PHYSICAL REVIEW D **85,** 094017 (2012)

$N\bar{D}$ system: A challenge for $\bar{\text{P}}$ANDA

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(Received 21 March 2012; published 18 May 2012)

We study the $N\bar{D}$ system by means of a 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 the charm -1 two-hadron systems. The presence of a heavy antiquark makes the interaction rather simple. We have found sharp quark-Pauli effects in some particular channels, generating repulsion, due to lacking degrees of freedom to accommodate the light quarks. Our results point to the existence of an attractive state, the $\Delta\bar{D}^*$ with $(T, S) = (1, 5/2)$, presenting a resonance close to threshold. It would show up in the scattering of $\bar{D}$ mesons on nucleons as a D wave state with quantum numbers $(T)J^P = (1)5/2^-$. This resonance resembles our findings in the $\Delta\Delta$ system, that offered a plausible explanation to the cross section of double-pionic fusion reactions through the so-called Abashian-Booth-Crowe effect. The study of the interaction of D mesons with nucleons is a goal of the $\bar{\text{P}}$ANDA collaboration at the European facility FAIR and, thus, the present theoretical study is a challenge to be tested at the future experiments.

DOI: 10.1103/PhysRevD.85.094017 PACS numbers: 14.40.Lb, 12.39.Pn, 12.40.−y

I. INTRODUCTION

The study of the interaction between charmed mesons and nucleons has become an interesting subject in several contexts [1]. It is particularly interesting for the study of chiral symmetry restoration in a hot and/or dense medium [2]. It will also help in the understanding of the suppression of the J/Ψ production in heavy ion collisions [3]. Besides, it may shed light on the possible existence of exotic nuclei with heavy flavors [4]. Experimentally, it will become possible to analyze the interaction of charmed mesons with nucleons inside nuclear matter with the operation of the FAIR facility at the GSI laboratory in Germany [1]. There are proposals for experiments by the $\bar{\text{P}}$ANDA 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. The $\bar{D}$ mesons then have a chance to interact with the nucleons inside the target material. Since the D mesons are produced in pairs in the antiproton-nucleon annihilation process, the appearance of one of those D mesons can be used as tag to look for such reactions. 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 [5]. Thus, a good knowledge of the interaction of charmed mesons with ordinary hadrons, like nucleons, is a prerequisite.

Before one can infer in a sensitive way changes of the interaction in the medium, 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 work, to make use of a chiral constituent quark model describing the NN interaction [6] as well as the meson spectrum in all flavor sectors [7] to obtain parameter-free predictions that may be testable in the future experiments of the $\bar{\text{P}}$ANDA collaboration. Such a project was already undertaken for the interaction between charmed mesons with reasonable predictions [8], which encourages us in the present challenge.

The paper is organized as follows. In Sec. II we present a description of the quark-model wave function for the baryon-meson system, centering our attention in its short-range behavior looking for quark-Pauli effects. In Sec. III we briefly revise the interacting potential. Section IV deals with the solution of the two-body problem by means of the Fredholm determinant. In Sec. V we present our results. We will discuss in detail the baryon-meson interactions emphasizing those aspects that may produce different results from purely hadronic theories. We will analyze the character of the interaction in the different isospin-spin channels, looking for the attractive ones that may lodge resonances to be measured at $\bar{\text{P}}$ANDA. We will also compare with existing results in the literature. Finally, in Sec. VI we summarize our main conclusions.

II. 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:

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

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 $\bar{D}$ or $\bar{D}^*$ mesons, the antisymmetrization operator is given by

$$\mathcal{A} = \left(1 - \sum_{i=1}^{3} P_{ij}^{\rm LST}\right), \tag{2}$$

where $P_{ij}^{\rm LST}$ exchanges a pair of identical quarks i and j, and j stands for the light quark of the charmed meson. If we assume Gaussian $0s$ wave functions for the quarks inside the hadrons, the normalization of the baryon-meson wave function $\Psi_{B_iM_j}^{\rm LST}(\vec{R})$ of Eq. (1) can be expressed as

$$\mathcal{N}_{B_iM_j}^{\rm LST}(R) = \mathcal{N}_{\rm di}^{L}(R) - C(S,T)\mathcal{N}_{\rm ex}^{L}(R). \tag{3}$$

$\mathcal{N}_{\rm di}^{L}(R)$ and $\mathcal{N}_{\rm ex}^{L}(R)$ stand for the direct and exchange radial normalizations, respectively, whose explicit expressions are

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

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 may impose antisymmetry requirements not present in a hadronic description. Such effects, if any, will be prominent for $L = 0$. Using the asymptotic form of the Bessel functions, $i_{L+1/2}$, we obtain,

$$\mathcal{N}_{\rm di}^{L=0} \xrightarrow[R\to 0]{} 4\pi\left\{1-\frac{R^2}{8}\left(\frac{4}{b^2}+\frac{1}{b_c^2}\right)\right\}\left[1+\frac{1}{6}\left(\frac{R^2}{8}\left(\frac{4}{b^2}+\frac{1}{b_c^2}\right)\right)^2+\ldots\right],$$
$$\mathcal{N}_{\rm ex}^{L=0} \xrightarrow[R\to 0]{} \left\{1-\frac{R^2}{8}\left(\frac{4}{b^2}+\frac{1}{b_c^2}\right)\right\}\left[1+\frac{1}{6}\left(\frac{R^2}{8b_c^2}\right)^2+\ldots\right]. \tag{5}$$

Finally, the S wave normalization kernel, Eq. (3), can be written in the overlapping region as

$$\mathcal{N}_{B_iM_j}^{L=0ST} \xrightarrow[R\to 0]{} \left\{1-\frac{R^2}{8}\left(\frac{4}{b^2}+\frac{1}{b_c^2}\right)\right\}\left\{[1-C(S,T)] + \frac{1}{6}\left(\frac{R^2}{8b_c^2}\right)^2[\gamma^2 - C(S,T)] + \ldots\right\}, \tag{6}$$

where $\gamma = 1 + [(4b_c^2)/b^2]$. Thus, the closer the value of $C(S,T)$ to 1 the larger the suppression of the normalization of the wave function at short distances, generating Pauli repulsion. In particular, if $C(S,T) = 1$ the norm goes to zero for $R \to 0$, which is called Pauli blocking [9]. $C(S,T)$ is a coefficient depending on the total spin (S) and isospin (T) of the B_iM_j two-hadron system and is given by

$$\begin{aligned} C(S,T) = 3\left(\frac{2S_1+1}{4}\right) &\times \sum_{\chi_i=\eta_i=0}^{1}\left\langle\left(\left(\chi_1,\frac{1}{2}\right),S_1;\left(\frac{1}{2},\frac{1}{2}\right),S_2;S,M_S\right|P_{k\ell}^{S}\left|\left(\chi_2,\frac{1}{2}\right),S_1;\left(\frac{1}{2},\frac{1}{2}\right),S_2;S,M_S\right\rangle \\ &\times\left\langle\left(\left(\eta_1,\frac{1}{2}\right),T_1;\left(0,\frac{1}{2}\right),\frac{1}{2};T,M_T\right|P_{k\ell}^{T}\left|\left(\eta_2,\frac{1}{2}\right),T_1;\left(0,\frac{1}{2}\right),\frac{1}{2};T,M_T\right\rangle, \end{aligned} \tag{7}$$

where the subindices k and ℓ in the exchange operator P^{ST} denote a quark of the baryon B_i and a quark of the meson M_j, respectively. χ_i (η_i) stands for the coupled spin (isospin) of two quarks inside the baryon, (S_1, T_1) are the spin and isospin of the baryon N or Δ, and S_2 is the spin of the meson $\bar{D}$ or $\bar{D}^*$.

The values of $C(S,T)$ are given in Table I. Similarly to Pauli blocked channels, corresponding to $C(S,T) = 1$, we will call Pauli suppressed channels those where $C(S,T)$ is close to 1. We can see that there is one system showing Pauli blocking; this is the $\Delta\bar{D}^*$ with $(T,S) = (2, 5/2)$. This is easily understood due to lacking degrees of freedom to accommodate the light quarks present on this configuration. The interaction in this channel will be strongly repulsive [9], as we will discuss in Sec. V. The channel $N\bar{D}^*$ with $(T,J) = (0, 1/2)$, where $C(S,T) = 2/3$ is Pauli suppressed, the norm kernel gets rather small at short distances giving rise to Pauli repulsion at short distances as we will see in Sec. V. We show in Fig. 1 the normalization kernel given by Eq. (3) for $L = 0$ and four different channels: $\Delta\bar{D}^*$ with $(T,J) = (2, 5/2)$, $N\bar{D}^*$ with $(T,J) = (0, 1/2)$, $N\bar{D}$ with $(T,J) = (1, 1/2)$, and $N\bar{D}^*$ with $(T,J) = (0, 3/2)$. In the first three cases $C(S,T)$ is positive, being exactly one in the first case and becoming smaller in the

TABLE I. $C(S,T)$ spin-isospin coefficients defined in Eq. (7).

	B_iM_j	$T=0$	$T=1$	$T=2$
$S=1/2$	$N\bar{D}$	0	1/3	⋯
	$N\bar{D}^*$	2/3	−1/9	⋯
	$\Delta\bar{D}^*$	⋯	1/9	−1/3
$S=3/2$	$N\bar{D}^*$	−1/3	5/9	⋯
	$\Delta\bar{D}$	⋯	−1/6	1/2
	$\Delta\bar{D}^*$	⋯	−1/18	1/6
$S=5/2$	$\Delta\bar{D}^*$	⋯	−1/3	1

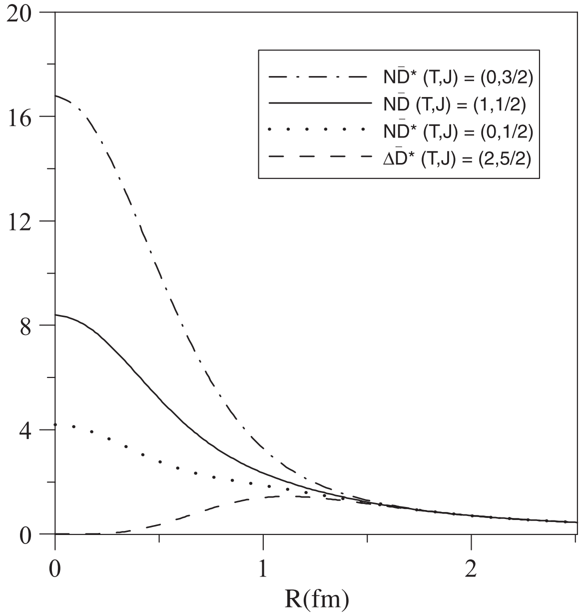

FIG. 1. Norm kernel defined in Eq. (3) for $L = 0$ and four different channels.

others, which makes the norm kernel to augment. In the last case $C(S, T)$ is negative, showing a large norm kernel at short distances and therefore one does not expect any Pauli effect at all.

III. 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 [6]. This model was proposed in the early 1990s in an attempt to obtain a simultaneous description of the nucleon-nucleon interaction and the baryon spectra. It was later on generalized to all flavor sectors [7]. In this model baryons 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 [10]. It reads,

$$V_{\mathrm{OGE}}(\vec{r}_{ij}) = \frac{\alpha_s}{4}\vec{\lambda}_i^c \cdot \vec{\lambda}_j^c \left\{ \frac{1}{r_{ij}} - \frac{1}{4}\left(\frac{1}{2m_i^2} + \frac{1}{2m_j^2} + \frac{2\vec{\sigma}_i \cdot \vec{\sigma}_j}{3m_i m_j}\right) \times \frac{e^{-r_{ij}/r_0}}{r_0^2 r_{ij}} - \frac{3S_{ij}}{4m_q^2 r_{ij}^3} \right\}, \tag{8}$$

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 [7],

$$\alpha_s(\mu) = \frac{\alpha_0}{\ln[(\mu^2 + \mu_0^2)/\gamma_0^2]}, \tag{9}$$

where $\alpha_0 = 2.118$, $\mu_0 = 36.976$ MeV, and $\gamma_0 = 0.113\ \mathrm{fm}^{-1}$. 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,

$$V_\chi(\vec{r}_{ij}) = V_{\mathrm{OSE}}(\vec{r}_{ij}) + V_{\mathrm{OPE}}(\vec{r}_{ij}), \tag{10}$$

where the scalar exchange (OSE) and the pseudoscaolar exchange (OPE) potentials are given by

$$\begin{aligned} V_{\mathrm{OSE}}(\vec{r}_{ij}) &= -\frac{g_{\mathrm{ch}}^2}{4\pi}\frac{\Lambda^2}{\Lambda^2 - m_\sigma^2} m_\sigma \left[Y(m_\sigma r_{ij}) - \frac{\Lambda}{m_\sigma} Y(\Lambda r_{ij}) \right], \\ V_{\mathrm{OPE}}(\vec{r}_{ij}) &= \frac{g_{\mathrm{ch}}^2}{4\pi}\frac{m_\pi^2}{12 m_i m_j}\frac{\Lambda^2}{\Lambda^2 - m_\pi^2} m_\pi \left\{ \left[Y(m_\pi r_{ij}) - \frac{\Lambda^3}{m_\pi^3} Y(\Lambda r_{ij}) \right] \vec{\sigma}_i \cdot \vec{\sigma}_j \right. \\ &\quad \left. + \left[H(m_\pi r_{ij}) - \frac{\Lambda^3}{m_\pi^3} H(\Lambda r_{ij}) \right] S_{ij} \right\} (\vec{\tau}_i \cdot \vec{\tau}_j). \end{aligned} \tag{11}$$

$g_{\mathrm{ch}}^2/4\pi$ 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 [11],

$$V_{\mathrm{CON}}(\vec{r}_{ij}) = \{-a_c(1 - e^{-\mu_c r_{ij}})\}(\vec{\lambda^c}_i \cdot \vec{\lambda^c}_j). \tag{12}$$

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

$$V_{q_i q_j}(\vec{r}_{ij}) = \begin{cases} [q_i q_j = nn] \Rightarrow V_{\mathrm{CON}}(\vec{r}_{ij}) + V_{\mathrm{OGE}}(\vec{r}_{ij}) + V_\chi(\vec{r}_{ij}) \\ [q_i q_j = cn] \Rightarrow V_{\mathrm{CON}}(\vec{r}_{ij}) + V_{\mathrm{OGE}}(\vec{r}_{ij}) \end{cases}, \tag{13}$$

TABLE II. Quark-model parameters.

$m_{u,d}$(MeV)	313
m_c(MeV)	1752
b(fm)	0.518
b_c(fm)	0.6
$\hat{r}_0$ (MeV fm)	28.170
a_c(MeV fm^{-1})	109.7
$g_{\rm ch}^2/(4\pi)$	0.54
m_σ(fm^{-1})	3.42
m_π(fm^{-1})	0.70
Λ(fm^{-1})	4.2
a_c (MeV)	230
μ_c (fm^{-1})	0.70

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. For the sake of completeness we compile the parameters of the model in Table II. The model guarantees a nice description of the baryon (N and Δ) [12] and the meson ($\bar{D}$ and $\bar{D}^*$) spectra [7]. Let us also note that the parameters of the model have been tuned in the meson and baryon spectra and the NN interaction, and therefore the present calculation is free of parameters.

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

$$V_{B_nM_m(\mathrm{LST})\to B_kM_l(L'S'T)}(R) = \xi_{\mathrm{LST}}^{L'S'T}(R) - \xi_{\mathrm{LST}}^{L'S'T}(\infty), \quad (14)$$

where

$$\xi_{\mathrm{LST}}^{L'S'T}(R) = \frac{\langle \Psi_{B_kM_l}^{L'S'T}(\vec{R}) \mid \sum_{i<j=1}^{5} V_{q_iq_j}(\vec{r}_{ij}) \mid \Psi_{B_nM_m}^{\mathrm{LST}}(\vec{R})\rangle}{\sqrt{\langle \Psi_{B_kM_l}^{L'S'T}(\vec{R}) \mid \Psi_{B_kM_l}^{L'S'T}(\vec{R})\rangle}\sqrt{\langle \Psi_{B_nM_m}^{\mathrm{LST}}(\vec{R}) \mid \Psi_{B_nM_m}^{\mathrm{LST}}(\vec{R})\rangle}}. \quad (15)$$

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 $\Psi_{B_nM_m}^{\mathrm{LST}}(\vec{R})$ for the baryon-meson system has been discussed in detail in Sec. II. We show in Fig. 2 the different diagrams contributing to the baryon-meson interaction. The contributions on the first line are pure hadronic interactions, while the diagrams in the second and third lines contain quark exchanges and therefore will not be present in a hadronic description. In Sec. V we will analyze in detail the different diagrams to make clear the contributions arising from pure quark effects.

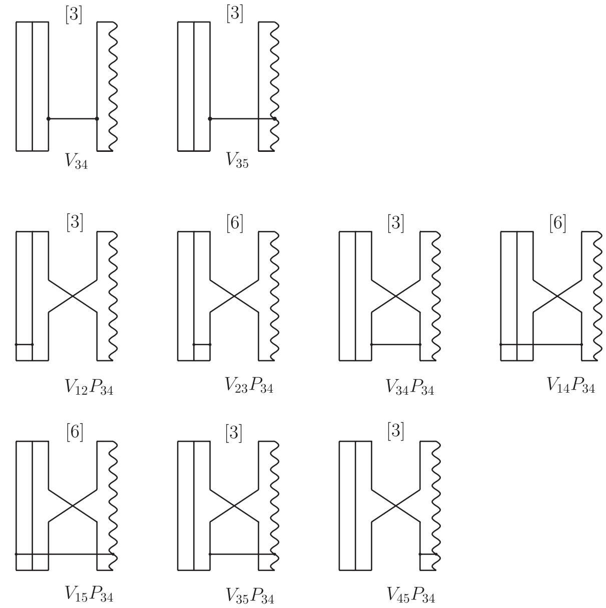

FIG. 2. Different diagrams contributing to the baryon-meson interaction. The vertical, solid lines represent a light quark, u or d, while the wavy line represents the charm antiquark. The horizontal line denotes the interacting quarks. The number between square brackets stands for the number of diagrams topologically equivalent.

IV. INTEGRAL EQUATIONS FOR THE TWO-BODY SYSTEMS

To study the possible existence of exotic states made of a light baryon, N or Δ, and a charmed meson, $\bar{D}$ or $\bar{D}^*$, we have solved the Lippmann-Schwinger equation for negative energies using the Fredholm determinant. 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 [8]. We consider a baryon-meson system B_iM_j ($B_i = N$ or Δ, and $M_j = \bar{D}$ or $\bar{D}^*$) 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_iM_j D wave. Moreover, the baryon-meson system can couple to other baryon-meson states. We show in Table III the coupled channels in the isospin-spin (T, J) basis. Thus, if we denote the different baryon-meson systems as channel A_i, the Lippmann-Schwinger equation for the baryon-meson scattering becomes

TABLE III. Interacting baryon-meson channels in the isospin-spin (T, J) basis.

	$T=0$	$T=1$	$T=2$
$J=1/2$	$N\bar{D}-N\bar{D}^*$	$N\bar{D}-N\bar{D}^*-\Delta\bar{D}^*$	$\Delta\bar{D}^*$
$J=3/2$	$N\bar{D}^*$	$N\bar{D}^*-\Delta\bar{D}-\Delta\bar{D}^*$	$\Delta\bar{D}-\Delta\bar{D}^*$
$J=5/2$		$\Delta\bar{D}^*$	$\Delta\bar{D}^*$

$$t_{\alpha\beta;TJ}^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}(p_\alpha, p_\beta; E) = V_{\alpha\beta;TJ}^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}(p_\alpha, p_\beta) + \sum_{\gamma=A_1,A_2,\cdots}\sum_{\ell_\gamma=0,2}\int_0^\infty p_\gamma^2 dp_\gamma V_{\alpha\gamma;TJ}^{\ell_\alpha s_\alpha,\ell_\gamma s_\gamma}(p_\alpha, p_\gamma) \times G_\gamma(E; p_\gamma) t_{\gamma\beta;TJ}^{\ell_\gamma s_\gamma,\ell_\beta s_\beta}(p_\gamma, p_\beta; E),\ \alpha, \beta = A_1, A_2, \cdots, \tag{16}$$

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, 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

$$G_\gamma(E; p_\gamma) = \frac{2\mu_\gamma}{k_\gamma^2 - p_\gamma^2 + i\epsilon}, \tag{17}$$

with

$$E = \frac{k_\gamma^2}{2\mu_\gamma}, \tag{18}$$

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

$$p_\gamma = d\frac{1 + x_\gamma}{1 - x_\gamma}, \tag{19}$$

where d is a scale parameter, and the same for p_α and p_β, we can write Eq. (16) as

$$t_{\alpha\beta;TJ}^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}(x_\alpha, x_\beta; E) = V_{\alpha\beta;TJ}^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}(x_\alpha, x_\beta) + \sum_{\gamma=A_1,A_2,\cdots}\sum_{\ell_\gamma=0,2}\int_{-1}^1 d^2\left(\frac{1 + x_\gamma}{1 - x_\gamma}\right)^2 \frac{2d}{(1 - x_\gamma)^2} dx_\gamma \times V_{\alpha\gamma;TJ}^{\ell_\alpha s_\alpha,\ell_\gamma s_\gamma}(x_\alpha, x_\gamma) G_\gamma(E; p_\gamma) t_{\gamma\beta;TJ}^{\ell_\gamma s_\gamma,\ell_\beta s_\beta}(x_\gamma, x_\beta; E). \tag{20}$$

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

$$\sum_{\gamma=A_1,A_2,\cdots}\sum_{\ell_\gamma=0,2}\sum_{m=1}^N M_{\alpha\gamma;TJ}^{n\ell_\alpha s_\alpha,m\ell_\gamma s_\gamma}(E) t_{\gamma\beta;TJ}^{\ell_\gamma s_\gamma,\ell_\beta s_\beta}(x_m, x_k; E) = V_{\alpha\beta;TJ}^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}(x_n, x_k), \tag{21}$$

with

$$M_{\alpha\gamma;TJ}^{n\ell_\alpha s_\alpha,m\ell_\gamma s_\gamma}(E) = \delta_{nm}\delta_{\ell_\alpha\ell_\gamma}\delta_{s_\alpha s_\gamma} - w_m d^2\left(\frac{1 + x_m}{1 - x_m}\right)^2 \frac{2d}{(1 - x_m)^2} \times V_{\alpha\gamma;TJ}^{\ell_\alpha s_\alpha,\ell_\gamma s_\gamma}(x_n, x_m) G_\gamma(E; p_{\gamma m}), \tag{22}$$

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. (19). 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., $|M_{\alpha\gamma;TJ}(E_B)| = 0$. We took the scale parameter d of Eq. (19) as $d = 3\ \text{fm}^{-1}$ and used a Gauss-Legendre quadrature with $N = 20$ points.

V. RESULTS AND DISCUSSION

In Figs. 3 and 4 we show some representative potentials of the baryon-meson system under study. In Fig. 3 we have depicted for two channels the contribution of the several pieces of the interacting potential: OGE, OSE, and OPE. In Fig. 4 we have separated for two different channels the contribution of the different diagrams of Fig. 2 in two groups: direct terms (those shown in the first line that do not contain quark exchanges) and exchange terms (diagrams in the second and third lines of Fig. 2 where quark-exchange contributions appear).

Regarding Fig. 3, there are general trends that have already been noticed for other two-hadron systems and we can briefly summarize. For very long distances ($R >$ 4 fm) the interaction is determined by 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

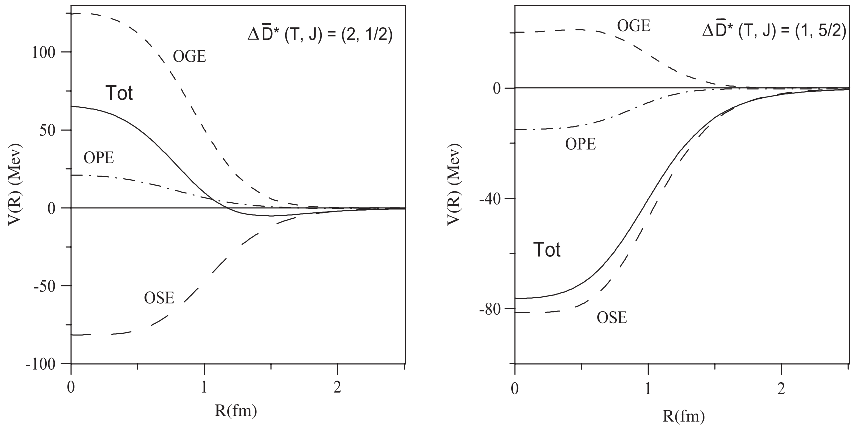

FIG. 3. Contributions to the $\Delta\bar{D}^*$ interaction from the different pieces of the potential detailed in Sec. III in two uncoupled channels: $(T, J) = (2, 1/2)$ (left panel) and $(T, J) = (1, 5/2)$ (right panel). Tot stands for the sum of all contributions (the total potential).

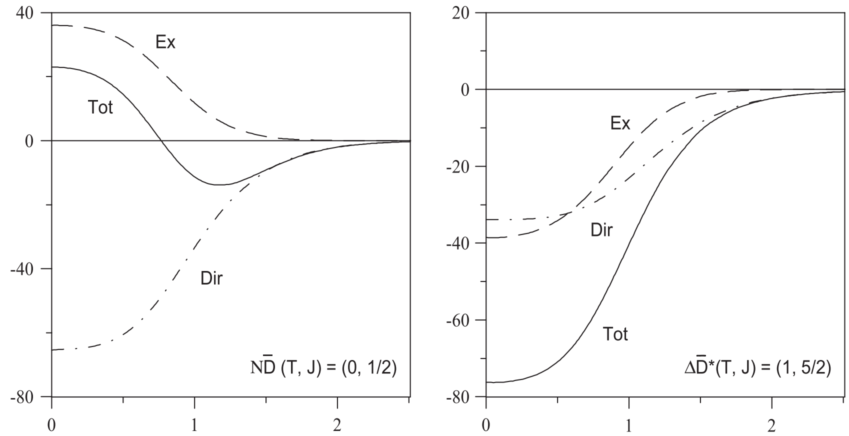

FIG. 4. Direct and exchange contributions to the $N\bar{D}$ $(T, J) = (0, 1/2)$ (left panel) and $\Delta\bar{D}^*$ $(T, J) = (1, 5/2)$ (right panel) potentials. The dash-dotted line represents the contribution of the direct terms (Dir), the dashed line stands for the exchange terms (Ex), and the solid line indicates the total potential (Tot).

repulsive or attractive depending on the balance between the OGE and OPE. In general, one can say that the OGE is mainly repulsive while the OSE is mostly attractive; see the two examples in Fig. 3. Thus, in most cases the character of the OPE at short range determines the final character of the total potential.

It is also interesting to analyze the interaction in terms of the different diagrams plotted in Fig. 2, as it has been done in Fig. 4. The dynamical effect of quark antisymmetrization can be estimated by comparing the total potential with the one arising from the diagrams in the first line of Fig. 2, which are the only ones that do not include quark exchanges. Let us note, however, that Pauli correlations are still present through the norm in the denominator of Eq. (15). To eliminate the whole effect of quark antisymmetrization one should eliminate quark-Pauli correlations from the norm as well. By proceeding in this way one gets a genuine baryonic potential that we call direct potential. The comparison of the total and direct potentials reflects the quark antisymmetrization effect beyond the one-hadron structure. In Fig. 4 we have separated the contribution of the direct and exchange terms for two different partial waves. As can be seen the direct potential is always attractive, due both to the contribution of the scalar exchange interaction and the absence of the one-gluon exchange (contributing only through quark-exchange diagrams). However, the character of the exchange part, containing quark-exchange diagrams that would therefore be absent in a pure hadronic description, depends on the sign of the color-spin-isospin coefficients. The sign of the dominant quark-exchange diagram, $V_{34}P_{34}$, is crucial for determining whether the exchange contributions are attractive or repulsive. Thus, quark-exchange diagrams give repulsion for the $N\bar{D}$ $(T, J) = (0, 1/2)$ channel (see left panel of Fig. 4), but attraction for the $\Delta\bar{D}^*$ $(T, J) = (1, 5/2)$ (see right panel of Fig. 4), determining the final character of the

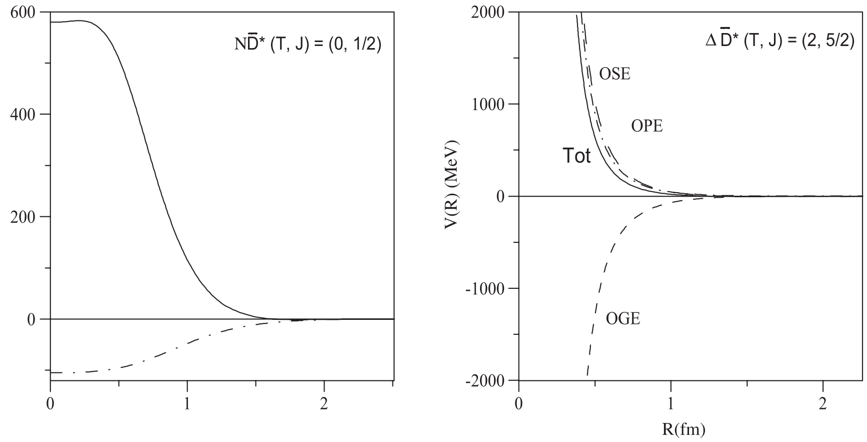

FIG. 5. Interacting potential in two different channels showing Pauli effects: $N\bar{D}^*$ $(T, J) = (0, 1/2)$ (left panel) and $\Delta\bar{D}^*$ $(T, J) = (2, 5/2)$ (right panel). See text for details.

interacting potential. Thus, dynamical quark-exchange effects play a relevant role in the $N\bar{D}$ interaction with observable consequences, as we will discuss below, being responsible for the repulsive character of the $N\bar{D}$ interaction at short distances in some partial waves.

The potential becomes strongly repulsive in those cases where we have Pauli blocked, $C(S, T) = 1$, or Pauli suppressed channels, $C(S, T)$, close to 1. In the left panel of Fig. 5 we have drawn the $N\bar{D}^*$ $(T, J) = (0, 1/2)$ interaction. As seen in Table I the value of $C(S, T) = 2/3$ is close to 1, suppressing the overlapping of the wave function at short distances; see Fig. 1. This gives rise to a strong repulsion at short range. The dash-dotted line in this figure represents the direct potential, the one obtained at hadronic level neither considering exchange diagrams in the norm kernel nor in the interacting potential. As can be seen, the effect of quark exchanges is rather important in what we have called Pauli suppressed partial waves. Hadronic interactions would therefore not be able to account for the consequences of Pauli effects as it was already noticed in a comparative study of the $N\Delta$ interaction by means of hadronic or quark-based models [13]. In the right panel of Fig. 5 we have drawn the $\Delta\bar{D}^*$ $(T, J) = (2, 5/2)$ interaction. In this case $C(S, T)$ is exactly 1, forbidding the overlapping of the two hadrons for $R = 0$. All contributions are very strong at short distances due to the behavior of the norm kernel and the total interaction becomes strongly repulsive. These forbidden states would have to be eliminated by hand in the resonating group method treatment of the two-hadron systems [14]. The existence of such a strong repulsion has also been observed in the $N\Delta$ system and may be concluded from the $N\Delta$ phase shift behavior derived from the πd elastic scattering data [15].

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, $\bar{D}$ or $\bar{D}^*$, as explained in Sec. IV. The existence of bound states or resonances will generate exotic states with charm -1 that could be identified in future experiments of the $\bar{\text{P}}$ANDA collaboration at the FAIR facility [1]. In Table IV we summarize the character of the interaction in the different (T, J) channels. It can be observed how all (T, J) channels containing Pauli blocked or Pauli suppressed states are repulsive: $(2, 5/2)$, $(0, 1/2)$, and $(1, 1/2)$. Thus, the Pauli principle at the level of quarks plays an important role in the dynamics of the $N\bar{D}$ system.

In Figs. 6–8 we have plotted the Fredholm determinant and some of the potentials of some representative channels. Figure 6 shows the weakly attractive $(T, J) = (0, 3/2)$ and the weakly repulsive $(T, J) = (1, 1/2)$ Fredholm determinants. In Fig. 7 we have plotted all potentials contributing to the $(T, J) = (1, 1/2)$ channel that, as seen in the right panel of Fig. 6, is weakly repulsive. The interaction in the lightest two-hadron systems contributing to this channel, $N\bar{D}$ and $N\bar{D}^*$, has a repulsive character and there is a weak coupling to the attractive but heavier two-hadron system $\Delta\bar{D}^*$, giving an overall repulsive channel. Figure 8 shows the Fredholm determinant for the two attractive channels of the $N\bar{D}$ system: $(T, J) = (2, 3/2)$ and $(T, J) = (1, 5/2)$. The $(T, J) = (1, 5/2)$ is the most attractive one, showing a bound state with a binding energy of 3.87 MeV. This channel corresponds to a unique physical system, $\Delta\bar{D}^*$. This situation is rather similar to the one we found in the $\Delta\Delta$ system [16], predicting an S wave resonance with maximum spin. Experimental evidence for such a reso-

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

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

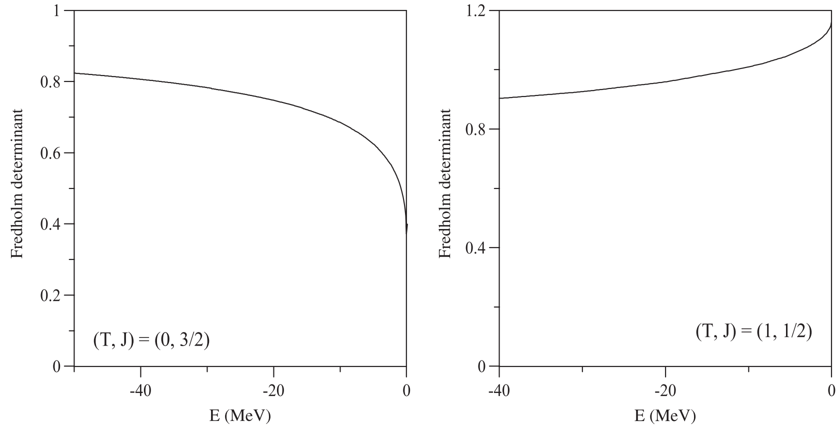

FIG. 6. $(T, J) = (0, 3/2)$ (left panel) and $(T, J) = (1, 1/2)$ (right panel) Fredholm determinant.

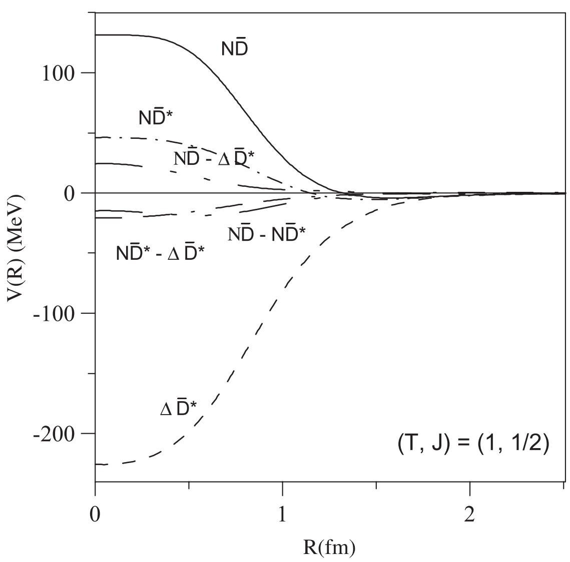

FIG. 7. Baryon-meson potentials contributing to the $(T, J) = (1, 1/2)$ channel.

nance was already reported in the NN scattering data, where the resonance appeared in the $3D_3$ partial wave, thus as a D wave in the NN system [17]. This prediction has been recently used as a possible explanation of the measured cross section of the double-pionic fusion of nuclear systems through the so-called Abashian-Booth-Crowe (ABC) effect [18]. The formation of an intermediate $\Delta\Delta$ resonance with the isospin, spin, parity, and mass, found in Ref. [16] [$(T)J^P = (0)3^+$ and $M = 2.37$ GeV], allowed to describe the cross section of the double-pionic fusion reaction $pn \rightarrow d\pi^0\pi^0$. In the present case, the bound state in the $(T, J) = (1, 5/2)$ $\Delta\bar{D}^*$ channel would appear in the scattering of $\bar{D}$ mesons on nucleons as a D wave resonance, which could in principle be measured at $\bar{\text{P}}$ANDA. Such a state will have quantum numbers $(T)J^P = (1)5/2^-$. Just to make sure that our conclusion does not depend on our choice of the parameter of the charm quark wave function that does not come out from the solution of the meson spectra, we have calculated the binding energy as a function of b_c. The result appears in Fig. 9, showing a

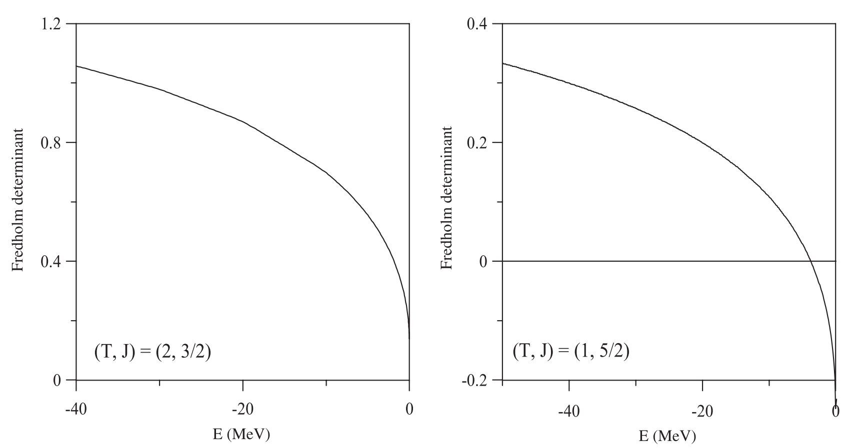

FIG. 8. $(T, J) = (2, 3/2)$ (left panel) and $(T, J) = (1, 5/2)$ (right panel) Fredholm determinant.

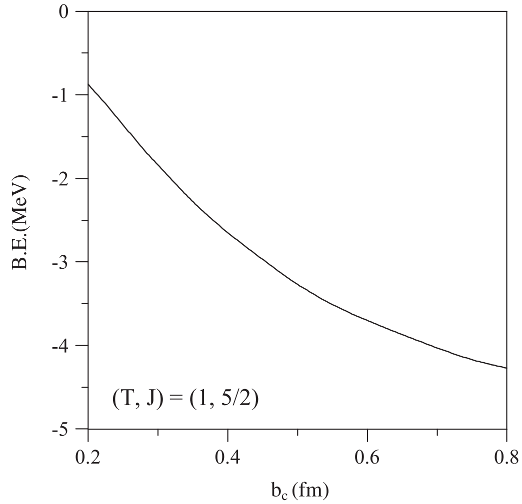

FIG. 9. Binding energy, in MeV, of the $(T, J) = (1, 5/2)$ channel as a function of the harmonic oscillator parameter used for the charm quark.

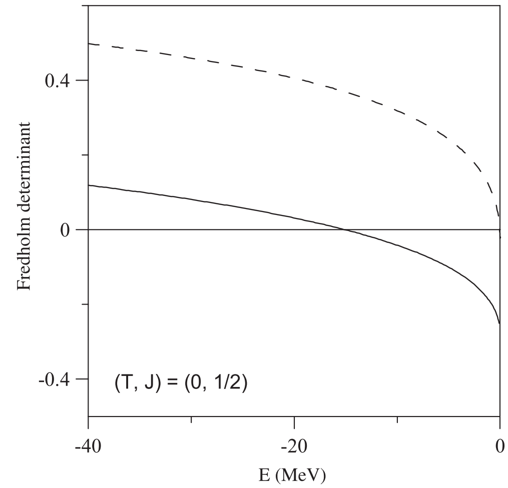

FIG. 10. $(T, J) = (0, 1/2)$ Fredholm determinant obtained with the direct potentials. The dashed line stands for the single channel problem $N\bar{D}$ and the solid line indicates the coupled-channel problem using also the $N\bar{D}^*$ state.

bound system for any reasonable choice of such a parameter, giving us confidence to our prediction.

In Refs. [19,20] the $N\bar{D}$ system has been analyzed by means of a hadronic model using Lagrangians satisfying heavy quark symmetry and chiral symmetry. They arrive to the conclusion that the $(T)J^P = (0)1/2^-$ channel is the most attractive one presenting a bound state of around 1.4 MeV. To emphasize the importance of quark-exchange effects we have repeated for this channel the calculations explained in Sec. IV using only the direct potentials (those without quark-exchange effects). These potentials, that would correspond to a purely baryonic interaction, are drawn as a dash-dotted line in the left panels of Fig. 4 [the $(T, J) = (0, 1/2)$ $N\bar{D}$ interaction] and Fig. 5 [the $(T, J) = (0, 1/2)$ $N\bar{D}^*$ interaction]. In both cases the importance of quark-exchange effects can be seen. These interactions would be the hadronic potentials of an effective theory without quark degrees of freedom. The results obtained neglecting the contribution of quark-exchange effects are shown in Fig. 10. As we can see in both cases, single-channel or coupled-channel calculation, a bound state appears in agreement with the conclusions of Refs. [19,20]. Once again, this comparison makes evident the great importance that quark-exchange effects may have in the system under study and it also represents a sharp example of a system where the quark-exchange dynamics may have observable consequences. A future effort in the study of the $N\bar{D}$ system will provide us with evidence to learn about the importance of quark-exchange dynamics.

Let us finally mention that there are recent estimations in the literature [21] about the $N\bar{D}$ cross section based on hybrid models considering meson exchanges supplemented with a short-range quark-gluon dynamics. The important conclusion of that work is that the predicted $N\bar{D}$ cross sections are of the same order of magnitude as those for the NK system, but with average values of 20 mb, roughly a factor of 2 larger than for the latter system. Thus, the study of the $N\bar{D}$ interaction could be a plausible challenge for the $\bar{P}$ANDA collaboration.

VI. SUMMARY

Summarizing, we have studied the $N\bar{D}$ system at low energies by means of a chiral constituent quark model that describes the baryon and meson spectra as well as the NN interaction. Because of the presence of a heavy antiquark the interaction becomes rather simple and parameter-free predictions can be obtained for the $N\bar{D}$ system. We have analyzed in detail the interaction in the different isospin-spin channels, emphasizing characteristic features consequence of the contribution of quark-exchange dynamics. Quark-Pauli effects generate a strong repulsion in some particular channels due to lacking degrees of freedom to accommodate the light quarks. Such effects have observable consequences generating repulsion in channels that otherwise would be attractive in a hadronic description. We have traced back our results to the previous analysis of the $N\Delta$ and $\Delta\Delta$ systems with peculiar predictions supported by the experimental data. We have found only one bound state in the $\Delta\bar{D}^*$ $(T, S) = (1, 5/2)$ system. This state, $(T)J^P = (1)5/2^-$, will show up in the scattering of $\bar{D}$ mesons on nucleons as a D wave resonance. Such a resonance resembles our findings in the $\Delta\Delta$ system that offered a plausible explanation to the cross section of double-pionic fusion reactions through the so-called ABC effect. The existence of this state is a sharp prediction of quark-exchange dynamics because in a hadronic model the attraction appears in different channels. Finally, it is important to emphasize that theoretical estimations indicate that the $N\bar{D}$ cross section may be attainable in the future facility FAIR and therefore the predicted resonance may be a challenge for the study of the $\bar{\rm P}$ANDA collaboration. This objective may constitute a helpful tool in discriminating among the different scenarios used to describe the dynamics of heavy hadron systems.

ACKNOWLEDGMENTS

This work has been partially funded by the Spanish Ministerio de Educación y Ciencia and EU FEDER under Contract No. FPA2010-21750, and by the Spanish Consolider-Ingenio 2010 Program CPAN (CSD2007-00042).

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