D mesons in nuclear matter: A DN coupled-channel equations approach

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A set of coupled two-body scattering equations is solved for the DN system embedded in an isosymmetric nuclear matter. The in-medium behavior of charmed D mesons, (D^+, D^0), is investigated from the self-consistent solution within this scheme. The effective meson-baryon Lagrangian in charm quantum number one sector, the key ingredient in the present study, is adopted from a recent model by Hofmann and Lutz that has aimed at combining the charmed meson degree of freedom in a consistent manner with chiral unitary models. After a critical examination, the original model is modified in several important aspects, such as the method of regularization, to be more consistent and practical for our objective. The resultant interaction is used to reproduce the position and width of the s-wave $\Lambda_c(2593)$ resonance in the isospin zero DN channel. In the isospin one channel, it generates a rather wide resonance at $\sim$2770 MeV. The corresponding in-medium solution is then sought by incorporating Pauli blocking and the D- and π-meson dressing self-consistently. At normal nuclear matter density, the resultant $\Lambda_c(2593)$ is found to stay narrow and shifted at a lower energy, whereas the $I = 1$ resonance is lowered in position as well and broadened considerably.

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I. INTRODUCTION

The present article is devoted to the study on the behavior of the D meson in a cold symmetric (viz. total isospin zero) nuclear matter by employing a set of in-medium DN coupled-channels equations to be solved self-consistently.

Of numerous theoretical investigations to date on the properties of mesons in nuclear medium, one may notice a recent interest directed toward the in-medium behavior of the (open charm) D mesons in studying the existing, anticipated, or speculative experimental outcome in relativistic heavy-ion collisions, antiproton reactions with nuclei, possible D-nuclear bound states, etc. [1–7]. The primary theoretical effort in this regard has been to understand how the mass of the D meson gets modified in nuclear matter: either normal or hot and/or dense. The main objective is to know, for example, if the D-meson mass is reduced significantly in a medium formed by heavy-ion collisions. If this were the case, then that could lead to an enhanced D production during the processes, bringing a possible conventional hadron physics scenario for the suppression of the J/Ψ production, often attributed to the long-time speculated and more exotic process of the formation of the quark-gluon plasma due to deconfinement. In the present work we shall adopt a self-consistent many-body coupled-channels method based on a hadronic effective Lagrangian that has enjoyed its success in studying the physics of $\bar{K}$ and K mesons in nuclear matter. As a matter of fact the methods of study on this subject as employed in the above publications follow rather closely the ones as applied to the study of these mesons in nuclear medium which was initially triggered by the issue of possible kaon condensate [8,9]. They are (i) the quantum chromodynamics (QCD) sum rule method (QCDSR), (ii) the nuclear mean-field approach (NMFA), and (iii) the so-called self-consistent coupled-channels method (SCCM). However, there are approaches based on effective quark potentials such as in Refs. [10] and [11] which are in a way complementary to these three approaches. We do not discuss those quark-model methods here.

The organization of the present article goes as follows. In Sec. II we have a critical retrospect of the related works within the three methods stated above to motivate the present one. Section III is devoted to critically reviewing a series of works [7,12] which had motivations close to ours and then explain why we have come to adopt a somewhat different method by modifying what was used in those works. Section IV presents the results of our study of the DN interaction obtained from a coupled-channels equation in free space. The implementation of various medium effects on the properties of the D meson in nuclear matter is discussed in Sec. V and our results are presented in Sec. VI. Section VII is devoted to our conclusion and final remarks. Those who are familiar with the subject might skip some parts of the next section.

II. CRITICAL RETROSPECT

In what follows, we outline each of the methods mentioned above (QCDSR, NMFA, and SCCM) as the basis of why we are motivated to take the steps presented in this work. We believe this to be appropriate because, to date, no such account has been given to compare different approaches. For this objective our subsequent discussion will be heavily inclined to review the approaches employed in the physics of $\bar{K}$ in nuclear matter because the methods listed above for the D meson have been extensively used for the former. This is due to the apparent similarity between the $\bar{K}$ and D [6] as discussed later in the section.

To begin with we note that in all these three approaches (except for Refs. [3,4], to be touched on later) the central entity is the meson propagator (or the correlator) in the nuclear environment, $\langle T[D(x)\bar{D}(y)]\rangle_{\rm nuc}$, where $D(x)$ is the D-meson interpolating field.

A. QCDSR

The method [1,2] exploits the quark-hadron duality to calculate the D-meson propagator in two different ways. On the one hand it is written in terms of the underlying quark fields. Then by means of short-distance operator product expansion (OPE) within perturbative QCD, it is expressed in terms of basic QCD constants and condensates. On the other hand, within the hadronic picture it is expressed in a spectral representation with a few adjustable parameters. Then the two sides are matched in deep Euclidean region to extract hadronic quantities. In the OPE expression it is shown [2] that the essential ingredient is the product of the charm quark mass, m_c, and the *in-medium* light scalar $\bar{q}q$ condensate, $\langle\bar{q}q\rangle_{\rm nuc}$. Following Ref. [13], the in-medium $\bar{q}q$ condensate is approximated by a sum of the vacuum part and the in-medium correction, the latter being the product of nuclear density and the nucleon matrix element of $\bar{q}q$,

$$\langle\bar{q}q\rangle_{\rm nuc} \approx \langle 0|\bar{q}q|0\rangle + \rho_N\langle N|\bar{q}q|N\rangle. \tag{1}$$

On the other hand, by a factorization ansatz combined with a linear density approximation [14] valid for low nuclear density, the in-medium D propagator is written as a sum of the free space part and the in-medium correction that is proportional to the nuclear density ρ_N times the free space DN scattering amplitude, which is further approximated by the DN scattering length. Then QCDSR is eventually used to express the free-space DN scattering length in terms of a few QCD parameters, $m_c\langle N|\bar{q}q|N\rangle$ in particular [1]. Finally, by matching, the in-medium meson mass shift is obtained as proportional to the nuclear density times the thus obtained DN scattering length. This relation is naturally expected from optical models in standard scattering theory in which one obtains the in-medium meson self-energy as the product (or convolution) of nuclear density and the meson-nucleon scattering amplitude. Similar quantities appear in NMFA and SCCM to be discussed later. In the present approach the medium effect enters through the linear dependence in ρ_N only. At normal nuclear matter density the isospin averaged D mass shift is obtained as ≈ -50 MeV. In Ref. [2] an additional mass shift due to the time component of the in-medium vector $\bar{q}q$ condensate is reported as $\sim\pm 25$ MeV for $\bar{D}(D)$. A rather strong sensitivity to the assumed high energy behavior of the spectral function is noted in Ref. [2], which may be related to the difficulty in determining the free space D mass in QCDSR [15]. So along with various approximations stated earlier, results mentioned here should be regarded as semiquantitative. However, an important finding is the large (attractive) contribution induced by $\approx m_c\langle N|\bar{q}q|N\rangle$ from OPE. This term, which is of scalar-isoscalar in nature, enters just like the familiar pion-nucleon σ term. But it is at least two orders of magnitude larger because the charm quark mass: $m_c \approx 1400$ MeV, is multiplied instead of the average light quark mass.

B. NMFA

We first outline prototype of this approach used in the first kaon condensation study in Refs. [8,9]. There, the leading term in the Lagrangian is from the non-linear realization of chiral symmetry for the interaction of Goldstone bosons with octet baryons. This is supplemented by symmetry breaking terms linear in the quark mass matrix. With N and K being the nucleon and kaon fields, respectively, the resulting kaon-nucleon interaction reads,

$$\mathcal{L}_{KN} = -i\frac{3}{8f^2}\bar{N}\gamma^\mu N\bar{K}\overleftrightarrow{\partial}_\mu K + \frac{\Sigma_{KN}}{f^2}\bar{N}N\bar{K}K, \tag{2}$$

where $\bar{K}\overleftrightarrow{\partial}_\mu K \equiv \bar{K}\partial_\mu K - (\partial_\mu\bar{K})K$ and f is the Goldstone boson decay constant. The first term is the Tomozawa-Weinberg (T-W) vector interaction. The second term provides a scalar-isoscalar attraction characterized by Σ_{KN}. This quantity, called the Σ term, is expressed by three *low-energy constants* that may be written in terms of the πN and $KN\,\sigma$ terms: σ_π and $\sigma_K^{(i)}(i=1,2)$, which are the measures of chiral symmetry breaking. In NMFA, the meson (here it is the kaon) self-energy, $\Pi_K(p_0,\vec{p})$, which provides the dispersion equation relating the energy (p_0) and momentum ($\vec{p}$), is density times the nuclear expectation value of the above meson-nucleon interaction: $-\rho_N\langle\mathcal{L}_{KN}\rangle_{\rm nuc}$. With a simple Fermi gas model for an isosymmetric nuclear matter, the energy of the kaon at rest is obtained as (notice the difference between the in-medium kaon and antikaon due to the T-W vector interaction),

$$p_{0(K,\bar{K})} = \sqrt{\left(m_K^s\right)^2 + \left(\frac{3\rho_0}{8f^2}\right)^2} \pm \frac{3\rho_0}{8f^2}, \tag{3}$$

where $(m_K^s)^2 = m_K^2 - \rho_s\Sigma_{KN}/f^2$ is the square of the in-medium (scalar) kaon mass with ρ_s being the nuclear scalar density, and the ordinary nuclear matter (vector) density is $\rho_0 \equiv \rho_N$. In the non-relativistic limit, $\rho_0 = \rho_s$. From the above result, a considerable reduction in the kaon mass (more precisely the energy at zero momentum) in a high-density medium might be expected (for both K and $\bar{K}$) if the strength Σ_{KN} becomes sizable, leading even to a possible kaon condensation. Further investigation on the K mesons in hot and/or dense nuclear medium with refinement may be found, for example, in Refs. [16,17]. See Ref. [18] for an extensive set of references.

The above method has been extended to the study of the $D(\bar{D})$ mesons in an isospin symmetric nuclear matter in Ref. [5]. Two steps are required to reach the goal. The first is to describe the static nuclear properties, viz. binding energy per nucleon, compressibility, etc., in a mean-field approach imposing SU(3) symmetry with the non-linear realization of chiral symmetry [19]. It is an extention of the original Walecka "σ-ω" model [20]. The optimal solution is found by imposing a stationary condition on the free energy of the system with respect to the variation of the mean scalar- and vector-meson fields for a given value of nuclear density. The meson-nucleon coupling constants are fixed at the normal nuclear saturation density. The second step is to construct an *in-medium D-meson-nucleon* interaction. The approach used for the in-medium kaon is extended to incorporate the D meson by using gauged SU(4) vector mesons, a method adopted in part from Ref. [21]. The outcome is a Lagrangian similar to Eq. (2) in Refs. [8,9], but with a few additional terms (see

Ref. [17] for the case of kaons in matter). Of particular interest are the ones describing the interaction of the D with *mean* scalar- (σ) and vector- (ω) meson fields. They are made to contribute to the in-medium D-meson interaction Lagrangian as

$$g_{D\sigma}\bar{D}D\sigma - ig_{D\omega}\bar{D}\overset{\leftrightarrow}{\partial}_{\mu}D\omega^{\mu}, \tag{4}$$

where the mean density-dependent scalar- and vector-meson fields are complicated functions of nuclear density, effective (mean field) meson-nucleon coupling constants, and so on. Those D interactions with in-medium mesons contribute to the meson self-energy. For example, the mean ω field contribution $g_{D\omega}\omega^0$ is to be added to the one from the T-W interaction, viz. the $3\rho_0/(8f^2)$ term in Eq. (3).

In the zero-nuclear-density limit, the model is constructed such that the in-medium scalar meson contribution in Eq. (4) reduces to the meson mass term in the total D-meson Lagrangian, viz. $m_D^2\bar{D}D$, as the free D Lagrangian should have only the kinetic energy part to begin with. This is consistent with the corresponding limit in the light quark condensate: $\langle\bar{q}q\rangle_{\rm nuc} \to \langle 0|\bar{q}q|0\rangle$ in QCDSR; recall Eq. (1). However, there is a touchy issue that needs clarification regarding the D interaction with the mean vector-meson field. At zero-nuclear-density limit it is plausible that this contribution vanishes so only the T-W interaction remains. But as touched on later, the latter arises from vector-meson exchanges between D and N in the low-energy and low-momentum-transfer limit as inferred from the hidden local symmetry picture of vector mesons [22] or from the success, for example, of the ρ-ω model for the KN interaction [23] (see also Ref. [24]). So at finite nuclear density there may well be some double counting in the vector-meson exchange contribution. In this respect we refer to an interesting finding in Ref. [17]. There a transport equation simulation for *heavy ion on heavy ion* was compared with available data in the spectra of K^+ and K^- produced in the reactions. The model-data consistency has found to become troublesome upon including the in-medium kaon-ω interaction. This might actually point to an inadequacy of including the mean vector-meson interactions for the case of the D meson as well.

In the end, the quantitative details of the NMFA prediction on the D mass shift vary depending on the details of models adopted (see Ref. [17]). However, a global feature is characterized by moderate drop in the mass, obtained for the simple chiral Lagrangian of Eq. (2), viz. about 70 MeV for D and 20 MeV for $\bar{D}$ at normal nuclear matter density. The average of those two values is a rough measure of the scalar-isoscalar interaction. So just as from QCDSR, one sees a potentially important role played by this attractive force between the D and nucleon.

Another mean-field approach that directly solves for the D-meson binding in nuclear medium is presented in Refs. [3,4] within the quark-meson coupling model [25]. Schematically, it is the "σ-ω" at the level of the u and d quarks confined in the nucleon and D-meson bags. More concretely, on optimizing the mean scalar- (σ) and vector- (ρ and ω) meson fields by reproducing the static properties of nuclei (or nuclear matter), those meson fields are used to describe the in-medium D interaction at the light quark level. Note that by construction there is no explicit T-W interaction between D and N. The in-medium $D(\bar{D})$ mass is obtained by a stationary condition on varying the heavy meson bag radius. It is found that the magnitude of the vector and scalar contributions to the mass shift are comparable in magnitude. In particular, the average mass shift of the D and $\bar{D}$ mesons due to the scalar meson interaction is about 60 MeV downward for normal nuclear matter density. Again, the importance of the attractive scalar-isoscalar interaction (represented here by the σ-meson exchange) is visible here.

Before reviewing SCCM, we simply summarize the common feature of the result from QCDSR and NMFA discussed so far: (i) the D-meson interactions are due to both static in-medium scalar- and vector-type interactions and (ii) a large reduction in the D-meson mass to which a scalar-isoscalar attraction appears to play an important role.

C. SCCM

The approaches discussed so far are *static* in that the meson-nuclear interaction is introduced such that it does not disturb the mean nuclear configuration. This is achieved when the meson scattering is elastic and nearly forward by each nucleon (and by mean meson fields) in nuclear medium. Also the meson-nucleon scattering should be reasonably weak. In this respect we recall that in NMFA the meson self-energy is obtained from the meson-nucleon interaction (or potential) rather than from its full iteration: the T matrix. Therefore, the methods would become inappropriate when the two-body meson-nucleon interaction is (i) strong and, in particular, dominated by intermediate bound or resonant states and/or (ii) strongly coupled to other meson-baryon channels. The low-energy $\bar{K}N$ interaction is a typical case that does not fit into the static mean-field description. Its coupling to other meson-baryon channels such as $\pi\Sigma$, $\eta\Lambda, \ldots$, etc., is strong. Close to threshold, this interaction is dominated by the nearby $\Lambda(1405)$ resonance, which is now strongly believed to be a combination of $\bar{K}N$ and $\pi\Sigma$ s-wave molecules [26,27] embedded in the continuum of lower-threshold channels. Because of the apparent similarity between the $\bar{K}N$ and DN systems in their coupled-channels nature as well as their association with the Λ-like resonances, $\Lambda(1405)$ and $\Lambda_c(2593)$, we think it very useful to outline the SCCM used in the study of $\bar{K}$ in nuclear matter. A good part of its practical aspects is effective for our present in-medium D study. See a prototype of this approach in Ref. [28].

Here again one needs two steps to achieve the goal. But unlike NMFA the first step is for the two-body aspect, then the many-body aspect enters later. First, the free-space $\bar{K}N$ multi-channel Bethe-Salpeter equation:

$$T = V + VGT, \tag{5}$$

is solved. Here, T is the transition operator matrix, G a diagonal matrix each element of which is the product of single-particle propagators for a meson and a baryon, and the potential (driving term) V is a matrix whose elements are T-W type meson-baryon interactions from the lowest-order terms in the nonlinear chiral Lagrangian. In on-shell approximation the solution T was shown to successfully reproduce the $\Lambda(1405)$ resonance and other reaction observables [29,30].

Next, the same set of equations is solved in nuclear medium. The underlying assumption is that the potential term V is unaltered in nuclear medium, but that all the medium effects enter through the intermediate meson-baryon propagators [31–36]. This in-medium propagator includes the effects from (i) Pauli blocking, (ii) binding of baryons by nuclear mean field, and (iii) dressing (or self-energy) of intermediate state mesons ($\bar{K}$, π, or η) due to their interactions with surrounding nucleons. The resultant quantity is denoted as $\tilde{G}$. The in-medium equation now reads

$$\tilde{T} = V + V\tilde{G}\tilde{T}. \tag{6}$$

Because of the meson dressing, particularly the dressing of $\bar{K}$ that creates nested $\bar{K}N$ interactions, this in-medium equations must be solved *self-consistently*. The thus obtained diagonal amplitude, $\tilde{T}(\bar{K}N \to \bar{K}N)$, already demonstrates certain essential features of in-medium $\bar{K}$. But a more suitable quantity to study is the kaon spectral function $S_{\bar{K}(K)}(p_0, \vec{p})$ which is proportional to the imaginary part of the kaon self-energy $\Pi_{\bar{K}(K)}(p_0, \vec{p})$. In free space, we have a trivial on-mass shell relation: $S_{\bar{K}(K)}(p_0, \vec{p}) = \delta[p_0 - E(\vec{p}\,)]/(2p_0)$, with $E(\vec{p}) = \sqrt{m_K^2 + \vec{p}^{\,2}}$. In the nuclear medium, this structure changes substantially, such that the $\bar{K}$ *meson mass pole* is not only shifted somewhat downward but also broadened. Moreover, there is an additional structure due to the in-medium $\Lambda(1405)$ resonance [31–36]. These aspects cannot be obtained from the NMFA or QCDSR approaches discussed earlier. Notably, within SCCM there has been no indication of a possible onset of kaon condensation even at higher densities. So one sees the possible importance of medium effects taken care of self-consistently which have made the difference. However, this difference might also be due, in part or to a good extent, to the fact that so far the equations in SCCM have been driven only by the T-W vector interaction without any additional ones such as the attractive Σ_{KN}-term in the diagonal $\bar{K}N$ channel present in NMFA. Recently, there have been several works on the *improved* coupled $\bar{K}N$ equations in free space which incorporate the *next-to-leading-order* interactions, including the corresponding scalar-isoscalar (or Σ_{KN} term) contribution [37–40]. Such additional terms have certainly improved the fit to available data by ≈20% thanks to several additional parameters related to them. With such a new type of interactions, one may wonder if the kaon mass could reduce sufficiently in the nuclear medium to give rise to kaon condensation. However, according to a recent work on in-medium $\bar{K}$ in isospin asymmetric matter which incorporates this type of contribution [35], kaon condensation does not appear to set in below eight times the nuclear matter saturation density. This is well beyond the limit of applicability of the model; certainly around this *critical* density new degrees of freedom, both hadronic and subhadronic, will have to be taken into account.

We now refer to a couple of exploratory works on the D meson within the same framework. First, a coupled-channels calculation for the D mesons in cold nuclear matter was done [6], motivated by the similarity between the DN and $\bar{K}N$ systems once the s quark in the later is replaced by a c quark. This is further reinforced by an apparent correspondence between the two $I = 0$, s-wave resonances: $\Lambda(1405)$ and $\Lambda_c(2593)$, in the coupled $\bar{K}N$ and DN channels, respectively. To make this analogy more concrete, free space amplitudes are obtained from a set of separable coupled-channels potentials simulating the T-W type interactions to reproduce the $I = 0$ $\Lambda_c(2593)$ as a DN s-wave *hadronic molecular* state of binding energy of ≈200 MeV with a width of ≈3 MeV, sitting very close to the $\pi\Sigma_c$ threshold. In the $I = 1$, s-wave channel, the model appears to have generated a resonance at about 2800 MeV but is less conspicuous than the fitted one in the $I = 0$ channel. Then the corresponding interaction is fed into the in-medium equation. A notable feature is the relative importance of the intermediate state pion dressing. The final result has found a slight upward-shifted and broadened D meson *pole* with a wiggle in the spectral function at normal nuclear matter density. One of the peaks in the wiggle corresponds to the shifted-broadened D pole, whereas the other seems to originate from the resonant structure in the in-medium $I = 1$ DN amplitude. Somewhat surprising is the apparent absence of an anticipated peak due to the $\Lambda_c(2593)$ resonance. The peak is visible when pion dressing is ignored in the coupled-channels problem. This work has been recently extended to finite temperature in Ref. [41].

Here a question remains as to if the simple prescription of $s \to c$ quark replacement be adequate to model the DN coupled-channels interaction. By so doing all the two-body channels with strangeness, such as $D_s\Lambda$, have been excluded. However, as we will show later, they have an important effect in the DN coupled-channels problem. In addition, from the point of view of symmetries, one must recall the following well-known fact: although the light Goldstone bosons such as π and K mesons are dictated by chiral symmetry, the charmed mesons such as D are quite heavier and obey the *heavy-quark symmetry*, the extreme opposite to the former. A blind $s \to c$ replacement breaks both of those symmetries.

A different approach, which respects the proper symmetries, was attempted in Ref. [42]. There, charmed baryon resonances are generated dynamically from the scattering of Goldstone bosons off ground-state charmed baryons with $J^P = \frac{1}{2}^+$. The $C = 1$, $S = I = 0$ resonance found at 2650 MeV was identified with the $\Lambda_c(2593)$ in spite of the fact that the width, due to the strong coupling to $\pi\Sigma_c$ states, is obtained as more than 20 times the experimental value of about 4 MeV. The trouble with this model is that couplings to DN and D_sY are completely absent. We recall that in Ref. [6] the former channel is essential in the formation of the $\Lambda_c(2593)$.

A satisfactory improvement came in a recent work [12] where the alleged shortcomings have been overcome by exploiting the universal vector-meson coupling hypothesis to break the SU(4) symmetry in a convenient and well-defined manner. More precisely, this is done via t-channel exchange of vector mesons between pseudoscalar mesons and baryons in such a way to respect chiral symmetry for the light meson sector and the heavy quark symmetry for charmed mesons, as well as to maintain the interaction to be of the T-W vector type. The model generates the $\Lambda_c(2593)$ resonance in the $C = 1$, $S = I = 0$ s-wave channel, as well as an almost degenerate s-wave resonance at 2620 MeV in the $C = 1$, $I = 1$, $S = 0$ channel not found experimentally so far. An application of this model to a preliminary study of D and D_s mesons in nuclear matter may be found in Ref. [7]. There the D-meson spectral function is found to have a peak, embedding the two resonances mentioned above, and another structure signaling the D-meson pole position, which is shifted upward by about 30 MeV and has a finite width. So, with respect to the D-pole mass shift, the two SCCM results show the opposite tendency to the one found both in QCDSR and NMFA, upon disregarding other attributes such as finite widths, and so on. Then one may wonder from where this difference originates: is it attributable to the coupled-channels aspect or to the absence of an extra scalar-isoscalar attraction in the two SCCM models?

Having discussed salient aspects of various models to date, we have naturally adopted SCCM in the present study of the D meson in symmetric nuclear matter. In particular, we employed a modified version of the coupled-channels method developed by Hofmann and Lutz [12] after its critical analysis in the next section. Specifically, we introduced a cut-off regularization as well as an extra phenomenological scalar-isoscalar attraction in the diagonal DN channel in a simplistic manner to study its implication.

III. WHY HAVE WE ADOPTED BUT MODIFIED THE HOFMANN-LUTZ MODEL?

The original Hofmann-Lutz model [12] is ambitious enough to include all the $J^P = \frac{1}{2}^-$ s-wave pseudoscalar-baryon interactions with charm quantum number values up to 3 in an attempt to interpret/predict various baryon resonances as molecular states. Our present interest in this model is only in the sectors with quantum numbers $C = 1$, $S = 0$, $I = 0$ and $I = 1$ which are associated with the DN channel. As for the $\bar{D}N$ sector with quantum numbers $C = -1$, $S = 0$, QCDSR [1,2], the quark-meson coupling model [3,4], as well as the simple quark model suggest that the interaction is weak and quite likely repulsive. See also Ref. [7]. Hence we are not concerned with this sector in the present work.

The Hofmann-Lutz model connects two sets of characteristic pseudoscalar meson-baryon sectors with one charmed quark belonging either to mesons, e.g., DN, or to baryons, such as $\pi\Sigma_c$, etc., by means of the universal vector-meson coupling hypothesis equipped with the Kawarabayashi-Suzuki-Fayyazuddin-Riazudden (KSFR) condition [43]. Its modern theoretical support is offered in the hidden local symmetry picture of vector mesons; see Ref. [22]. In this respect the model is an important first step for improvement because heavy-quark effective theory equipped with chiral symmetry, chiral heavy-quark effective theory (χHQET), can deal only with Goldstone bosons interacting with charmed baryons (or charmed mesons); see Refs. [42,44–46]. It cannot be applied to channels of our interest such as DN. While retaining the physical hadron masses, the Hofmann-Lutz model uses SU(4) symmetry to construct the effective interaction between pseudoscalar mesons in 16-plet with baryons in 20-plet representations through a t-channel exchange of a 16-plet of vector mesons. The universal vector-meson coupling hypothesis provides the global interaction strength among the above SU(4) multiplets. Then, aided by the KSFR relation that is consistent with chiral symmetry at very low energy and momentum transfer, the resultant lowest-order meson-baryon interaction is found to take a near T-W form in the $t = 0$ limit. An interesting and important consequence of this picture is that, compared with the ones that exploit SU(4) symmetry alone, the T-W interactions resulting from an exchange of a charmed meson are reduced by an extra factor, $\sim(m_V/m_V^c)^2$, where the masses here are the typical (uncharmed) vector meson and a (singly) charmed meson, respectively. We discuss this aspect further below.

Now we are in the position to obtain the interaction V in our present work based, in good part, on the Hofmann-Lutz model. The two sectors of our interest are all s-wave and have $J^P = \frac{1}{2}^-$. We retain the following channels:

$$\pi\Sigma_c(2589),\quad \boldsymbol{DN}(2810),\quad \eta\Lambda_c(2835),\quad K\Xi_c(2960),\quad K\Xi_c'(3071),\quad \boldsymbol{D_s\Lambda}(3085),\quad \eta'\Lambda_c(3245)$$

for the $C = 1$, $I = S = 0$ sector, and

$$\pi\Lambda_c(2425),\quad \pi\Sigma_c(2589),\quad \boldsymbol{DN}(2810),\quad K\Xi_c(2960),\quad \eta\Sigma_c(3005),\quad K\Xi_c'(3071),\quad \boldsymbol{D_s\Sigma}(3160),\quad \eta'\Sigma_c(3415)$$

for the $C = 1$, $I = 1$, $S = 0$ sector. Here, channels containing charmed pseudoscalar mesons are denoted in bold letters. The values between parentheses following each channel in the above expressions are the corresponding channel thresholds in MeV. The transition interaction (potential) for $i \leftrightarrow j$ due to t-channel exchanges of vector meson "X" reads (note a somewhat different notation from Ref. [12])

$$V_{ij}(q_i, q_j; \sqrt{s}) = \sum_X g^2 C_{ij}^X \bar{u}(\vec{p}_i)\gamma^\mu \left[g_{\mu\nu} - \frac{k_{\bar{\mu}} k_\nu}{m_X^2} \right] \times \frac{1}{t - m_X^2} (q_i + q_j)^\nu u(\vec{p}_j), \tag{7}$$

where g ($\approx$6.6) is the universal vector-meson coupling constant, C_{ij}^X is the product of relevant SU(3) Clebsch-Gordan coefficients associated with quantum numbers of the vertices projected to the ij channels, and p_i, q_i, p_j, q_j are the four momenta of the baryon and meson in channels i and j, respectively. As usual, $s = (p_i + q_i)^2 = (p_f + q_f)^2$ and the momentum transfer is $k \equiv q_i - q_j = p_j - p_i$, with $t \equiv k^2 = k^\mu k_\mu$.

The next step is to expand $1/(t - m_X^2)$ in powers of t/m_X^2. The second term in this expansion, viz. t/m_X^2, tends to compensate to a good extent the term $k_\mu k_\nu/m_X^2$ in the numerator of the vector-meson propagator, so to a good approximation, the t dependence is at most $O[(t/m_X^2)^2]$. Thus for our objective it is consistent to disregard the terms $O(1/m_X^2)$ altogether and make the interaction to be of zero range. Note that the Hofmann-Lutz model has retained the $O(1/m_X^2)$ term in the numerator. The effect due to this additional contribution is examined later, particularly for charmed vector-meson exchanges. Then we adopt the average mass $\bar{m}_V$ for all the uncharmed nonet vector mesons as well as $\bar{m}_V^c$ for the anti-triplet charmed vector mesons. The difference caused by this simplification is found to be quite small. By dropping the p-wave contribution from $\bar{u}(\vec{p}_i)u(\vec{p}_j)$, we now have an s-wave, zero-range limit of the above interaction in the following T-W form upon adopting the normalization convention ($\bar{u}u = 1$) used in Ref. [26],

$$V_{ij}^I(\sqrt{s}) = -\frac{\kappa C_{ij}}{4f^2}(2\sqrt{s} - M_i - M_j)\left(\frac{M_i + E_i}{2M_i}\right)^{1/2} \times \left(\frac{M_j + E_j}{2M_j}\right)^{1/2}. \tag{8}$$

Here the left-hand side is the ij element of the (on-shell) interaction matrix V in Eqs. (5) and (6). On the right-hand side, $C_{ij} = \sum_X C_{ij}^X$, M_i and M_j, as well as E_i and E_j are the masses and energies of baryons in channels i and j, respectively. In addition, f is the pseudoscalar meson decay constant from the KSFR relation: $\bar{m}_V^2/g^2 = 2f^2$. We also introduced κ, a reduction factor which is *unity* for transitions $i \leftrightarrow j$, driven by (uncharmed) vector-meson exchanges (ρ, ω, ϕ, K^*), etc., but is equal to $\kappa_c = (\bar{m}_V/\bar{m}_V^c)^2$ for charmed vector-meson exchanges such as D^* and D_s^*. The transition coefficients $\tilde{C}_{ij} \equiv \kappa C_{ij}$, which are symmetric with respect to the indices, are listed in Tables I and II. The reader may notice that the thus obtained T-W interaction strengths C_{ij} are simply the consequence of SU(4) symmetry *modulo* sign convention and that the vector-meson exchange picture shows a definite pattern of breaking the SU(4) symmetric interaction.

Before testing and then using the resultant interaction, we want to check one important aspect of the approximation we have made to reach the T-W form. Let us take Eq. (7) for nondiagonal transition interactions by charmed meson exchanges and, as an extreme case, consider the one for $DN \leftrightarrow \pi\Lambda_c$. Here the variable $t = k^2$ in $1/[t - (m_V^c)^2]$ is far from zero but could be as large as $t \approx +M_N^2$, where M_N is the nucleon mass. Also, in such a transition, the $k_\mu k_\nu/(m_V^c)^2$ contribution in the numerator, which we have disregarded, might significantly affect the magnitude of the driving term: in the channel under consideration it reduces the size of the driving term by more than 50% in the energy range of our interest. However, we have confirmed numerically that in the same energy range those two contributions tend to mutually compensate such that neglecting the two together makes a maximum deviation of $O\left[(t^2/(\bar{m}_V^c)^4\right] \approx 25\%$ as compared with the original t-channel charmed vector-meson exchange interaction projected on to the s wave. Near and above the DN threshold it is only 10% or less. Thus as stated earlier, our procedure has turned out to be not only simpler but more consistent than what is adopted in Ref. [12]. Combined with the reduction factor κ_c already multiplied to this type of transitions, our ordinary T-W form of the interaction is consistent with the lowest-order chiral symmetry as well as heavy quark symmetry by use of the extended KSFR relation. In actual calculations we simply set $\kappa_c = 1/4$ and the resulting amplitudes are found stable against a small variation around this value.

TABLE I. The $\tilde{C}_{ij}$ coefficients for the $C = 1, S = 0$ meson-baryon interaction for isospin $I = 0$.

	$\pi\Sigma_c$	DN	$\eta\Lambda_c$	$K\Xi_c$	$K\Xi_c'$	$D_s\Lambda$	$\eta'\Lambda_c$
$\pi\Sigma_c$	4	$\sqrt{\frac{3}{2}}\kappa_c$	0	0	$\sqrt{3}$	0	0
DN		3	$-\frac{1}{\sqrt{2}}\kappa_c$	0	0	$-\sqrt{3}$	$-\kappa_c$
$\eta\Lambda_c$			0	$-\sqrt{3}$	0	$-\sqrt{\frac{2}{3}}\kappa_c$	0
$K\Xi_c$				2	0	$-\frac{1}{\sqrt{2}}\kappa_c$	0
$K\Xi_c'$					2	$-\sqrt{\frac{3}{2}}\kappa_c$	0
$D_s\Lambda$						1	$\frac{1}{\sqrt{3}}\kappa_c$
$\eta'\Lambda_c$							0

TABLE II. The $\tilde{C}_{ij}$ coefficients for the $C = 1, S = 0$ meson-baryon interaction for isospin $I = 1$.

	$\pi\Lambda_c$	$\pi\Sigma_c$	DN	$K\Xi_c$	$\eta\Sigma_c$	$K\Xi_c'$	$D_s\Sigma$	$\eta'\Sigma_c$
$\pi\Lambda_c$	0	0	$-\sqrt{\frac{3}{2}}\kappa_c$	1	0	0	0	0
$\pi\Sigma_c$		2	$1\kappa_c$	0	0	$\sqrt{2}$	0	0
DN			1	0	$\frac{1}{\sqrt{6}}\kappa_c$	0	-1	$\frac{1}{\sqrt{3}}\kappa_c$
$K\Xi_c$				0	0	0	$\sqrt{\frac{3}{2}}\kappa_c$	0
$\eta\Sigma_c$					0	$-\sqrt{3}$	$\sqrt{\frac{2}{3}}\kappa_c$	0
$K\Xi_c'$						0	$-\frac{1}{\sqrt{2}}\kappa_c$	0
$D_s\Sigma$							1	$-\frac{1}{\sqrt{3}}\kappa_c$
$\eta'\Sigma_c$								0

The next step is to confirm the relevance of the interaction as obtained above. We first check the resulting amplitudes in free space. Here an on-shell ansatz (equivalent to the N/D method) has been employed which allows for reducing the coupled integral equations, Eq. (5), to a single matrix equation whose solution may be written as

$$T = (I - V\hat{G})^{-1}V, \tag{9}$$

where $\hat{G}$ is a diagonal matrix whose elements are now four-momentum integrated propagators of the channels involved. Momentum integrations have been regularized by a dimensional method as found in Refs. [26,47]:

$$\begin{aligned} G_i(\sqrt{s}) &= i2M_i \int \frac{d^4q}{(2\pi)^4} \frac{1}{(P-q)^2 - M_i^2 + i\varepsilon} \frac{1}{q^2 - m_i^2 + i\varepsilon} \\ &= \frac{2M_i}{16\pi^2}\left[a_i(\mu) + \ln\frac{M_i^2}{\mu^2} + \frac{m_i^2 - M_i^2 + s}{2s}\ln\frac{m_i^2}{M_i^2} \right. \\ &\quad \left. + \frac{\bar{q}_i}{\sqrt{s}}\ln\frac{M_i^2 + m_i^2 - s - 2\bar{q}_i\sqrt{s}}{M_i^2 + m_i^2 - s + 2\bar{q}_i\sqrt{s}}\right], \end{aligned} \tag{10}$$

where $\bar{q}_i$ is the on-shell momentum and m_i and M_i are the meson and baryon masses in channel i. With the regularization scale set to $\mu = 1.0$ GeV and imposing the subtraction points following [12], we found that all the corresponding subtraction constants $a_i(\mu)$ stay close to *the natural size*, viz. ≈ -2.0

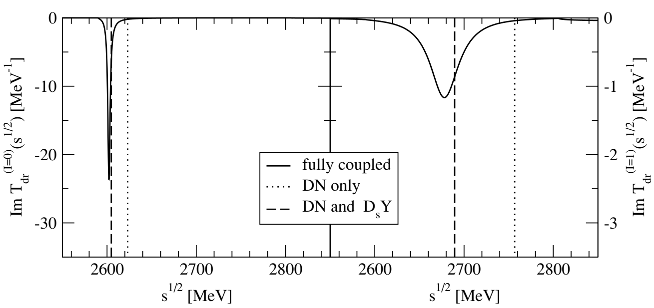

FIG. 1. Imaginary part of the DN amplitude in the dimensional regularization scheme as a function of $\sqrt{s}$ for $I = 0$ (left panel) and $I = 1$ (right panel).

[26,47]. With this we calculate the $C = 1$, $S = 0$, $I = 0, 1$ amplitudes and find the positions of the Λ_c and Σ_c resonances at ≈2620 and ≈2680 MeV, respectively, as in Ref. [12] for the model they refer to as the SU(4) *symmetric* case. With a very small change in the value of the subtraction constant for the DN channel in the $I = 0$ sector, namely $a_{DN}(\mu)$: $-1.92 \to -1.97$, we were even able to adjust the position of the resonance to the empirical one for $\Lambda_c(2593)$. However, a larger change in the corresponding $a_{DN}(\mu)$ value appears to have been required to shift the position of the $I = 1$ resonance down to 2620 MeV as predicted in Ref. [12] for the model in which SU(4) *symmetry* is broken upon shifting the value of the universal vector-meson coupling constant by up to ≈20% in some channels. In the present work, we have *not* adopted any such modifications. Our results are shown by the solid lines in Fig. 1. The meaning of the other two vertical lines is discussed in the next section.

As for the widths of those resonances, our values are far larger than those by Hofmann and Lutz for both isospin sectors. The Hoffman-Lutz prediction for the $I = 0$ resonance is no more than 0.2 MeV, whereas our estimate is ∼3.0 MeV, closer to the experimental value. A more dramatic difference is found in the $I = 1$ sector: the Hofmann-Lutz prediction is ∼3.3 MeV or lower, which is in sharp contrast to our large estimated value of ∼35 MeV. We have been able to trace the origin of this difference to the $k_\mu k_\nu/(m_V^c)^2$ term, retained in the Hofmann-Lutz model. Because this term can reduce the charm exchange transition interaction by about a factor of 2 or more, the width of the DN-type resonances decaying into $\pi\Lambda_c$, $\pi\Sigma_c$ states are correspondingly smaller. As discussed earlier, this term should not be retained for consistency, hence our larger resonance widths should be preferred. The prediction for the position of the $I = 1$ resonance is revisited in the next section, where we will introduce an explicit cut-off regularization.

Upon confirming that we have properly adopted the Hofmann-Lutz model with a few simplifications, some of which have led to improvements, we go on to include in-medium effects in the amplitude. To achieve this we have followed Refs. [33] and [7] and solved

$$\tilde{T} = T + T(\tilde{G} - G)\tilde{T}, \tag{11}$$

which results from combining Eqs. (5) and (6), where we have included the Pauli blocking effect as well as the D-meson dressing. Much of the method of solution has been taken from Ref. [32] and will be briefly described in Sec. V. Equation (11) must be solved self-consistently just like Eq. (6) as the dressed propagator $\tilde{G}$ contains the solution $\tilde{T}$. So we have iterated the equation by starting from the free-space solution T. Although the input T has been obtained by dimensional regularization, it is found that an explicit cutoff must be introduced in calculating $(\tilde{G} - G)$ to extract a tempered in-medium solution. The exception to this is when only Pauli blocking is taken into account. Simply this is because of the vanishing contribution to $(\tilde{G} - G)$ from momenta outside the nuclear Fermi sea. In Refs. [7] and [33] an explicit cut-off value of 800 MeV was used, which appears reasonable as an educated guess. However, it is important to check the stability of the solution against the change in the cut-off value. So we have varied the upper limit of the momentum integration within the range accepted in effective hadron physics. As demonstrated in Fig. 2, the physics extracted in this way varies wildly as a function of the cut-off value, thus no reliable prediction of in-medium amplitudes appears possible.

In an attempt to overcome this trouble, we went back and took another look at the two original equations, Eqs. (5) and (6). Regarding the free-space one, viz. Eq. (5), one must remember that the resonance positions and widths which it generates depend inherently on how one regularizes the divergent integral of the loop G. Employing a dimensional regularization scheme, as in Ref. [12] and in several other works, has the advantage that the divergent part can be isolated analytically. But for the in-medium equation, Eq. (6), the divergent part cannot be identified unambiguously as $\tilde{G}$ is available only numerically, so one is forced to introduce an explicit cutoff. Then in the absence of any well-defined rule to relate the two regularization schemes, one would have no idea as to what cut-off value is relevant, nor to what extent the properties of the free space solution T might have been carried over to the in-medium solution $\tilde{T}$. Furthermore, mathematically the free

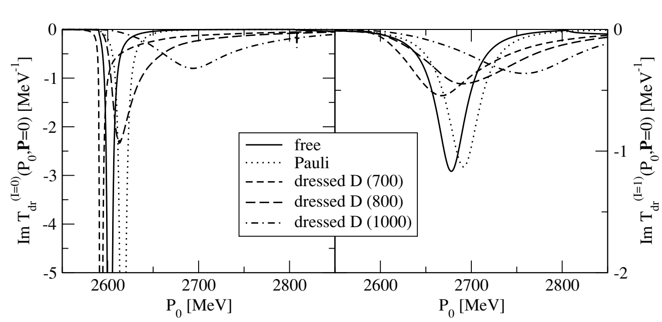

FIG. 2. Imaginary part of the in-medium DN amplitude at total momentum $\vec{P} = 0$ as a function of total energy P_0 for $I = 0$ (left panel) and $I = 1$ (right panel). The free amplitude in the dimensional regularization scheme of Fig. 1 (solid lines) is modified by medium corrections calculated according to Eq. (11) with various values of the cutoff: 700 MeV (short-dashed lines), 800 MeV (long-dashed lines), and 1000 MeV (dot-dashed lines). The dependence of the corrections on the cutoff is evident, except when one only considers Pauli blocking (dotted line) where the correction affects only loop momenta up to $p_F = 270$ MeV, well below any of the cut-off values explored.

and in-medium equations cannot be combined into a single one, viz. Eq. (11), when the support of the integration in the two equations is not identical. Hence, the method employed above to solve Eq. (11) is an inconsistent one. In view of this, we discarded the dimensional method but adopted the *conventional* cut-off method for the solution to the free-space equation as well. This is described in the next section.

IV. FREE SPACE AMPLITUDES IN CUT-OFF METHOD

From our discussion in the last section, we conclude that only the direct cut-off regularization method is left to us as appropriate and practical for our later study of the in-medium D properties. So here we apply it identically to both free and in-medium equations in a manner used in Ref. [32] for the study of $\bar{K}$ in nuclear matter. In this section we construct a set of free-space amplitudes in this scheme. A novel feature here is that we supplement the T-W vector interaction discussed in the previous section, viz. Eq. (8), with a scalar-isoscalar attraction: recall again our discussion on its possible importance in Sec. II. Here we follow a simple and conventional treatment of this term used in the kaon condensate studies and write it as (see Ref. [5,18] and also the last term in Eq. (2)):

$$\mathcal{L}_\Sigma \equiv \frac{\Sigma_{DN}}{f_D^2}\bar{N}N\bar{D}D. \tag{12}$$

In the above expression f_D is the D meson weak decay constant, and Σ_{DN} is the strength of this interaction. Note that for simplicity we introduce this only in the diagonal DN interaction: presumably similar terms might claim their right in the diagonal D_sY, $(Y = \Lambda, \Sigma)$ interactions because these channels couple strongly to DN as understood from Tables I and II. However, we want to look for possible effects from such scalar-isoscalar attractions in a semi-quantitative manner, thus preferring to contain the number of parameters. The s-wave projection of this interaction is simply equal to

$$V_\Sigma(\sqrt{s}) = -\frac{\Sigma_{DN}}{f_D^2}\left(\frac{M_N+E}{2M_N}\right), \tag{13}$$

for both $I = 0, 1$ DN channels. Concerning f_D, its most recent extraction from the branching ratios of $D^+ \to \mu^+\nu$ is about 223 MeV [48]. Various calculations and measurements as cited in this reference do seem to agree within about 10%. For simplicity we adopted $f_D = 200$ MeV. As for the value of Σ_{DN}, we simply follow what QCDSR [1] and NMFA of Ref. [5] suggest and estimate it conservatively as $\Sigma_{DN} \approx 2000$ MeV. We *also* accommodate the case where no such attraction is added, hence $\Sigma_{DN} = 0$ then.

With the above preparation, we solve the coupled-channels equations in free space and reproduced the $\Lambda_c(2593)$ resonance in the $I = 0$ sector, as seen on the left panel of Fig. 3, which shows the imaginary part of the diagonal $I = 0$ DN amplitude for both models. The parameters for those two cases are as follows: model A: $f = 1.15f_\pi$, $\Sigma = \Sigma_{DN}/f_D^2 = 0.05$ MeV^{-1}, $\Lambda = 727$ MeV, and model B: $f = 1.15f_\pi$, $\Sigma = \Sigma_{DN}/f_D^2 = 0$ MeV^{-1}, $\Lambda = 787$ MeV, where Λ is the ultraviolet cut-off value for the integration in the loop G (the value of f is loosely fixed in conformity with Ref. [29]). The width of the $I = 0$ resonance is found to be $\sim$4 MeV for model A and $\sim$5 MeV for model B, respectively. Note that in the same isospin sector, another resonance very close to $\Lambda_c(2593)$ but far wider has been identified, as in Ref. [12] that is due to the chiral excitation in the charmed hyperon channels [42] and is connected to the DN channel by a charmed vector-meson exchange. So this is more easily seen in the $\pi\Sigma_c$ channel. The same two sets of parameters are now used in the $I = 1$ coupled channels, and we found a somewhat different prediction as seen in the right panel of Fig. 3 for the diagonal $I = 1$ DN amplitude. Although a resonance is generated in both models, the one with an extra attractive Σ_{DN} interaction (model A) pulls the resonance lower to about 2770 MeV (with a width $\sim$25 MeV), whereas the one without it (model B) keeps the resonance position at $\approx$2795 MeV, not very far from the DN threshold: $\approx$2810 MeV, together with a width of $\sim$20 MeV. Tentatively we call this $I = 1$ resonance as $\Sigma_c(2770)$.

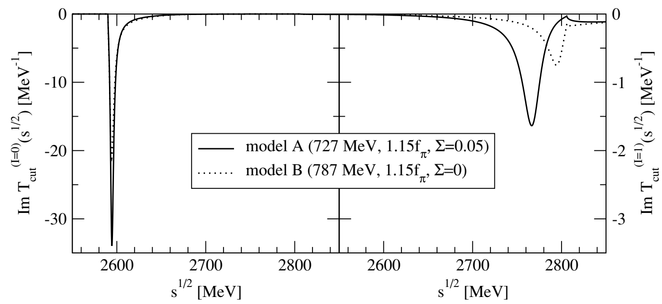

FIG. 3. Imaginary part of the DN amplitude as a function of $\sqrt{s}$ for $I = 0$ (left panel) and $I = 1$ (right panel), obtained by using the cut-off scheme with model A (solid lines) and model B (dotted lines).

To understand how the difference in the two isospin sectors comes about, it should be useful to have an *anatomical study* of the underlying mechanism for resonance formation. The following argument can be made qualitatively, based on Tables I and II, and can then be substantiated quantitatively by actual calculations. First, we see that the T-W diagonal DN interaction is attractive in both isospin channels, hence when channel couplings are turned off, the s-wave DN state may form a bound state in both $I = 0$ and $I = 1$ sectors, as seen by the dotted lines in Fig. 4, corresponding to model A,

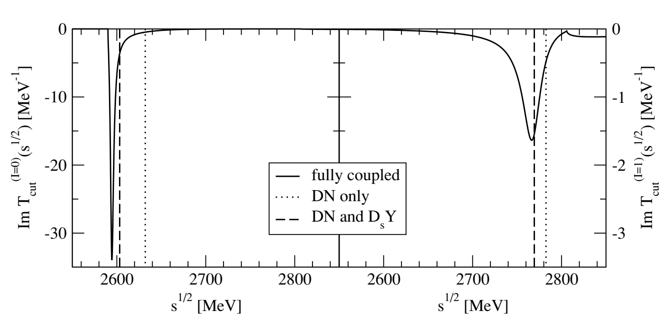

FIG. 4. Imaginary part of the DN amplitude as a function of $\sqrt{s}$ for $I = 0$ (left panel) and $I = 1$ (right panel) obtained by using the cut-off scheme (model A).

although the same behavior is found for model B, viz. with and without the additional attraction by Σ_{DN}. The binding is, of course, deeper for $I = 0$ because the T-W interaction is three times stronger. Next, we introduce channel couplings DN-$D_sY(Y = \Lambda, \Sigma)$. Note that this type of coupling is absent in Ref. [6], where particles with strangeness are excluded. In both isospin channels the corresponding strengths are comparable to the DN diagonal interaction. Because the thresholds for these D_s involved channels are higher than those with D, the channel coupling brings additional attraction to the DN binding, an effect that can be easily justified from second-order perturbation arguments. Third, the remaining couplings with the DN channel are from those connected by a charm transfer, hence suppressed by $\kappa_c \approx 1/4$, so the additional shift in the bound-state pole positions is smaller. The only apparent effect from coupling to κ_c-suppressed channels on the diagonal DN amplitudes is the finite width due to those channels with lower thresholds, hence transforming the bound states to resonances.

By comparing the results of our explicit cut-off models as in Figs. 3 and 4 with the ones from the dimensional scheme shown in Fig. 1, one sees rather large differences in the $I = 1$ sector. In the latter the downshift in the DN *bound-state* pole position due to coupling to the D_sY channel is found larger than in the $I = 0$ sector, just the opposite to what one finds in the explicit cut-off scheme. In addition, the dimensional scheme finds the $I = 1$ resonance position at 2680 MeV, quite lower than the one from the cut-off method. To understand the possible origin of the differences, we look at the choice of the subtraction point in the dimensional approach as discussed in Ref. [12], which was taken in each sector (with a set of definite quantum numbers) at $\mu = \sqrt{m^2 + M^2}$, where m and M are the meson and baryon masses in the channel in which $m + M$ is minimum. In the present $C = 1$, $S = 0$ case, $m = m_\pi$ for both isospin sectors, whereas $M = M_{\Sigma_c}$ for $I = 0$ and $M = M_{\Lambda_c}$ for $I = 1$. This definition is somewhat different from the one used in Refs. [49,50]; presumably, the choice in Ref. [12] would make more sense numerically when both m and M are relatively large and/or comparable in size but is not an absolute measure. In due consideration of such a difference, there seems to be no reason to insist on choosing different subtraction points for the two present isospin sectors where the difference in the value of the subtraction constants is practically the pion mass, $M_{\Sigma_c} - M_{\Lambda_c} \sim m_\pi$. In addition, as stated in our anatomy study above, the $I = 1$ *bound* or resonant state should come visibly higher than the one for the $I = 0$ sector. So it should make sense to adopt the subtraction point for $I = 0$ also in the $I = 1$ sector *as viewed from* the DN channel. In addition, such a minor change will not disturb the approximate crossing symmetry as promoted in Ref. [12]. In fact, when we made this increase in μ: $\mu_{I=1} \to \mu_{I=0}$, then the $I = 1$ channel resonance position goes up from 2680 to ~2750 MeV, the new value being closer to the one produced in our cut-off scheme. The sensitivity of the $I = 1$ resonance pole position to a relatively small shift in the subtraction point had been least anticipated and thus is a little surprising.

Before finishing this section, devoted to our study of the two-body input for the in-medium calculation, we note that the Belle Collaboration has recently measured in this energy range an isotriplet of excited charmed baryons decaying into $\Lambda_c^+\pi^-$, $\Lambda_c^+\pi^0$, and $\Lambda_c^+\pi^+$ [51]. It is interpreted as a new charmed baryon, the $\Sigma_c(2800)$, having a width of around 60 MeV, measured with more than 50% error. This baryon has been tentatively identified with a d-wave resonance to conform to quark-model predictions [52], although the expected width $\Gamma \sim 15$ MeV [53] is smaller than the observed one. Actually, the fits performed in Ref. [51] were not too sensitive to varying the signal parametrization using s-wave or p-wave Breit-Wigner functions, hence this resonance could still qualify as an s-wave meson-baryon molecule of the type found in the present work around the same energy and having a width of ~40 MeV, which is compatible with the experimental one.

V. DN COUPLED-CHANNELS EQUATION IN NUCLEAR MATTER

The first obvious medium effect to be included in the DN coupled-channels equations is the Pauli blocking on the intermediate nucleon states. This is a particularly important one in the vicinity of a dynamically generated resonance, as was explicitly shown in Ref. [28] for the $\Lambda(1405)$ in the context of $\bar{K}N$ scattering. Here, intermediate nucleons need more energy to access states that are not occupied, so the resonance is generated at higher energies, moving from below to above the $\bar{K}N$ threshold. This induces, in turn, strong changes in the $\bar{K}N$ amplitude near the threshold. Namely, the threshold behavior of the amplitude changes from repulsive in free space to attractive in the medium upon including Pauli corrections. When this attraction felt by the $\bar{K}$ is now fed into the solution of the in-medium amplitude, the resonance moves back to lower energies. This thus necessitates a self-consistent calculation, which was done in Ref. [31], where the resonance was found practically at the same location as in free space. This behavior was confirmed in Ref. [32], which incorporated also the self-energy of the pions present in the coupled-channels problem along with the baryon binding.

Based on the above observation on the in-medium $\bar{K}$ behavior, we also consider those medium effects in our present study on the properties of the D meson following the approach presented in Ref. [32]. For this purpose, we have only to incorporate them in the meson and baryon propagators of the loop function G, which is then denoted as $\tilde{G}$. Note that here we have not included baryon binding energies altogether due to our lack of knowledge of the charmed baryon mean-field potentials.

The effects of Pauli blocking are simply included by replacing the free nucleon propagator by the in-medium one,

$$G_N(p_0, \vec{p}) = \frac{1 - n(\vec{p})}{p_0 - E_N(\vec{p}) + i\varepsilon} + \frac{n(\vec{p})}{p_0 - E_N(\vec{p}) - i\varepsilon}, \qquad (14)$$

where $n(\vec{p})$ is the nucleon occupation with value 1 (0) for nucleons below (above) the Fermi momentum and $E_N(\vec{p})$ is the nucleon energy. For the D (and π) mesons we incorporate the corresponding self-energy (dressing) in the propagator

$$D_D(q_0, \vec{q}) = \frac{1}{q_0^2 - \vec{q}^{\,2} - m_D^2 - \Pi_D(q_0, \vec{q}, \rho)}, \qquad (15)$$

which is done, in practice, through the corresponding Lehmann representation:

$$D_D(q_0,\vec{q}\,) = \int_0^\infty d\omega \frac{S_D(\omega,\vec{q}\,)}{q_0-\omega+i\varepsilon} - \int_0^\infty d\omega \frac{S_{\bar{D}}(\omega,\vec{q}\,)}{q_0+\omega-i\varepsilon}, \tag{16}$$

where $S_{D(\bar{D})}(q_0,\vec{q}\,)$ is the spectral function of the $D(\bar{D})$ meson. In free space it is simply

$$S_{D(\bar{D})}(q_0,\vec{q}\,) = \Theta(q_0)\delta\big(q_0^2-\vec{q}^{\,2}-m_D^2\big), \tag{17}$$

where $\Theta(q_0)$ is the Heaviside step function. In the nuclear medium the spectral function becomes

$$\begin{aligned} S_D(q_0,\vec{q}\,) &= -\frac{1}{\pi}\mathrm{Im}D_D(q_0,\vec{q}\,) \\ &= -\frac{1}{\pi}\frac{\mathrm{Im}\Pi_D(q_0,\vec{q}\,)}{\big|q_0^2-\vec{q}^{\,2}-m_D^2-\Pi_D(q_0,\vec{q}\,)\big|^2}. \end{aligned} \tag{18}$$

With these medium modifications the propagator loop function $\tilde{G}$ reads:

$$\begin{aligned} \tilde{G}_{DN}(P_0,\vec{P}\,) = & \int_{|\vec{q}|<q_{\rm max}^{\rm lab}} \frac{d^3q}{(2\pi)^3}\frac{M_N}{E_N(\vec{P}-\vec{q}\,)} \\ & \times\Bigg[\int_0^\infty d\omega S_D(\omega,\vec{q}\,) \\ & \times \frac{1-n(\vec{P}-\vec{q}\,)}{P_0-\omega-E_N(\vec{P}-\vec{q}\,)+i\varepsilon} \\ & + \int_0^\infty d\omega S_{\bar{D}}(\omega,\vec{q}\,) \\ & \times \frac{n(\vec{P}-\vec{q}\,)}{P_0+\omega-E_N(\vec{P}-\vec{q}\,)-i\varepsilon}\Bigg], \end{aligned} \tag{19}$$

for DN states and

$$\begin{aligned} \tilde{G}_{\pi Y_c}(P_0,\vec{P}\,) = & \int_{|\vec{q}|<q_{\rm max}^{\rm lab}} \frac{d^3q}{(2\pi)^3}\frac{M_{Y_c}}{E_{Y_c}(\vec{P}-\vec{q}\,)}\int_0^\infty d\omega S_\pi(\omega,\vec{q}\,) \\ & \times \frac{1}{P_0-\omega-E_{Y_c}(\vec{P}-\vec{q}\,)+i\varepsilon}, \end{aligned} \tag{20}$$

for $\pi\Lambda_c$ or $\pi\Sigma_c$ states, where $P=(P_0,\vec{P}\,)$ is the total two-particle four-momentum and $\vec{q}$ is the meson momentum in the nuclear matter rest frame.

For $\eta(\eta')Y_c$, $K\Xi_c(\Xi_c')$, and D_sY states, the corresponding meson lines are not dressed by self-energy insertions, viz. we use the loop integral for free space. The reasons are as follows: (i) the coefficients coupling the $\eta(\eta')Y_c$ channels with the DN channel are small, as shown in Tables I and II, and (ii) containing an $\bar{s}$-quark, the K couples weakly to nucleons and its spectral function may be approximated by the free-space one, viz. by a delta function. Note that this last prescription applies also to $S_{\bar{D}}(\omega,\vec{q})$ in Eq. (19). As for the spectral function of the D_s^+ meson appearing in the in-medium D_sY channels, it has been shown [7] that, in addition to the quasi-particle peak, it presents a lower energy mode associated with an exotic resonance predicted around 75 MeV below the D_s^+N threshold [12]. Therefore, with large coupling coefficients for transitions $DN \leftrightarrow D_sY$, as seen in Tables I and II, one may eventually have to solve an extended in-medium self-consistent coupled-channels problem combining the $C=1$, $S=0(DN)$ and $C=1$, $S=1(D_sN)$ sectors. This arduous task will be relegated to a future work.

Now the in-medium amplitudes $\tilde{T}$ are obtained by directly solving the coupled-channels Eq. (6) with the medium modified loop function $\tilde{G}$ discussed above or from the equivalent Eq. (11), where, formally, the medium correction appears as the second term on the right hand side.

The in-medium D self-energy is finally obtained by integrating $\tilde{T}_{DN}$ over the nucleon Fermi sea as

$$\begin{aligned} \Pi_D(q_0,\vec{q}\,) &= 2\int\frac{d^3p}{(2\pi)^3}n(\vec{p}\,)\big[\tilde{T}_{Dn}(P_0,\vec{P}\,)+\tilde{T}_{Dp}(P_0,\vec{P}\,)\big] \\ &= \int\frac{d^3p}{(2\pi)^3}n(\vec{p}\,)\big[\tilde{T}^{(I=0)}(P_0,\vec{P}\,) \\ &\quad +3\tilde{T}^{(I=1)}(P_0,\vec{P}\,)\big], \end{aligned} \tag{21}$$

where $P_0=q_0+E(\vec{p}\,)$ and $\vec{P}=\vec{q}+\vec{p}$ are the total energy and momentum of the DN pair in the nuclear matter rest frame and the values $(q_0,\vec{q}\,)$ stand for the energy and momentum of the D meson also in this frame. Recall that $\Pi_D(q_0,\vec{q})$ must be determined self-consistently because it is obtained from the in-medium amplitude $\tilde{T}_{DN}$, which contains the DN loop function $\tilde{G}_{DN}$, and this last quantity itself is a function of $\Pi_D(q_0,\vec{q}\,)$.

VI. RESULTS AND DISCUSSION

Before starting our discussion on the various medium effects, let us first explore what is the mass shift of the D meson in isospin symmetric nuclear matter for various simple approximations. We define the effective mass as $m_D^*=\omega_{qp}(\vec{q}=0)$, where the quasi-particle energy, ω_{qp}, is the solution to the equation

$$\omega_{qp}(\vec{q}\,)^2=\vec{q}^{\,2}+m_D^2+\mathrm{Re}\,\Pi_D(\omega_{qp}(\vec{q}\,),\vec{q}\,). \tag{22}$$

The mass shift can then be approximated by the D-meson optical potential at zero momentum:

$$m_D^*-m_D\simeq U_D(0)\equiv\frac{\mathrm{Re}\,\Pi_D(m_D^*,\vec{q}=0)}{2m_D}. \tag{23}$$

First, when one adopts a mean-field type of approximation and replaces $\tilde{T}_{DN}$ by the diagonal V_{DN} in Eq. (21), the mass shift is found as $\sim$−60 MeV for model A that considers $\Sigma_{DN}\neq 0$ or −43 MeV for model B, viz. in the absence of the Σ_{DN} term. This amount of attraction is comparable to what was obtained in QCDSR and NMFA models as discussed in Sec. II [1–3,5]. This may be regarded as the calibration of our diagonal DN interaction with respect to those earlier works.

Next we go one step to replace V_{DN} by the free-space T_{DN}. Due to the presence of resonances generated, the latter quantity is obviously distinct from the former. So we expect a better description than the first one by using T_{DN} in Eq. (21), a procedure that is referred to as the $T\rho_N$ approximation to the D-meson self-energy. With this the D meson feels a repulsion of $\sim$25 MeV, rather than attraction, indicating the importance of a non-perturbative treatment of the problem

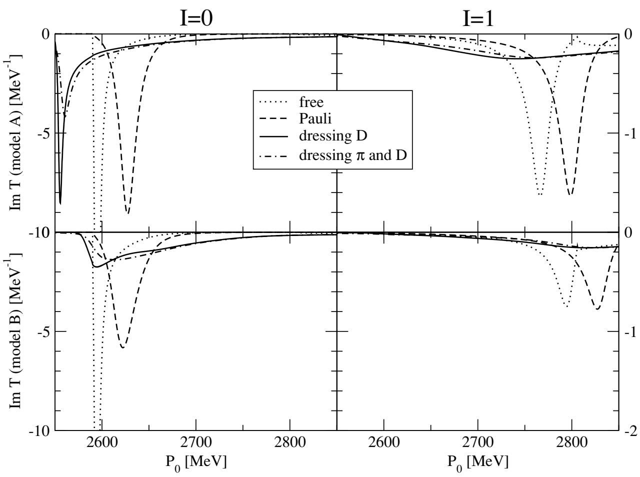

FIG. 5. Imaginary part of the in-medium DN amplitude for total momentum $P = 0$ as a function of the total energy P_0 for $I = 0$ (left panels) and $I = 1$ (right panels), obtained in the cut-off scheme using model A (upper panels) and model B (lower panels).

once resonances are present. As discussed later, in the present model this drastic change is caused mainly by the isospin one $\Sigma_c(2770)$ resonance which lies below but close to the DN threshold.

Then we go further to consider the medium effects, Pauli blocking, and the self-consistent inclusion of the D-meson self-energy in the coupled-channels equations. When Pauli blocking alone is considered, the amount of repulsion increases up to 40 MeV, whereas when meson dressings are incorporated in addition, the repulsive mass shift goes down to 5 MeV. Note, however, that the actual shift of the D-meson mass is eventually determined by the position of the quasi-particle peak in the spectral function, which has a complex structure determined both by the real and the imaginary parts of the self-energy. We will come back to this point toward the end of this section.

Being convinced of the importance of a proper treatment of in-medium effects, we now discuss more explicitly the changes induced in the DN amplitude and, as a result, on the D-meson self-energy and its spectral density. In Fig. 5 we display the imaginary part of the $I = 0$ (left panels) and $I = 1$ (right panels) DN amplitude at normal nuclear matter density, $\rho_0 = 0.17$ fm^{-3}, as a function of the total energy P_0 and total momentum $\vec{P} = 0$. The results of model A and B are presented in the two upper and two lower panels, respectively. With Pauli blocking (dashed lines), the resonances are produced at higher energies than in the free amplitudes (dotted lines), in exact analogy to the behavior of the $\bar{K}N$ in-medium amplitudes described at the beginning of this section. When one adds the D-meson dressing, the $I = 0$ and $I = 1$ resonances move down (solid lines), even below their corresponding free space location. This effect is especially pronounced for model A, which contains a nonvanishing Σ_{DN} term. In particular, the $\Lambda_c(2593)$ appears about 50 MeV lower in energy than in free space. The reason for this additional attraction when the D self-energy is included self-consistently is that the DN amplitudes develop strengths at much lower energies than their free-space thresholds, starting actually at the threshold for the $\pi\Lambda_c$ states, 2422 MeV. This enhances the phase space for intermediate states, inducing effectively a strong attraction in the coupled-channels equations. The result is the lowering of resonance positions below their free-space counterparts. These in-medium resonances will be denoted from now on as $\tilde{\Lambda}_c(2593)$ and $\tilde{\Sigma}_c(2770)$. In particular, the mass of the $\tilde{\Lambda}_c(2593)$ is lowered by about 50 MeV in nuclear matter at normal density. Note that the width of the in-medium $\tilde{\Lambda}_c(2593)$ is not zero even if it now appears below the free space $I = 0$ coupled-channels threshold, $\pi\Sigma_c$. This is due to the process $\tilde{\Lambda}_c(2593)N \to \pi N\Lambda_c$, which opens up as soon as the in-medium D-meson self-energy is incorporated. A similar argument holds for explaining the much larger width of the $I = 1$ resonance, which can also decay through nucleon absorption processes of the type $\tilde{\Sigma}_c(2770)N \to \pi N\Lambda_c, \pi N\Sigma_c$. When the pions are also dressed (dot-dashed lines in the same figure), the tendency does not change much, even if two-nucleon absorption channels, $\tilde{\Lambda}_c(2593)NN \to \Lambda_c NN$, $\Sigma_c NN$, are now possible through the coupling of the pion to particle-hole configurations. This is in contrast to what is observed for $\bar{K}N$ dynamics [32] and also to what is found in Ref. [6] for in-medium D mesons. The reason is that, in the present model, the interaction $DN \to \pi Y_c$ is reduced by the factor $\kappa_c \simeq 1/4$ originating from the t-channel exchange of a charmed vector meson.

The results for models A and B are qualitatively similar. The only noticeable difference is that, due to the absence of attractive scalar-isoscalar Σ_{DN} term, model B produces in-medium resonances at higher energies. Correspondingly, their widths are larger due to the increased decaying phase space.

To illustrate the effect of each approximation on the sign and strength of the in-medium DN amplitude, we display its real part in Fig. 6 for the same cases as in Fig. 5 but focusing on the energy region close to the free-space DN threshold. The bare interaction V, represented by the thin solid line, is attractive both in the $I = 0$ and $I = 1$ channels. The dotted

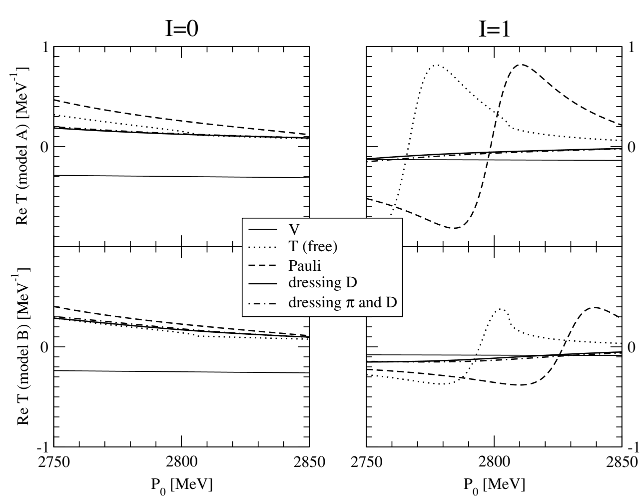

FIG. 6. Real part of the in-medium DN amplitude for total momentum $\vec{P} = 0$ as a function of total energy P_0 for $I = 0$ (left panels) and $I = 1$ (right panels) obtained in the cut-off scheme with model A (upper panels) and model B (lower panels). The selected range of energies covers values around the DN threshold.

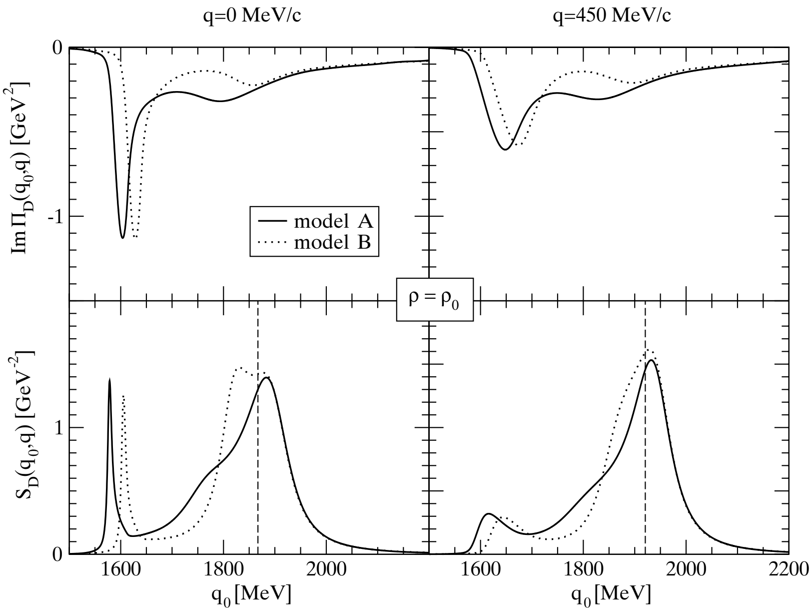

FIG. 7. Imaginary part of the D-meson self-energy (upper panels) and corresponding D-meson spectral function (lower panels) as functions of energy for two values of the meson momentum, $q = 0$ and $q = 450$ MeV at normal nuclear matter density. Results are shown for model A (solid lines) and model B (dotted lines). Thin vertical lines display the location of the D-meson pole in free space.

line represents the free DN amplitude, for which we see that the repulsive effect induced by the $I = 0$ $\tilde{\Lambda}_c(2593)$ is still visible at energies around 2800 MeV. This effect is more pronounced for the $I = 1$ channels because the $\tilde{\Sigma}_c(2770)$ resonance is just below the DN threshold. When Pauli blocking effects are included (dashed lines), we observe the same qualitative behavior except that the amplitudes are shifted to higher energies. The fully self-consistent amplitudes, both with pion dressing (dot-dashed lines) or without (solid lines) show similar features in the $I = 0$ channels as in the other approximations. But in the $I = 1$ channel a drastic dilution of the $\tilde{\Sigma}_c(2770)$ resonance ends up producing a mildly attractive interaction that partly compensates the repulsion found in the $I = 0$ sector.

Finally, in Figs. 7 and 8 we present our results for the imaginary part of the D-meson self-energy (upper panels), as well as the corresponding spectral function (lower panels), as functions of the meson energy q_0 for nuclear matter densities $\rho = \rho_0$ and $2\rho_0$, respectively. We show results for two values of the meson momentum, $q = 0$ (left panels) and $q = 450$ MeV (right panels) and for models A (solid lines) and B (dotted lines). The self-energy presents two peaks, the more pronounced one at lower energy is built up from $\tilde{\Lambda}_c N^{-1}$ configurations, whereas the peak at higher energy is due to the coupling of the D meson to $\tilde{\Sigma}_c N^{-1}$ states.

Each peak of the imaginary part of the D-meson self-energy has a direct association with the structure observed in the spectral function at a slightly lower energy, as seen in the lower panels. The narrower peak is the one associated with $\tilde{\Lambda}_c N^{-1}$ configurations and, for model A, it lies about 50 MeV below what would have been the corresponding location if the $\Lambda_c(2593)$ had not been modified by medium effects. The

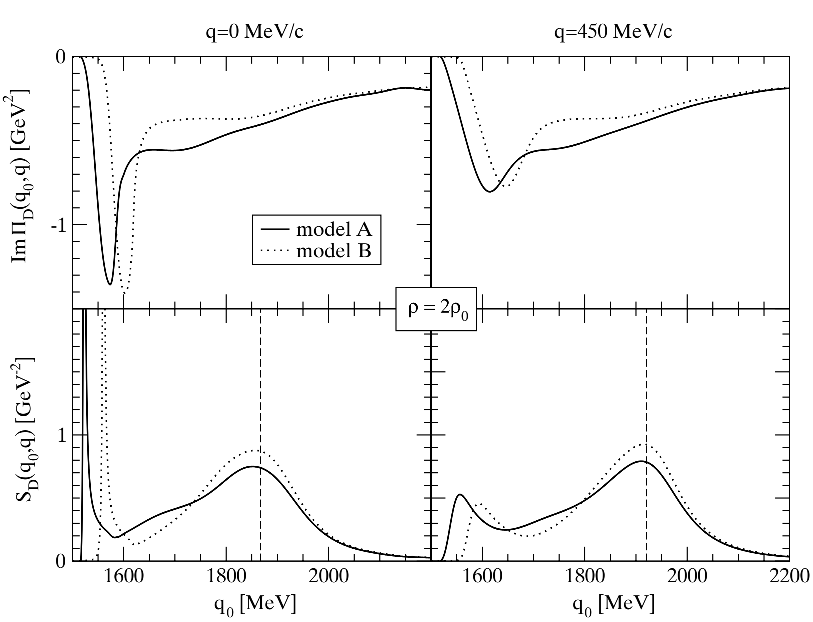

FIG. 8. The same as described in the legend to Fig. 7 for nuclear matter density $\rho = 2\rho_0$.

in-medium attraction of this lower energy mode is more moderate for model B. By inspecting Fig. 8, we also observe that the amount of attraction of these $\tilde{\Lambda}_c N^{-1}$ configurations increases by an additional 50 MeV or slightly more when going to twice nuclear matter density. The structure in the spectral function connected to the bump in the self-energy associated with the $I = 1$ resonance is very faint. It stands out only in model B and at normal nuclear matter density because of its proximity to the D-meson quasi-particle peak.

The location of the quasi-particle peak is mainly driven by the value of quasi-particle energy [Eq. (22)]. At normal nuclear matter density, the quasi-particle energy was found to be 5 MeV above the D-meson mass. However, as commented at the beginning of this section, the actual quasi-particle peak of the spectral function can appear slightly shifted due the energy dependence of the imaginary part of the D-meson self-energy. At normal nuclear matter density and for both models A and B, we find the quasi-particle peak about 20 MeV above the D-meson pole position in free space (denoted by thin vertical lines in the figures). At twice nuclear matter density, the quasi-particle peak is shifted mildly below the free D-meson pole in model A, while staying practically at the same position in model B.

Our results are qualitatively similar to those found in Ref. [6] but differ considerably from those in Ref. [7], especially at $2\rho_0$ where the latter reference finds 60 MeV repulsion for the quasi-particle peak of the in-medium D meson. This is due to the fact that the $I = 1$ resonance is found at a much lower energy than in the present work, influencing differently the DN interaction around threshold, which becomes substantially more repulsive. Correspondingly, the contribution of the $I = 1$ DN amplitude will affect the D-meson spectral function at lower energies. In fact, due to the almost degeneracy between the $I = 0$ and $I = 1$ charmed resonances found in the model of Ref. [12], the lower mode of the D-meson spectral function in Ref. [7] is dominated by $\tilde{\Sigma}_c N^{-1}$ configurations (recall the weight factor 3 for $I = 1$ contributions with respect to $I = 0$ ones).

VII. CONCLUSION

We now summarize our present work. First, we critically reviewed several effective hadronic approaches used to investigate the properties of the D meson in nuclear matter. This then lead us to adopt a recent model by Hofmann and Lutz [12] but with a few important modifications: (i) some simplifications in the form of the interactions which have turned out to be more consistent on reduction from t-channel vector-meson exchanges to a zero-range Tomozawa-Weinberg (T-W) form; (ii) introduction of a supplementary scalar-isoscalar interaction in the diagonal DN channel, which we call the Σ_{DN} term, apparently prevalent both in the QCD sum rule (QCDSR) and mean-field (NMFA) approaches to the problem; and (iii) switching to a conventional momentum cut-off regularization that was found to be more consistent than the dimensional method in view of its application to meson-baryon scattering in nuclear medium.

In free space, the coupled-channels equations resulting from the meson-baryon interactions thus obtained were regularized to reproduce the position and width of the $\Lambda_c(2593)$ resonance in the $I = 0$ DN channel. In the $I = 1$ channel, the same interactions were found to generate a wide resonance, which we term $\Sigma_c(2770)$.

In nuclear matter, the DN diagonal element of the interaction was tested in the lowest-order approximation to the D-meson self-energy based on the simplest mean-field picture, thus with an equation of the type Eq. (3) for D. A D mass reduction of $\sim$60 MeV was found. This is consistent with the consequence in QCDSR [1,2] and some NMFA results [3–5]. If this consistency can be regarded as important, that would support, at least in part, the introduction of the Σ_{DN} in our scheme, which we have taken to be quite conservative. It would be interesting to develop models that could provide a more precise value for this term.

Our in-medium study finds that, once the fully self-consistent coupled-channels equations are solved, including Pauli blocking and meson dressing effects, the situation changes quite drastically. Namely, at normal nuclear matter density the quasi D-meson peak in the D-meson spectral function is found about 20 MeV higher than the corresponding free-space pole position. This appears roughly independent of whether the Σ_{DN} term is present. The primary cause is found in the fact that the $\tilde{\Sigma}_c(2770)$ resonance gets extremely broadened due to the medium effects. As the nuclear matter density increases, the upward shifting of the D-meson quasiparticle peak slows down and can even reverse if $\Sigma_{DN} \neq 0$. So it may be worth exploring how this trend continues at even higher densities. The $\tilde{\Lambda}_c(2593)$ resonance remains narrow and lowers its position by almost 50 MeV at normal nuclear matter density when the Σ_{DN} term is retained in the diagonal DN channel. This tendency persists when the matter density is doubled. It may be appropriate to stress that, unlike in the case of the in-medium $\bar{K}$, the role of the intermediate pion dressing has been found to be of minor importance here. This makes a marked difference from the result reported in Ref. [6], where the $\tilde{\Lambda}_c(2593)$ appears to be washed out by the effect of the pion dressing.

In the Introduction, we stated that one of the primary motivations for studying the behavior of the D meson in nuclear matter is an attempt to understand (even partially) the reduction of the J/Ψ charmonium production observed in the ultra-relativistic heavy-ion reactions, and so on. We are fully aware that the present work is just a first step toward that goal based on the effective hadronic picture. To be more realistic, one of the principal aspects that we will need to investigate is the implementation of the finite temperature effect. As mentioned above, the finite nuclear density makes the $\tilde{\Lambda}_c(2593)$ resonance to move to sensibly lower energies mantaining its narrow width. So it should be important to study if it may survive the thermal agitation of the order of $\sim$100 MeV. In such a case, the excited charmonia such as $\chi_{c\ell}(1P)$ $(\ell = 1, 2)$, could decay strongly through this in-medium resonance, thereby reducing the usual supply of J/Ψ mesons coming from their radiative decay [54]. In addition to a realistic implementation of the temperature effect, one should also consider the nuclear mean-field binding of ground-state charmed baryons, a more extensive study of the scalar attraction characterized by the Σ_{DN} term, inclusion of reaction channels with charmed vector mesons, etc., even within the context of the effective hadronic picture.

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