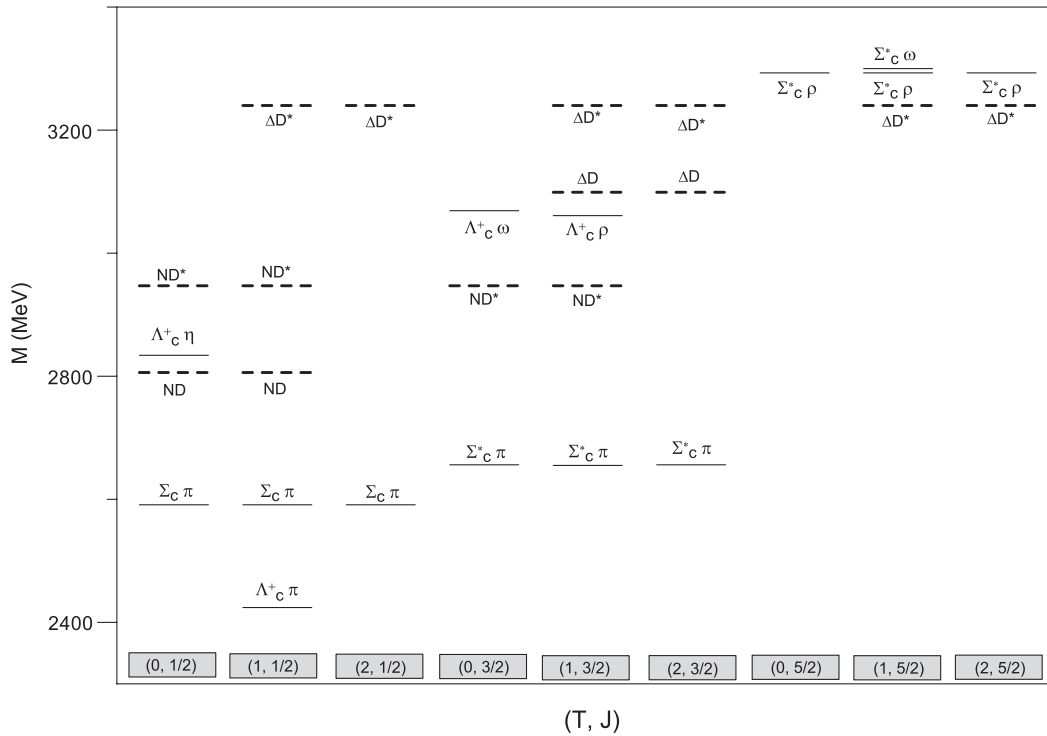

FIG. 5. Different two-body channels contributing to each set of (T, J) quantum numbers as shown in Table I.

TABLE III.   Character of the interaction in the different baryon-meson (T, J) channels.

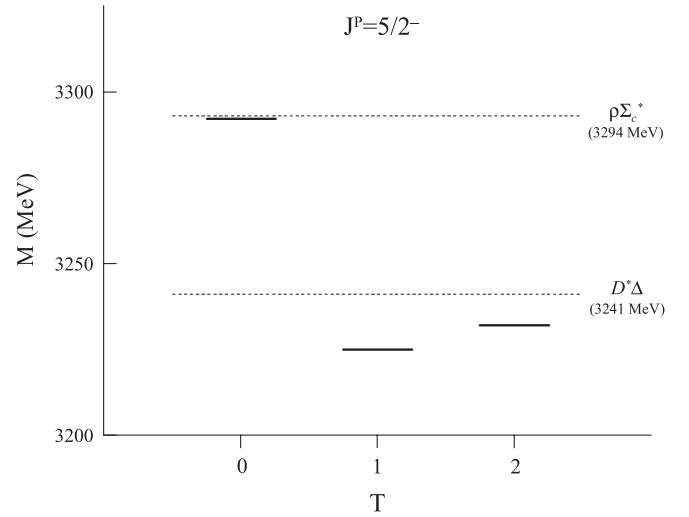|	$T = 0$	$T = 1$	$T = 2$
$J = 1/2$	Weakly attractive	Weak	Strongly repulsive
$J = 3/2$	Weak	Weak	Weak
$J = 5/2$	Attractive	Attractive	Attractive

FIG. 6.   Masses and quantum numbers of compact five-quark states. The dashed lines stand for the different two-hadron thresholds.

five-quark states beyond simple ND resonances [43]. The study of these contributions requires from a full coupled-channel approach including all possible physical states contributing to a given set of quantum numbers (T, J), as has been demonstrated in Ref. [44] for the charmonium spectrum. Standard mesons $(q\bar{q})$ and baryons (qqq) are the only clusters of quarks where it is not possible to construct a color singlet using a subset of their constituents. Thus, $q\bar{q}$ and qqq states are proper solutions of the two- and three-quark Hamiltonian, respectively, corresponding in all cases to bound states. This, however, is not the case for multi-quark combinations, and in particular for five-quark states addressing the baryon spectrum. Thus, when dealing with higher order Fock space contributions to baryon spectros-copy, one has to discriminate between possible five-quark bound states or resonances and simple pieces of the baryon-meson continuum. For this purpose, one has to analyze the two-hadron states that constitute the possible thresholds for each set of quantum numbers. These thresholds have to be determined assuming quantum number conservation within exactly the same scheme (parameters and interactions) used for the five-body calculation. Working with strongly interacting particles, a baryon-meson state should have well-defined total angular momentum (J) and parity (P). If noncentral forces are not considered, orbital angular momentum (L) and total spin (S) are also good quantum numbers. We have represented in Fig. 5 the different two-hadron thresholds contributing to each set of (T, J) quantum numbers.

Given a general five-quark state contributing to the ND wave function, $(nnnQ\bar{n})$ (in the following n stands for a light quark and Q for a heavy c or b quark), two different thresholds are allowed, $(nnn)(Q\bar{n})$ and $(nnQ)(n\bar{n})$. A very simple property [45] of the ground state solutions of the Schrödinger $(q_1\bar{q}_2)$ two-body problem is that they are concave in $(m_{q_1}^{-1} + m_{q_2}^{-1})$, and hence $M_{Q\bar{n}} + M_{\bar{Q}n} \geq M_{Q\bar{Q}} + M_{n\bar{n}}$. This property is enforced both by nature[1] and by all models in the literature unless forced to do otherwise. It implies that in all relevant cases the lowest two-meson threshold for any $(Qn\bar{Q}\,\bar{n})$ state will be the one made of quarkonium-light mesons, i.e., $(Q\bar{Q})(n\bar{n})$ (see Fig. 1 of Ref. [46]). A straightforward generalization of this property to the five-quark system could be obtained within a quark-diquark model if $m_{q_1} \leq m_{q_2} \leq m_{q_3}$. Then

$M_{q_3\bar{q}_2} + M_{q_1\bar{q}_1} \leq M_{q_3\bar{q}_1} + M_{q_1\bar{q}_2}$, because the intervals in $1/\mu$ of the left-hand side and right-hand side have the same middle, but the left-hand side one is wider that the right-hand side one. Now, in a crude quark-diquark model, one can translate this as $M_{q_3q_1q_1} + M_{q_1\bar{q}_1} \leq M_{q_3\bar{q}_1} + M_{q_1q_1q_1}$, as it is observed in Fig. 5 except for the higher spin states where the angular momentum coupling rules impose further restrictions.[2]

Hence, as we have already illustrated in Fig. 4 an important source of attraction might be the coupled-channel effect of the two thresholds, $(nnn)(Q\bar{n}) \leftrightarrow (nnQ)(n\bar{n})$ [39]. Thus, to check the efficiency of this mechanism, we have repeated the calculation of Sec. II C but considered all physical states reflected in Table I. We have represented in Fig. 5 the lowest baryon-meson thresholds contributing to each set of (T, J) quantum numbers. In Table III we have summarized the character of the interaction in the different (T, J) channels. When the $(nnn)(Q\bar{n})$ and $(nnQ)(n\bar{n})$ thresholds are sufficiently far away, the coupled-channel effect is small, and bound states are not found. However, when the thresholds move closer, the coupled-channel strength is increased, and bound states may appear for a subset of quantum numbers. Hence, threshold vicinity is a required but not sufficient condition to bind a five-quark state. Under these conditions, the channels with high spin $J^P = 5/2^-$ are the only ones that may lodge a compact five-quark state for all isospins as it is shown in Fig. 6. The reason stems on the reverse of the ordering of the thresh-olds, being the lowest threshold $(nnn)(Q\bar{n})$ the one with the more attractive interaction. In the other cases, the break

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[1]$M_{D^*} + M_{\bar{D}^*} = 4014$ MeV $\geq M_{J/\psi} + M_\omega = 3879$ MeV.

[2]We thank J. M. Richard for this simple and nice argument that does not make any assumption on the shape of the interaction, linear or not, although it assumes a quark-diquark ansatz.
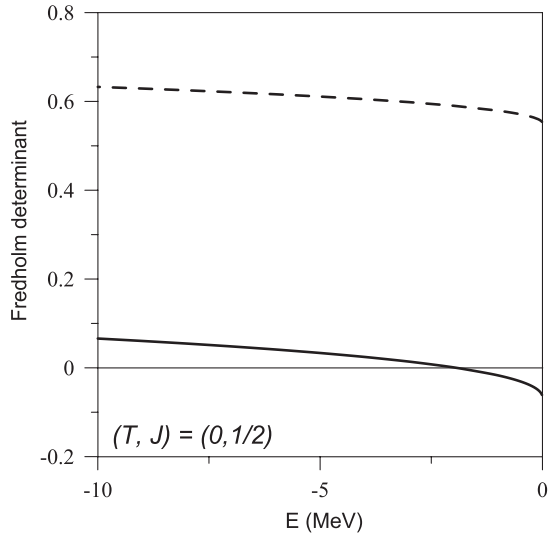
FIG. 7. Fredholm determinant of the $(T)J^P = (0)1/2^-$ channel considering only the ND contributions of Table I (solid line) and considering all contributions (dashed line). E indicates the energy below the corresponding lowest threshold, ND for the solid line and $\Sigma_c \pi$ for the dashed one, as can be seen in Fig. 5.

apart threshold $(nnQ)(n\bar{n})$, weakly interacting, is the lowest one destroying the possibility for any resonance. This is illustrated in Fig. 7 where we show the Fredholm determinant for the $(T)J^P = (0)1/2^-$ channel. When one only considers ND channels (solid line) the interaction is attractive and it has a bound state. However, when the lowest break apart threshold is considered (dashed line) the bound state does not appear any more.

As already advertised in the first part of our discussion, of particular interest is the $(T)J^P = (2)5/2^-$ state that survives the consideration of the break apart thresholds. It may correspond to the $\Theta_c(3250)$ pentaquark found by the QCD sum rule analysis of Ref. [11] when studying the unexplained structure with a mass of 3250 MeV/c^2 in the $\Sigma_c^{++}\pi^-\pi^-$ invariant mass reported recently by the *BABAR* collaboration [40]. Such state could be also detected by the propagation of D mesons in nuclear matter as an S wave ΔD^* system and it thus constitutes a challenge for the PANDA collaboration.

IV. SUMMARY

Summarizing, we have studied higher order Fock space components on the charm baryon spectrum. For this purpose we have used two different approaches. In a first step we did a coupled-channel calculation of the ND system looking for molecular states. In a second step we allowed for the coupling to a heavy baryon—light meson system looking for compact exotic five-quark states. Both calculations have been done within the framework of a full-fledged chiral constituent quark model. This model, tuned in the description of the baryon and meson spectra as well as the NN interaction, provides parameter-free predictions for charm $+1$ molecular or compact two-hadron systems. Unlike the $N\bar{D}$ system, no sharp quark-Pauli effects are found due to the nonexistence of quark-exchange diagrams and thus of the OGE contribution. The importance of the coupled-channel dynamics has been emphasized to connect with the result of other hadronic models. We have found several close to threshold resonances in the ND system that could be traced back to some of the measured charm baryon excited states. If the full dynamics of the five-quark system is considered the number of resonances is reduced and their energies augmented. Of particular interest is the prediction of a $(T)J^P = (2)5/2^-$ baryonic state that survives the consideration of the break apart thresholds. It may correspond to the $\Theta_c(3250)$ pentaquark found by the QCD sum rule analysis of Ref. [11] when studying the unexplained structure with a mass of 3250 MeV/c^2 in the $\Sigma_c^{++}\pi^-\pi^-$ invariant mass reported recently by the *BABAR* collaboration [40].

The advent of new experimental data on charm baryon spectroscopy will shed light about the structure of some of the already observed states that, otherwise, will also help us in understanding the short-range dynamics of many-quarks systems (confinement), either confirming the existence of ND resonances close-to-threshold or not. The first scenario will point to a two-hadron resonance while the second may be a hint for the presence of many-quark states. This objective may be attainable in several current and future collaborations like PANDA, LHCb, ExHIC or J-PARC.

[1] E. Braaten, Phys. Rev. Lett. **111**, 162003 (2013).

[2] S.-K. Choi *et al.* (Belle Collaboration), Phys. Rev. Lett. **91**, 262001 (2003).

[3] M. Gell-Mann, Phys. Lett. **8**, 214 (1964).

[4] G. T. Bodwin, E. Braaten, E. Eichten, S. L. Olsen, T. K. Pedlar, and J. Russ, arXiv:1307.7425, and references therein.

[5] E. Klempt and J.-M. Richard, Rev. Mod. Phys. 82, 1095 (2010).

[6] V. Crede and W. Roberts, Rep. Prog. Phys. 76, 076301 (2013).

[7] X. G. He, X. Q. Li, X. Liu, and X. Q. Zeng, Eur. Phys. J. C 51, 883 (2007).

[8] J. He and X. Liu, Phys. Rev. D 82, 114029 (2010).

[9] Y. Dong, A. Faessler, T. Gutsche, and V. E. Lyubovitskij, Phys. Rev. D 81, 074011 (2010).

[10] Y. Yamaguchi, S. Ohkoda, S. Yasui, and A. Hosaka, Phys. Rev. D 87, 074019 (2013).

[11] R. M. Albuquerque, S. H. Lee, and M. Nielsen, Phys. Rev. D 88, 076001 (2013).

[12] J.-R. Zhang, Phys. Rev. D 89, 096006 (2014).

[13] C. E. Jiménez-Tejero, A. Ramos, and I. Vidaña, Phys. Rev. C 80, 055206 (2009); C. García-Recio, V. K. Magas, T. Mizutani, J. Nieves, A. Ramos, L. L. Salcedo, and L. Tolos, Phys. Rev. D 79, 054004 (2009); O. Romanets, L. Tolos, C. García-Recio, J. Nieves, L. L. Salcedo, and R. G. E. Timmermans, Phys. Rev. D 85, 114032 (2012); M. Bayar, C. W. Xiao, T. Hyodo, A. Doté, M. Oka, and E. Oset, Phys. Rev. C 86, 044004 (2012); P. G. Ortega, D. R. Entem, and F. Fernández, Phys. Lett. B 718, 1381 (2013).

[14] P. González, J. Vijande, and A. Valcarce, Phys. Rev. C 77, 065213 (2008).

[15] U. Wiedner (PANDA Collaboration), Prog. Part. Nucl. Phys. 66, 477 (2011).

[16] J. Kogut, M. Stone, H. W. Wyld, W. R. Gibbs, J. Shigemitsu, S. H. Shenker, and D. K. Sinclair, Phys. Rev. Lett. 50, 393 (1983).

[17] T. Matsui and H. Satz, Phys. Lett. B 178, 416 (1986).

[18] C. B. Dover and S. H. Kahana, Phys. Rev. Lett. 39, 1506 (1977).

[19] A. Yokota, E. Hiyama, and M. Oka (to be published).

[20] A. Khodjamirian, Ch. Klein, Th. Mannel, and Y.-M. Wang, Eur. Phys. J. A 48, 31 (2012).

[21] J. He, Z. Ouyang, X. Liu, and X.-Q. Li, Phys. Rev. D 84, 114010 (2011).

[22] J. Haidenbauer and G. Krein, Phys. Lett. B 687, 314 (2010).

[23] S. Cho et al. (ExHIC Collaboration), Phys. Rev. Lett. 106, 212001 (2011).

[24] S. Cho et al. (ExHIC Collaboration), Phys. Rev. C 84, 064910 (2011).

[25] D. Milanes (LHCb Collaboration), Nucl. Phys. B, Proc. Suppl. 234, 139 (2013).

[26] T. A. Aaltonen et al. (CDF Collaboration), Phys. Rev. D 89, 072014 (2014).

[27] K. Saito, K. Tsushima, and A. W. Thomas, Prog. Part. Nucl. Phys. 58, 1 (2007).

[28] T. F. Caramés, C. E. Fontoura, K. Tsushima, G. Krein, and A. Valcarce (to be published).

[29] A. Valcarce, H. Garcilazo, F. Fernández, and P. González, Rep. Prog. Phys. 68, 965 (2005).

[30] J. Vijande, F. Fernández, and A. Valcarce, J. Phys. G 31, 481 (2005).

[31] T. Fernández-Caramés, A. Valcarce, and J. Vijande, Phys. Rev. Lett. 103, 222001 (2009).

[32] T. F. Caramés and A. Valcarce, Phys. Rev. D 85, 094017 (2012).

[33] A. de Rújula, H. Georgi, and S. L. Glashow, Phys. Rev. D 12, 147 (1975).

[34] G. S. Bali, Phys. Rep. 343, 1 (2001).

[35] A. Valcarce, H. Garcilazo, and J. Vijande, Phys. Rev. C 72, 025206 (2005).

[36] H. Garcilazo, J. Phys. G 13, L63 (1987).

[37] A. Valcarce, A. Faessler, and F. Fernández, Phys. Lett. B 345, 367 (1995).

[38] C. Nakamoto and H. Toki, Prog. Theor. Phys. 99, 1001 (1998).

[39] M. F. M. Lutz and E. E. Kolomeitsev, Nucl. Phys. A755, 29 (2005).

[40] J. P. Lees et al. (BABAR Collaboration), Phys. Rev. D 86, 091102(R) (2012).

[41] S. Capstick and N. Isgur, Phys. Rev. D 34, 2809 (1986); D. Ebert, R. N. Faustov, and V. O. Galkin, Phys. Lett. B 659, 612 (2008); A. Valcarce, H. Garcilazo, and J. Vijande, Eur. Phys. J. A 37, 217 (2008).

[42] A. Valcarce, H. Garcilazo, and J. Vijande, Phys. Lett. B 733, 288 (2014).

[43] J. Vijande and A. Valcarce, Phys. Rev. C 80, 035204 (2009).

[44] T. F. Caramés, A. Valcarce, and J. Vijande, Phys. Lett. B 699, 291 (2011).

[45] R. A. Bertlmann and A. Martin, Nucl. Phys. B168, 111 (1980); S. Nussinov, Z. Phys. C 3, 165 (1979).

[46] T. F. Caramés, A. Valcarce, and J. Vijande, Phys. Lett. B 709, 358 (2012).