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Hadronic molecules with a $\bar{D}$ meson in a medium

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We study the effect of a hot and dense medium on the binding energy of hadronic molecules with open-charm mesons. We focus on a recent chiral quark-model-based prediction of a molecular state in the $N\bar{D}$ system. We analyze how the two-body thresholds and the hadron-hadron interactions are modified when quark and meson masses and quark-meson couplings change in a function of the temperature and baryon density according to predictions of the Nambu–Jona-Lasinio model. We find that in some cases the molecular binding is enhanced in medium as compared to their free-space binding. We discuss the consequences of our findings for the search for exotic hadrons in high-energy heavy-ion collisions as well as in the forthcoming facilities FAIR or J-PARC.

I. INTRODUCTION

Recent developments in hadron physics have been motivated by the observation of exotic hadrons [1–4]. Most of them lie near open heavy-flavor thresholds implying that they may form the so-called hadronic molecules, that are colorless hadronic clusters loosely bound by a relatively weak residual interaction. This could be a possible situation of what is expected to occur, in general, for multiquark systems. In this respect, the recently discovered five-quark baryonic resonances by the LHCb Collaboration at the Large Hadron Collider (LHC) at CERN, $P_c(4380)^+$ and $P_c(4450)^+$ [5], have been interpreted as simple baryon-meson bound states. Also, some of the exotic mesonic states discovered in the hidden-charm and hidden-beauty sectors [6–9] might be well understood as meson-meson resonances. However, a comprehensive theoretical explanation of the nature of these exotic states is still missing [3,10,11].

The interest in structures containing hadrons with heavy flavors has also been recently reinvigorated by several studies about the existence of nuclear bound states with heavy mesons [12,13] and baryons [14–17], the latter ones already predicted soon after the discovery of baryons

possessing net charm [18–20]. In fact, there are theoretical estimations of the production cross sections as well as experimental requirements for producing charmed hypernuclei by means of charm exchange reactions on nuclei [21]. Last but not least, it is also worth mentioning the suggestion of possible bound states of charmonium in nuclei due to multiple gluon exchange [22,23]. Such a possibility has been recently revisited by means of effective Lagrangians [24,25] and effective Gaussian potentials [26,27], stimulated by the lattice QCD suggestion of a weakly attractive interaction between charmonium and nucleons [28].

On the experimental side, there are exciting perspectives at extant and forthcoming facilities. At the LHC, all four collaborations, ALICE, ATLAS, CMS and LHCb, are engaged in searches for exotic hadrons. Particularly interesting is the possibility of the production of exotic hadrons in the hot and dense environment created in a high-energy heavy-ion collisions. In such an environment, heavy quarks are abundantly produced and they can pick up light-flavor quarks and antiquarks during the evolution of the medium and form multiquark states through a coalescence mechanism [29,30]. The production of bound states of D mesons with nucleons and nuclei can be achieved in different laboratories worldwide. There are planned experiments by the $\bar{P}$ANDA Collaboration to produce them by annihilating antiprotons on nuclei at the Facility for Antiproton Ion Research (FAIR) [31,32]. The planned installation of a 50 GeV high-intensity proton beam at Japan Proton Accelerator Research Complex (J-PARC) [33–35] provides

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an additional opportunity. A Super *B* collider [36] offers similar possibilities.

Thus, the physics of charm in the nuclear medium is becoming a hot topic and is expected to bring further progress in our understanding of the basic theory of the strong interaction, quantum chromodynamics (QCD). First-principles, analytical calculations within QCD of nuclear processes are presently impossible and, consequently, the interpretation of nuclear reaction results will always be afflicted by large uncertainties. The complete lack of experimental information on elementary interactions of charmed hadrons with nucleons in free space imposes additional difficulties in accessing in-medium effects. In order to make progress, one way to proceed, as advocated in Refs. [37–41], is to use models constrained as much as possible by symmetry arguments, analogies with other similar processes, and the use of different degrees of freedom. This is particularly true when it comes to quark-model studies of the production of bound states in a hot and dense medium, as possible in-medium changes of the properties of the light constituent quarks must be taken into account.

The behavior of quarkonia in a hot medium has attracted much interest. It was for instance suggested a long time ago, that in a color-deconfined medium with a temperature above a critical value, T_c, charmonia will melt due to the color Debye screening, and thus serve as a signal for the formation of quark-gluon plasma [42]. For the bottomonium spectrum, a similar phenomenon may occur. Indeed, it was recently reported that bottomonium spectra have received significant modifications when comparing their yield in proton-proton and Pb-Pb collisions at the LHC [43]. The investigation of the nuclear medium effects on charmed hadron bound states is the main objective of this work. For this purpose we will select a charmed hadron molecule whose existence is determined by pure quark effects, the $(I)J^P = (1)5/2^- \ \Delta\bar{D}^*$ molecule [41]. The bound state is determined by an attractive short-distance quark-exchange interaction, a feature due to the Pauli principle at the quark level that cannot be captured by an effective Lagrangian employing low-dimension hadronic operators. Given the prominent role played by quark-exchange effects in free space, we investigate the impact of in-medium changes in the parameters of the chiral quark model used in the evaluation of the binding energy of the molecule. To assess the required in-medium dependence on the constituent quark masses and the coupling constants of the light quarks to the π and σ mesons, we use the Nambu–Jona-Lasinio (NJL) model [44,45].

The paper is organized as follows. We use Sec. II for describing the in-medium dependence of the basic ingredients of the constituent chiral quark model: quark and meson masses as well as quark-meson couplings within a NJL framework. In Sec. III we study the hadron masses in medium by means of the modifications we have derived for

the basic parameters of the quark model used. We discuss in Sec. IV the in-medium hadron-hadron interactions and we will briefly revise the solution of the two-body bound-state problem looking for bound states. We present and discuss our results in Sec. V. Finally, in Sec. VI we summarize our main conclusions.

II. MEDIUM DEPENDENCE OF QUARK AND MESON MASSES AND QUARK-MESON COUPLINGS

Within the perspective of the chiral quark model, changes in the masses of the light hadrons and their mutual interactions at finite temperature (T) and baryon density (ρ_B) are driven by the change of the order parameter of dynamical chiral symmetry breaking, the quark condensate. For sufficiently large values of T and ρ_B, the (absolute value of the) in-medium light quark condensate, $\langle\langle\bar{q}q\rangle\rangle$, becomes very small; in the chiral limit, it can actually vanish. For zero baryon density, lattice QCD simulations [46] at almost physical pion masses ($m_\pi = 161$ MeV) have shown a drastic decrease of $|\langle\langle\bar{q}q\rangle\rangle|$ around a temperature of $T_{pc} = 154 \pm 9$ MeV. For finite baryon densities, the combined T and ρ_B behavior of $\langle\langle\bar{q}q\rangle\rangle$ is presently unknown; the main reason for the lack of this knowledge is due to difficulties of using the Monte Carlo methods of lattice QCD due to the sign problem [47]. On the other hand, for low T and ρ_B, there are model-independent predictions [48–50] for $\langle\langle\bar{q}q\rangle\rangle$:

$$\frac{\langle\langle\bar{q}q\rangle\rangle}{\langle\bar{q}q\rangle} = 1 - \sum_h \frac{\Sigma_h}{f_\pi^2 m_\pi^2}\rho_s^h = 1 - \frac{T^2}{8f_\pi^2} - \frac{1}{3}\frac{\rho_B}{\rho_0}, \quad (1)$$

where $\Sigma_h = m_q \partial m_h/\partial m_q$, ρ_s^h is the scalar density of hadron h in matter, m_q is the current quark mass, $\langle\bar{q}q\rangle$ is the vacuum light quark condensate, f_π is the pion leptonic decay constant, and ρ_0 is the baryon saturation density of nuclear matter.

While changes in the constituent quark mass for low values of T and ρ_B could be directly related to the T and ρ_B dependence of the condensate, the calculation of corresponding changes in the masses of the π and σ mesons and their couplings to the constituent quarks requires a model. In this work we employ the NJL model; in addition to reproducing the result in Eq. (1), its bosonized version with π and σ mesons [51] has the same Yukawa quark-meson couplings as those in the chiral constituent quark model (CCQM) of Ref. [41], and it gives very simple expressions for the masses and couplings (for reviews on this and other applications of the model in different problems in hadron and nuclear physics, see e.g. Refs. [52–55]).

To make the paper self-contained and set the notation, we review the basic features of the NJL model relevant for our purposes here. The results we use are derived from the Lagrangian density:

$$\mathcal{L}_{NJL} = \bar{q}(i\partial \!\!\!/ - m_q)q + G[(\bar{q}q)^2 + (\bar{q}i\gamma_5\boldsymbol{\tau}q)^2], \quad (2)$$

where m_q is the current quark mass; we will work in the isospin symmetric limit, $m_q = m_u = m_d$. At finite T and ρ_B, the meson masses m_σ and m_π and the quark-meson coupling constants $g_{qq\sigma}$ and $g_{qq\pi}$, defined respectively as the poles of the meson propagators and their residues, are obtained from the equations ($M = \sigma, \pi$)

$$1 - 2G\Pi_M(\omega^2 = m_M^2) = 0, \quad g_{qqM}^2 = \left[\frac{\partial \Pi_M(\omega^2)}{\partial \omega^2}\right]^{-1}_{\omega^2 = m_M^2}, \quad (3)$$

where $\Pi_M(\omega^2)$ is the meson polarization function

$$\Pi_M(\omega^2) = 2N_c N_f \int \frac{d^3\boldsymbol{k}}{(2\pi)^3} \frac{1 - [n_q^+(\boldsymbol{k}) + n_q^-(\boldsymbol{k})]}{E_q(\boldsymbol{k})} F_M(\omega^2), \quad (4)$$

with $N_c = 3$ and $N_f = 2$ being the number of colors and flavors, $E_q(\boldsymbol{k}) = (M_q^2 + \boldsymbol{k}^2)^{1/2}$,

$$F_\pi(\omega^2) = \frac{E_q^2(\boldsymbol{k})}{E_q^2(\boldsymbol{k}) - \omega^2/4}, \qquad F_\sigma(\omega^2) = \frac{E_q^2(\boldsymbol{k}) - M_q^2}{E_q^2(\boldsymbol{k}) - \omega^2/4}, \quad (5)$$

and M_q is the constituent quark mass, which is a solution of the gap equation that involves the quark condensate $\langle\!\langle \bar{q}q \rangle\!\rangle$:

$$M_q = m_q + 4G\langle\!\langle \bar{q}q \rangle\!\rangle, \quad (6)$$

$$\langle\!\langle \bar{q}q \rangle\!\rangle = N_c N_f M_q \int \frac{d^3\boldsymbol{k}}{(2\pi)^3} \frac{1 - [n_q^+(\boldsymbol{k}) + n_q^-(\boldsymbol{k})]}{E_q(\boldsymbol{k})}. \quad (7)$$

Here and in Eq. (4), $n_q^\pm(\boldsymbol{k})$ are the quark and antiquark Fermi-Dirac distributions

$$n_q^\pm(\boldsymbol{k}) = \frac{1}{1 + e^{\beta[E_q(\boldsymbol{k}) \mp \mu_B]}}, \quad (8)$$

with $\beta = k_B T$ and μ_B is the quark-baryon chemical potential. The baryon density ρ_B is given in terms of these distributions by

$$\rho_B = \frac{2N_c N_f}{3} \int \frac{d^3\boldsymbol{k}}{(2\pi)^3} [n_q^+(\boldsymbol{k}) - n_q^-(\boldsymbol{k})]. \quad (9)$$

The integral in Eq. (4) is to be understood as a principal-value integral when $\omega^2 > 4M_Q^2$. The temperature-independent part of the integrals in Eqs. (4) and (6) are ultraviolet divergent and need regularization. Since the model is nonrenormalizable, the regularization scheme is

part of the model; here we use a three-dimensional cutoff scheme parametrized by a cutoff Λ.

Vacuum quantities, which are used to fit the parameters of the model, are obtained from the above equations by setting the Fermi-Dirac distributions to zero. It is important to note that the couplings $g_{qq\sigma}$ and $g_{qq\pi}$ are not bare couplings; they incorporate the effects of dynamical chiral symmetry breaking (DCSB) and as such are different from each other. In the chiral quark model, such effects arise from corrections to the bare quark-meson vertices and are parametrized by phenomenological form factors. At high temperatures and densities, when chiral symmetry is restored, the couplings $g_{qq\sigma}$ and $g_{qq\pi}$ become equal to each other, as we discuss in the next section.

The free parameters of the NJL model are the current quark mass m_q, the coupling G and the cutoff Λ. They are fixed by fitting the vacuum values for the quark condensate $\langle \bar{u}u \rangle = \langle \bar{d}d \rangle$, the pion decay constant f_π and the pion mass m_π. Taking [53] $m_q = m_u = m_d = 5$ MeV, $G\Lambda^2 = 2.14$ and $\Lambda = 653$ MeV, one obtains $\langle \bar{u}u \rangle = \langle \bar{d}d \rangle = -(252 \text{ MeV})^3$, $f_\pi = 94$ MeV and $m_\pi = 142$ MeV. With such parameters, one obtains for the constituent quark mass $M_q = M_u = M_d = 328$ MeV and for the σ mass $m_\sigma = 663$ MeV.

In Fig. 1 we present the results for the masses of the constituent quarks M_q, the π and σ masses, $m_M, M = \pi, \sigma$, and quark-meson couplings g_{Mqq} as a function of temperature for different values of the chemical potential μ, i.e., (a) $\mu = 0$, (b) $\mu = 0.1$ GeV, (c) $\mu = 0.15$ GeV, and (d) $\mu = 0.2$ GeV. As expected, the constituent quark mass M_q and the σ mass $m_\sigma(\sim 2M_q)$ drop significantly at sufficiently high temperatures and densities, while the couplings become degenerate. We also note that the vacuum values of the masses of the constituent quarks and of the mesons differ from those in the CCQM—see Table II of Ref. [41]—by less than 5%. One could readjust the parameters of the NJL model to obtain even closer results but, as we are mainly interested in the medium dependence of these quantities, we simply use the ratios of the medium to vacuum values of the masses and couplings for calculating the baryon and meson masses and their interactions.

We close this section by reflecting on the limitations and advantages of the present calculation. Initially, it should be clear that both the NJL and chiral quark models are supposed to describe the chiral aspects of QCD in vacuum and at low temperature and baryon chemical potential. At sufficiently high temperature and chemical potential, chiral symmetry is restored and there is no quark condensate, constituent quarks, bound sigma mesons, effective quark-meson couplings, etc. The models break down and are not adequate to describe QCD in such regimes. The discrete poles in the σ and π meson correlation functions melt into a continuum above the (pseudo)critical temperature, describing correlations of essentially massless quark-antiquark
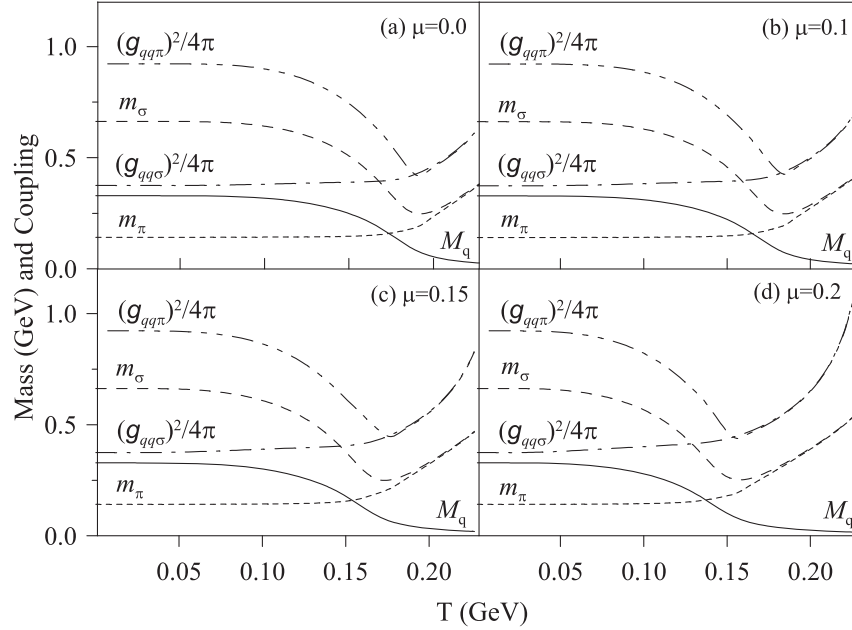
FIG. 1.   Temperature dependence of the masses of the constituent quark M_q (solid line), pion m_π (dotted line), and the sigma m_σ (dashed line), and the quark-meson couplings $g^2_{qq\sigma}/4\pi$ (dash-dotted line) and $g^2_{qq\pi}/4\pi$ (dash-double-dotted line) for different values of chemical potential μ, in GeV.

pairs with thermal masses that grow with temperature. The NJL model is not expected to describe the precise QCD behavior of these thermal masses and it would be desirable to employ a model that interpolates between the low- and high-energy regimes of QCD, like those that incorporate confinement, dynamical chiral symmetry breaking and asymptotic freedom, and describe baryon bound states [56,57]. Such models would allow one to calculate all meson-baryon properties, in vacuum and in medium, within a single framework. Medium effects can be incorporated in the spirit of the quark-meson coupling model, in which meson mean fields couple directly to current quarks in the hadron; for a review, see Ref. [58]. However, the study of hadron-hadron interactions, taking into account, in particular, quark-exchange effects, is difficult due to the use of an underlying soliton or bag model. In the context of a CCQM, an earlier investigation on nucleon and nuclear matter properties, including quark exchange effects, is the one of Ref. [59], but no application of the model to in-medium charmed hadrons is available. In view of this, and given the close relationship between the CCQM and the NJL model, we believe that our calculation captures the basic physics of chiral symmetry on quark exchange in the interaction of $\bar{D}$ mesons with nucleons.

III. HADRON MASSES IN MEDIUM

In the CCQM of Ref. [41], baryons are described as clusters of three interacting massive (constituent) quarks, with their mass coming from dynamical chiral symmetry

breaking in QCD. Short-distance perturbative QCD effects are taken into account through the one-gluon exchange (OGE). In addition to the masses for the constituent quarks, DCSB implies the presence of (pseudo-)Goldstone bosons; their effects are taken into account by introducing them as explicit degrees of freedom via π and σ fields. These fields introduce long-range interactions between the light u and d constituent quarks. Quark confinement is incorporated via an effective potential that contains string-breaking effects. The charm and light quarks interact only via one-gluon exchange and, of course, are subject to the same confining potential. For a review on the model as well as the technical details and methods to solve the two- and three-quark problems, see Refs. [60,61].

Given the temperature and chemical potential dependence of quark and meson masses and quark-meson couplings derived in the previous section, one can calculate the masses of the hadrons of interest: $\bar{D}$ and $\bar{D}^*$ mesons and N and Δ baryons. We show in Fig. 2(a) the variation of the $\bar{D}$ and $\bar{D}^*$ meson masses as a function of the temperature for the different chemical potentials. There is almost no variation of the mass for any chemical potential for a temperature below 0.1 GeV. For temperatures above this value, the masses of the pseudoscalar and vector mesons change in a rather similar manner, which makes manifest that the variation of the mass is a spin-independent effect. Being systems made of a light and a heavy quark, only confinement and the one-gluon exchange contribute to the mass of the $\bar{D}$ and $\bar{D}^*$ mesons. The dominant contribution to the variation of the masses with temperature comes from
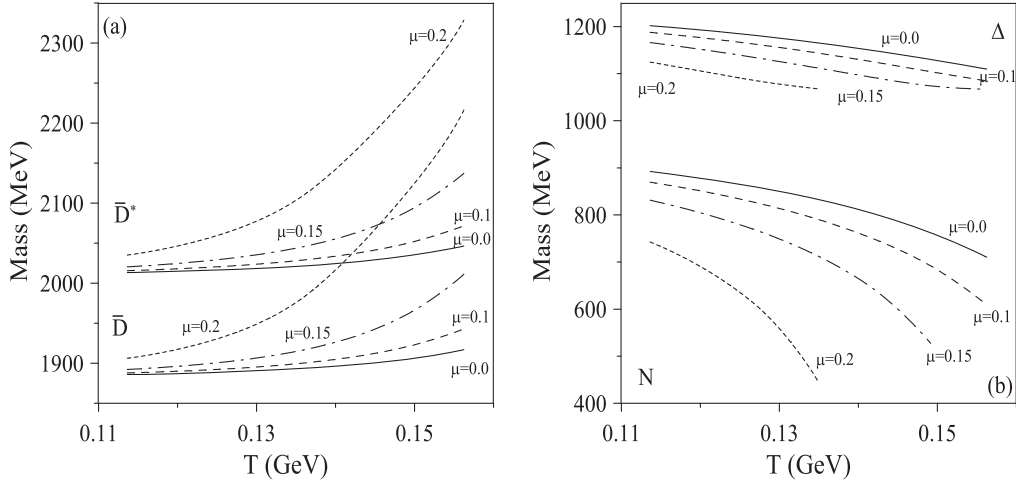
FIG. 2. (a) Masses of the $\bar{D}$ and $\bar{D}^*$ mesons as a function of the temperature for the different chemical potentials. (b) Same as (a) for the N and Δ baryons.

the kinetic energy due to the reduction of the mass of the light constituent quark. The mass of the charm quark is not modified by temperature, and the reduced mass of the heavy-light system approaches the mass of the light quark. It is therefore the kinetic energy that is mainly responsible for the change in the mass of the mesons. The modifications in the spin-dependent part, which are responsible for the mass difference between the $\bar{D}$ and the $\bar{D}^*$ mesons, are minimized due to the presence of the heavy-quark mass in the denominator of the one-gluon exchange through its $1/(m_{q_i} m_{q_j})$ dependence. Similar results have been obtained in the literature very recently for the variation of the $\bar{D}$ meson mass in the nuclear medium using QCD sum rules [62].

In Fig. 2(b) we depict the variation of the masses of the N and Δ baryons as a function of the temperature for the different chemical potentials. As we can see there are important differences as compared to the $\bar{D}$ meson case, due to the presence of three light quarks. Being a more involved system, it can be easily concluded that the diminishing of the mass of the nucleon is mainly due to the spin-dependent part of the one-gluon exchange, which also generates the decreasing of the mass of the Δ. While the $(\vec{\sigma} \cdot \vec{\sigma})(\vec{\tau} \cdot \vec{\tau})$ structure of the chiral pseudoscalar interaction gives attraction for symmetric spin-isospin pairs and repulsion for antisymmetric ones, which would augment the mass of the Δ, the $(\vec{\sigma} \cdot \vec{\sigma})(\vec{\lambda} \cdot \vec{\lambda})$ structure of the color-magnetic part of the OGE gives similar contributions in both cases, diminishing the mass of the N and the Δ. The effect is much more pronounced in the case of the N, due to the presence of a spin-zero diquark, where the OGE is attractive and this effect is increased when the mass of the quark is diminished, as it happens when the temperature and the chemical potential change.

We note that in principle the OGE and confining potentials in the chiral quark model are temperature

dependent due to Debye screening, as demonstrated e.g. by a recent lattice QCD calculation in Ref. [63]. However, such a temperature dependence has a minor impact on our calculations. This is because strong modifications of the potentials appear at long distances and only for temperatures well above the critical temperature, as shown in Fig. 10 of Ref. [63]. In fact, this has been analyzed in a phenomenological manner by some of the present authors in Ref. [64], where it was shown that the quarkonia ground-state masses are almost independent of the temperature until very close to the critical temperature, above which the hadrons melt. For even larger temperatures, as discussed in the previous section, the underlying models we use lose applicability. Moreover, as explained in detail in Ref. [41], in order to evaluate the interaction kernel between two hadrons (see next section) one must subtract the self-energy contributions from the kernel and, consequently, any possible modification of confinement would drop out in the calculation of the interaction between the two hadrons.

Once we have determined the effect of the temperature and the chemical potential on the hadron masses, we know the thresholds for the study of the possible existence of meson-baryon resonances in nuclear matter.

IV. IN-MEDIUM BINDING OF $N\bar{D}$ MOLECULES

Next, we investigate how the hadron-hadron interactions are modified in a medium at finite T and μ. For this purpose we follow exactly the same scheme that has been detailed in Ref. [41], evaluating the interacting potentials with the quark and meson masses and quark-meson coupling constants determined in Sec. II for the different temperatures and chemical potentials. We note that our calculation is particularly applicable for a medium similar to the one formed in a high-energy heavy-ion collision, in which quarks coalesce to form weakly bound hadron molecules
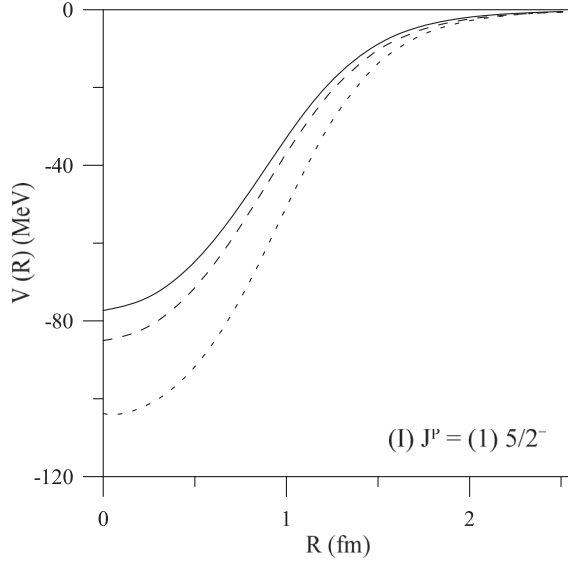
FIG. 3. $(I, J) = (1, 5/2)$ $\Delta\bar{D}^*$ interaction as a function of the temperature for different values of the chemical potential. The solid line stands for the free case $(T, \mu) = (0., 0.)$, the dashed line for $(T, \mu) = (0.12, 0.)$, and the dotted line for $(T, \mu) = (0.12, 0.15)$, where T and μ are given in GeV.

[30]. We will center our attention on the particular state highlighted in Ref. [41], the $\Delta\bar{D}^*$ state with $(I)J^P = (1)5/2^-$. We show in Fig. 3 the $\Delta\bar{D}^*$ interaction with $(I, J) = (1, 5/2)$ for some selected values of the

temperature and the chemical potential. As we can see, the effect of the medium is to strengthen the interaction. This is due to the fact that the interaction in this channel is controlled by the scalar exchange due to the almost exact cancellation of the repulsive one-gluon exchange and the attractive one-pion exchange [see Fig. 3(b) of Ref. [41]]. When increasing either the temperature or the chemical potential, the interactions grow, but the cancellation between the repulsive one-gluon exchange and the attractive one-pion exchange still remains, and the diminishing of the mass of the scalar boson with the temperature and the chemical potential, generates a stronger interaction.

To study the possible existence of an exotic state in this particular channel in the medium, we solve the Lippmann-Schwinger equation for negative energies by looking at the Fredholm determinant $D_F(E)$ at zero energy [65]. 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 $D_F(0) < 0$. We consider a baryon-meson system $Q_i R_j$ $(Q_i = N$ or Δ and $R_j = \bar{D}$ 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 $Q_i R_j$ D wave. Moreover, the baryon-meson system could couple to other baryon-meson states, $Q_k R_m$ (in the present case there would not be coupling between different physical systems). If we denote the different baryon-meson systems as channel A_i, the Lippmann-Schwinger equation for the baryon-meson scattering becomes

$$t_{\alpha\beta;IJ}^{\ell_\alpha s_\alpha, \ell_\beta s_\beta}(p_\alpha, p_\beta; E) = V_{\alpha\beta;IJ}^{\ell_\alpha s_\alpha, \ell_\beta s_\beta}(p_\alpha, p_\beta) + \sum_{\gamma = A_1, A_2, \cdots} \sum_{\ell_\gamma = 0, 2} \int_0^\infty p_\gamma^2 dp_\gamma V_{\alpha\gamma;IJ}^{\ell_\alpha s_\alpha, \ell_\gamma s_\gamma}(p_\alpha, p_\gamma)$$
$$\times G_\gamma(E; p_\gamma) t_{\gamma\beta;IJ}^{\ell_\gamma s_\gamma, \ell_\beta s_\beta}(p_\gamma, p_\beta; E), \quad \alpha, \beta = A_1, A_2, \ldots, \tag{10}$$

where t is the two-body scattering amplitude, I, 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

$$G_\gamma(E; p_\gamma) = \frac{2\mu_\gamma}{k_\gamma^2 - p_\gamma^2 + i\epsilon}, \tag{11}$$

with

$$E = \frac{k_\gamma^2}{2\mu_\gamma}, \tag{12}$$

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 $i\epsilon(\epsilon > 0)$ in the denominator. If we make the change of variables

$$p_\gamma = d\frac{1 + x_\gamma}{1 - x_\gamma}, \tag{13}$$

where d is a scale parameter, and the same for p_α and p_β, we can write Eq. (10) as

$$t_{\alpha\beta;IJ}^{\ell_\alpha s_\alpha, \ell_\beta s_\beta}(x_\alpha, x_\beta; E) = V_{\alpha\beta;IJ}^{\ell_\alpha s_\alpha, \ell_\beta s_\beta}(x_\alpha, x_\beta) + \sum_{\gamma = A_1, A_2, \cdots} \sum_{\ell_\gamma = 0, 2} \int_{-1}^1 d^2 \left(\frac{1 + x_\gamma}{1 - x_\gamma}\right)^2 \frac{2d}{(1 - x_\gamma)^2} dx_\gamma$$
$$\times V_{\alpha\gamma;IJ}^{\ell_\alpha s_\alpha, \ell_\gamma s_\gamma}(x_\alpha, x_\gamma) G_\gamma(E; p_\gamma) t_{\gamma\beta;IJ}^{\ell_\gamma s_\gamma, \ell_\beta s_\beta}(x_\gamma, x_\beta; E). \tag{14}$$

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

$$\sum_{\gamma=A_1,A_2,\cdots}\sum_{\ell_\gamma=0,2}\sum_{m=1}^{N} M_{\alpha\gamma;IJ}^{n\ell_\alpha s_\alpha,m\ell_\gamma s_\gamma}(E)t_{\gamma\beta;IJ}^{\ell_\gamma s_\gamma,\ell_\beta s_\beta}(x_m,x_k;E) = V_{\alpha\beta;IJ}^{\ell_\alpha s_\alpha,\ell_\beta s_\beta}(x_n,x_k),\tag{15}$$

with

$$M_{\alpha\gamma;IJ}^{n\ell_\alpha s_\alpha,m\ell_\gamma s_\gamma}(E) = \delta_{nm}\delta_{\ell_\alpha\ell_\gamma}\delta_{s_\alpha s_\gamma} - w_m d^2\left(\frac{1+x_m}{1-x_m}\right)^2\frac{2d}{(1-x_m)^2}\times V_{\alpha\gamma;IJ}^{\ell_\alpha s_\alpha,\ell_\gamma s_\gamma}(x_n,x_m)G_\gamma(E;p_{\gamma m}),\tag{16}$$

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. (13). If a bound state exists at an energy E_B, the determinant of the matrix $M_{\alpha\gamma;IJ}^{n\ell_\alpha s_\alpha,m\ell_\gamma s_\gamma}(E_B)$ vanishes, i.e., $|M_{\alpha\gamma;IJ}(E_B)| = 0$. We took the scale parameter d of Eq. (13) as $d = 3$ fm^{-1} and used a Gauss-Legendre quadrature with $N = 20$ points.

V. RESULTS AND DISCUSSION

The existence of charmed hadron molecules has been a topic of interest in recent years in different theoretical frameworks, as chiral quark-models [41,66], boson-exchange models [67], or effective Lagrangian approaches [68,69]. As already mentioned, our interest here is in the charmed hadron molecule $\Delta\bar{D}^*$ with isospin-spin quantum numbers $(I,J) = (1,5/2)$ that was recently predicted [41] within a chiral constituent quark model approach [60,61]. Our interest is motivated mainly by the crucial role played by an attractive short-distance quark-exchange interaction, which is a prominent feature due to the Pauli principle at the quark level. This is important because, in general, quark-exchange effects cannot be captured by an effective Lagrangian employing low-dimension hadronic operators. This feature was explicitly demonstrated in Ref. [41] for the case of the $N\bar{D}$ system by comparing predictions from the chiral quark model [60,61], and an effective Lagrangian [68,69] satisfying heavy-quark and chiral symmetries.
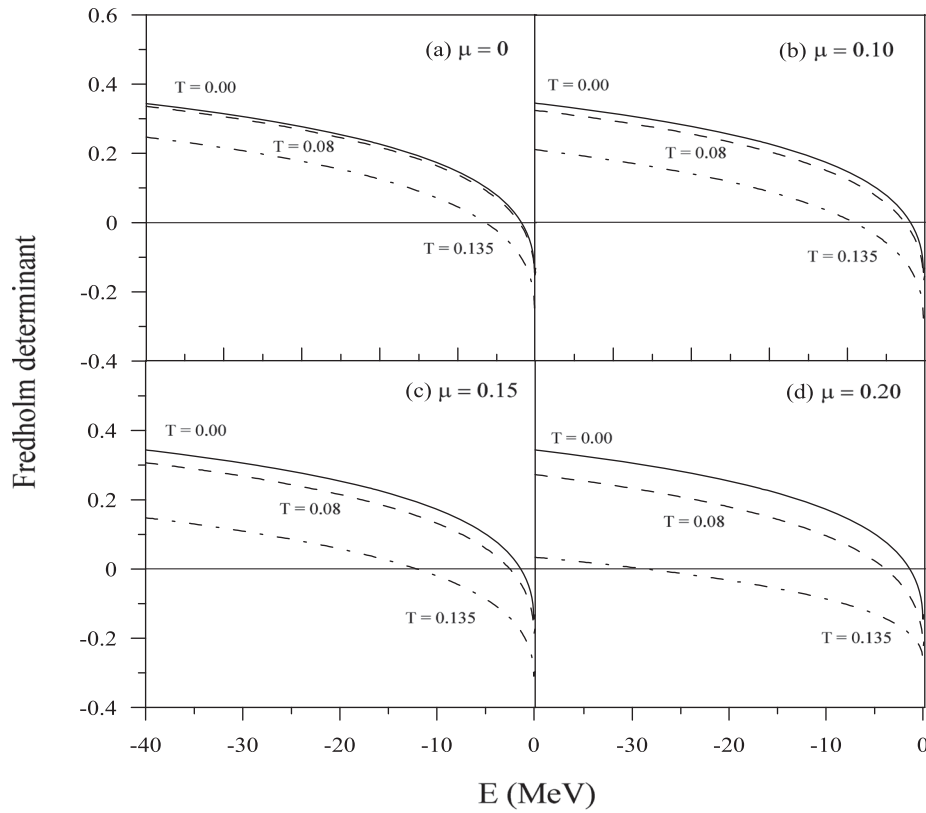

FIG. 4. $(I)J^P = (1)5/2^-$ $\Delta\bar{D}^*$ Fredholm determinant for selected temperatures of the different chemical potentials (in GeV).
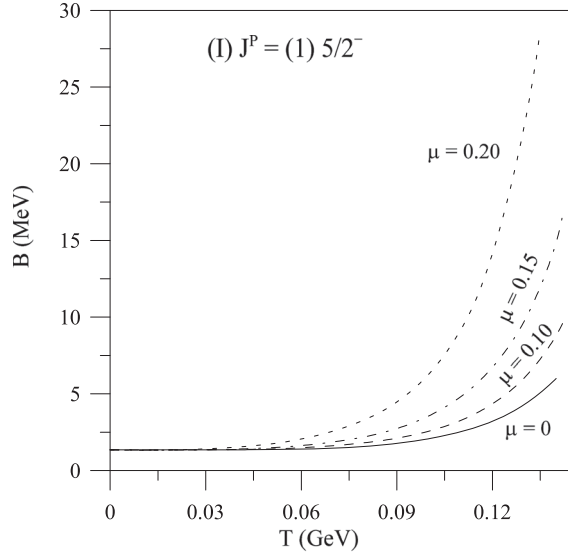
FIG. 5.   Binding energy of the $(I)J^P = (1)5/2^-$ $\Delta\bar{D}^*$ bound state, as a function of the temperature for different values of the chemical potential (in GeV).

Given the prominent role played by quark-exchange effects in free space, and the possibility that such a molecule can be formed in the environment of a heavy-ion collision [30], it is important to investigate the impact of in-medium changes in the parameters of the chiral quark model used in the evaluation of the binding energy of the molecule. The implications of our results in the coalescence dynamics in the formation of the molecule is left for a future publication.

Making use of the in-medium hadron masses and hadron-hadron interactions derived in the previous section, we have solved the Lippmann-Schwinger equation for the $(I)J^P = (1)5/2^-$ $\Delta\bar{D}^*$ system. We show in Fig. 4 the Fredholm determinant for selected temperatures of the different chemical potentials. In all cases $E = 0$ corresponds to the mass of the corresponding threshold, i.e. $M_{\bar{D}^*}(T, \mu) + M_\Delta(T, \mu)$. As pointed out in the determination of the in-medium masses of $\bar{D}$ mesons and N and Δ's, there is almost no variation of the binding energy for any chemical potential for a temperature below 0.1 GeV. For temperatures above this value it can be seen that the binding energy increases when increasing the temperature and/or the chemical potential. This is so in spite of the fact that the mass of the threshold diminishes, increasing in this way the kinetic energy. However the change of the interacting potential due to the diminishing of the mass of the scalar boson is capable of increasing the binding. In Fig. 5 we show the binding energy of the $(I)J^P = (1)5/2^-$ $\Delta\bar{D}^*$ state, as a function of the temperature for different values of the chemical potential. It varies between 1.35 MeV for the free case up to around 30 MeV for the harder system we have considered.

The bound state found in the $(I, J) = (1, 5/2)$ $\Delta\bar{D}^*$ channel would also appear in the scattering of D mesons on

nucleons as a D-wave resonance, which could in principle be measured in the near future. There are proposals for experiments by the $\bar{P}$ANDA Collaboration [31,38] to produce D mesons by annihilating antiprotons on the deuteron. They are based on recent estimations of the cross section for the production of $D\bar{D}$ pairs in proton-antiproton collisions [70,71]. The predicted bound state has quantum numbers $(I)J^P = (1)5/2^-$ and it constitutes a sharp prediction of quark-exchange dynamics because in a hadronic model the attraction appears in different channels [68,69].

The charmed molecule discussed in this paper, reflects the possible importance of quark antisymmetrization dynamics, which was already noted when studying the $\Delta\Delta$ system within the CCQM model [72]. In that case, an S-wave resonance with maximum spin was predicted. Experimental evidence of such a resonance was found in the NN scattering data, in the $^3D_3$ partial wave, and therefore as a D-wave resonance in the NN scattering [73]. This finding, the so-called $d^*(2380)$ resonance, has been used by the CELSIUS/WASA Collaboration [74] as a possible explanation of the cross section of the double-pionic fusion reaction $pn \to d\pi^0\pi^0$. The interpretation of the CELSIUS/WASA data is compatible with the formation of a $\Delta\Delta$ intermediate state, which also provides a plausible explanation for the Abashian, Booth and Crowe effect [75]: an unexpected enhancement in the double-pionic fusion of deuterons and protons to ^{3}He.

VI. SUMMARY

In brief, we have studied the effects of temperature and baryon density on the binding energy of hadronic molecules with open-charm mesons. We have selected a recent chiral quark-model-based prediction of a resonance in the $\Delta\bar{D}^*$ system with quantum numbers $(I)J^P = (1)5/2^-$. We have analyzed the modification of the thresholds and the hadron-hadron interactions when the basic parameters of the CCQM, which are the quark and meson masses and quark-meson couplings, change with temperature and baryon density as predicted by the Nambu–Jona-Lasinio model. We have found that the in-medium binding energy is enhanced as compared to its free-space binding by the presence of the hot and dense matter. This finding is relevant for ongoing heavy-ion experiments at the Relativistic Heavy-Ion Collider (RHIC) and the LHC, and for the planned experiments at FAIR and J-PARC.

As already mentioned, exotic hadrons like the hadron molecule $\Delta\bar{D}^*$ we studied here can be produced and realistically measured in high-energy heavy-ion collisions. In such collisions a hot and dense medium is created with abundant production of heavy quarks that, during the evolution of the medium, can pick up light-flavor quarks and antiquarks and form multiquark states through a coalescence mechanism [29,30]. Quark coalescence occurs

during the phase transition to the hadronic phase and it has been proven to give a very successful description of particle production in heavy-ion collisions at RHIC [76]. Coalescence depends very sensitively on the spatial extension of the produced hadrons [29] and as such, modifications to the binding energy of hadrons, in particular of molecular states, due to temperature and density will obviously affect their production rates in heavy-ion collisions. The implications of our results in the coalescence dynamics in the formation of the $\Delta \bar{D}^*$ molecule is out of the scope of the present paper and is left for a future publication.

The future research programs at different facilities like FAIR and J-PARC are expected to shed light on the uncertainties of our knowledge about the in-medium dependence of the hadron-hadron interaction with heavy flavors. While the scarce experimental information leaves room for some degree of speculation in the study of processes involving charmed hadrons, the situation can be ameliorated with the use of well-constrained models based as much as possible on symmetry principles and analogies with other similar processes. The present detailed theoretical investigation of the possible existence of bound states in a hot and dense medium is based on well-established models which provide the basic tools to make progress in the knowledge of properties of the $N\bar{D}$

interaction. It is hoped that our work is of help toward raising the awareness of experimentalists that it is worthwhile to investigate few-baryon systems, specifically because for some quantum numbers such states could be stable.

ACKNOWLEDGMENTS

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