

Open-charm systems in cold nuclear matter

M.F.M. Lutz^{a,*}, C.L. Korpa^b

^a *Gesellschaft für Schwerionenforschung (GSI), Planck Str. 1, 64291 Darmstadt, Germany*

^b *Department of Theoretical Physics, University of Pecs, Ifjusag u. 6, 7624 Pecs, Hungary*

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Abstract

We study the spectral distributions of charmed meson with $J^P = 0^-$ quantum numbers in cold nuclear matter applying a self-consistent and covariant many-body approach established previously for the nuclear dynamics of kaons. At leading orders the computation requires as input the free-space two-body scattering amplitudes only. Our results are based on the s-wave meson–nucleon amplitudes obtained recently in terms of a coupled-channel approach. The amplitudes are characterized by the presence of many resonances in part so far not observed. This gives rise to an intriguing dynamics of charmed mesons in nuclear matter. At nuclear saturation density we predict a pronounced two-mode structure of the D^+ mesons with a main branch pushed up by about 32 MeV. The lower branch reflects the coupling to two resonance-hole states that are almost degenerate. For the D^- we obtain a single mode pushed up by about 18 MeV relative to the vacuum mode. Most spectacular are the results for the D_s^+ meson. The presence of an exotic resonance-hole state gives rise to a rather broad and strongly momentum dependent spectral distribution.

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1. Introduction

It has been suggested to explore the possible mass modifications of D mesons with the FAIR project at GSI [1]. It is planned to study such effects by detecting D mesons in heavy-ion reactions but also in antiproton-nucleus reactions. The two approaches are complementary, with the second focusing on the properties of D mesons in cold nuclear matter. Thus it is of importance to develop a thorough theoretical understanding of how the D mesons may change their properties in nuclear matter.

At present there are only few works published on the properties of non-strange D mesons [2–5]. To the best of our knowledge, there is no attempt to predict the properties of the strange D_s^\pm mesons in nuclear matter so far. Except for the work by Tolos et al. [5] mean field parameterizations were considered. The model by Tsushima et al. [2] arrived at attractive mass shifts for the D^+ and D^- mesons of about 60 MeV at saturation den-

sity. A different mean field ansatz by Sibirtsev et al. [3] that considers also vector mean field contributions suggests an attractive mass shift of about 140 MeV for the D^+ meson, but a small repulsive mass shift of about 20 MeV for the D^- meson. A QCD sum rule analysis of Hayashigaki [4] predicted an attractive mass shift of about 50 MeV for the D^+ meson. These developments resemble to some extent the first attempts to predict the properties of kaons and antikaons in nuclear matter (see e.g. [6–8]), which also assumed a mean field type behavior for the mass shifts in nuclear matter. However, by now it is well established that for the antikaon such an ansatz is not valid due to complicated many-body dynamics induced by the presence of resonance-hole states (see e.g. [9,10]). Thus it may be crucial to study also the properties of D mesons in a more microscopic manner.

An important step towards a more realistic approach to the D^+ meson in nuclear matter was taken by Tolos et al. [5], who did not insist on a mean field approach exploring the possible influence of the $\Lambda_c(2594)$ resonance on the nuclear dynamics of the D^+ meson. The key ingredient of a realistic description are the $D^+N \rightarrow D^+N$ scattering amplitudes, that determine the D^+ meson self energy at least for dilute nuclear mat-

* Corresponding author.

E-mail address: m.lutz@gsi.de (M.F.M. Lutz).

ter [8,13]. In turn the results of [5] depend decisively on their particular form of the coupled-channel amplitudes that took into account the $\pi \Sigma_c$, DN , $\eta \Lambda_c$ channels for $I = 0$ and $\pi \Lambda_c$, $\pi \Sigma_c$, DN , $\eta \Sigma_c$ channels for $I = 1$. The $\Lambda_c(2594)$ resonance which carries $J^P = \frac{1}{2}^-$ quantum numbers was dynamically generated in [5] as suggested first in [11]. The interaction strengths were derived relying on a SU(3) symmetry for the u , d and c quarks. As a consequence of their amplitudes an attractive mass shift of about 10 MeV together with a significant broadening of the D^+ at nuclear saturation density was predicted. Recently a more complete computation [12] that incorporated the additional channels $K \Xi_c$, $K \Xi'_c$, $D_s \Lambda$ for $I = 0$ and $K \Xi_c$, $K \Xi'_c$, $D_s \Sigma$ for $I = 1$ arrived at scattering amplitudes that differ significantly from those of [5]. The interaction was saturated by a t-channel vector meson exchange. Comparing the different interactions [5,12] the work by Tolos et al. severely overestimates the charm-exchange channels. In [12] those channels are suppressed by a kinematical factor $m_\rho^2/m_D^2 \sim 0.2$. As an important new result of [12] the $I = 0$ amplitude shows two resonance structures around 2.6 GeV. A narrow state that couples dominantly to the DN , $D_s \Lambda$ channels and a broader state that is interpreted as a chiral excitation of the open-charm sextet $\frac{1}{2}^+$ ground states [11]. The narrow state should be identified with the observed state $\Lambda_c(2594)$. The broader one, which couples strongly to the $\pi \Sigma_c$ channels, awaits experimental confirmation and couples very weakly to the DN channel. It is the analogue of the $\Lambda(1405)$, which is a chiral excitation of the baryon octet $\frac{1}{2}^+$ ground states (see e.g. [14–17]). In contrast to [5] the isospin one amplitude of [12] reflects a narrow resonance of mass 2.62 GeV which couples dominantly to the DN , $D_s \Sigma$ channels. This additional resonance will affect the properties of D^+ mesons in nuclear matter significantly.

The purpose of this Letter is threefold. First we derive the properties of the D^+ meson in nuclear matter based on the improved understanding of the D^+N scattering amplitudes of [12]. Second, we present the mass shift of the D^- as predicted by the scattering amplitudes of [12]. Thirdly, we present for the first time predictions on how the strange D_s^\pm mesons change their properties in a dense nuclear environment. Again the scattering amplitudes as obtained in [12] are used. The s-wave $D_s^\pm N \rightarrow D_s^\pm N$ scattering amplitudes are characterized by exotic resonances at 2.89 GeV and 2.78 GeV. The computations are based on the self consistent and covariant many-body approach developed in [9,10]. It is important to perform such computations in a self consistent manner since the feedback of an altered meson spectral function on the resonance structure proved to be a decisive many-body affect [9,10,18].

2. Self consistent and covariant nuclear dynamics for charmed mesons

In this section we recall the self consistent and relativistic many-body framework required for the evaluation of the pseudo-scalar meson propagation in nuclear matter [10]. The key ingredient are the vacuum on-shell meson–nucleon scatter-

ing amplitudes

$$\begin{aligned} \langle D^j(\bar{q})N(\bar{p})|T|D^i(q)N(p) \rangle \\ = (2\pi)^4 \delta^4(q+p-\bar{q}-\bar{p}) \\ \times \bar{u}(\bar{p}) T_{DN \rightarrow DN}^{ij}(\bar{q}, \bar{p}; q, p) u(p), \end{aligned} \quad (1)$$

where $\delta^4(\dots)$ guarantees energy–momentum conservation and $u(p)$ is the nucleon isospin-doublet spinor. Here is $D = (D_+, D_0)$, (\bar{D}_0, D_-) , D_s^+ or D_s^- depending on which system we study. The scattering amplitudes are decomposed into their isospin subsystems. The reactions $D_s^\pm N \rightarrow D_s^\pm N$ carry $I = 1/2$ inheriting the isospin of the nucleons. The amplitudes involving the isospin doublets (D_+, D_0) or (\bar{D}_0, D_-) are decomposed into isospin zero and one components

$$\begin{aligned} T_{DN \rightarrow DN}^{ij}(\bar{q}, \bar{p}; q, p) &= T_{DN \rightarrow DN}^{(0)}(\bar{k}, k; w) P_{(I=0)}^{ij} \\ &\quad + T_{DN \rightarrow DN}^{(1)}(\bar{k}, k; w) P_{(I=1)}^{ij}, \\ P_{(I=0)}^{ij} &= \frac{1}{4}(\delta^{ij} 1 + (\vec{\tau})^{ij} \vec{\tau}), \\ P_{(I=1)}^{ij} &= \frac{1}{4}(3\delta^{ij} 1 - (\vec{\tau})^{ij} \vec{\tau}), \end{aligned} \quad (2)$$

where q, p, \bar{q}, \bar{p} are the initial and final meson and nucleon 4-momenta and

$$w = p + q = \bar{p} + \bar{q}, \quad k = \frac{1}{2}(p - q), \quad \bar{k} = \frac{1}{2}(\bar{p} - \bar{q}). \quad (3)$$

The in-medium meson–nucleon scattering amplitude is determined by the coupled-channel Bethe–Salpeter equation:

$$\begin{aligned} \mathcal{T}(\bar{k}, k; w) &= \mathcal{K}(\bar{k}, k; w) \\ &\quad + \int \frac{d^4 l}{(2\pi)^4} \mathcal{K}(\bar{k}, l; w) \mathcal{G}(l; w) \mathcal{T}(l, k; w), \end{aligned} \quad (4)$$

where the in-medium scattering amplitude $\mathcal{T}(\bar{k}, k; w, u)$ and the two-particle propagator $\mathcal{G}(l; w, u)$ depend on the 4-velocity u_μ characterizing the nuclear matter frame. For nuclear matter moving with a velocity \vec{v} one has

$$u_\mu = \left(\frac{1}{\sqrt{1 - \vec{v}^2/c^2}}, \frac{\vec{v}/c}{\sqrt{1 - \vec{v}^2/c^2}} \right), \quad u^2 = 1. \quad (5)$$

We emphasize that (4) is properly defined from a Feynman diagrammatic point of view even in the case where the in-medium scattering process is no longer well defined due to a broad meson spectral function.

To be consistent with the dynamics derived in [12] all channels must be considered in (4) that were incorporated in [12]. In this work we exclusively study the effect of an in-medium modified two-particle propagator, i.e. we identify the in-medium interaction kernel \mathcal{K} in (4) with its vacuum limit to be taken from [12]. We restrict ourselves further to the in-medium change of the D mesons only, i.e. all two-particle propagators in (4) that do not involve a D meson are taken to be of the form used in [12]. We write

$$\begin{aligned} \Delta S(p, u) &= 2\pi i \Theta(p \cdot u) \delta(p^2 - m_N^2) (\not{p} + m_N) \\ &\quad \times \Theta(k_F^2 + m_N^2 - (u \cdot p)^2), \end{aligned}$$

$$\begin{aligned}
S(p, u) &= S(p) + \Delta S(p, u), \\
\mathcal{D}(q, u) &= \frac{1}{q^2 - m_D^2 - \Pi(q, u)}, \\
\mathcal{G}_{DN}(l; w, u) &= -iS\left(\frac{1}{2}w + l, u\right)\mathcal{D}\left(\frac{1}{2}w - l, u\right), \quad (6)
\end{aligned}$$

where the Fermi momentum k_F parameterizes the nucleon density ρ with

$$\rho = -2 \text{tr} \gamma_0 \int \frac{d^4 p}{(2\pi)^4} i \Delta S_N(p, u) = \frac{2k_F^3}{3\pi^2 \sqrt{1 - \vec{v}^2/c^2}}. \quad (7)$$

In the rest frame of the bulk with $u_\mu = (1, \vec{0})$ one recovers with (7) the standard result $\rho = 2k_F^3/(3\pi^2)$. In this work we also refrain from including nucleonic correlation effects. The meson self energy $\Pi(q, u)$ is evaluated self consistently in terms of the in-medium scattering amplitudes

$$\begin{aligned}
\Pi(q, u) &= 2 \text{tr} \int \frac{d^4 p}{(2\pi)^4} i \Delta S(p, u) \\
&\quad \times \bar{T}\left(\frac{1}{2}(p - q), \frac{1}{2}(p - q); p + q, u\right), \\
\bar{T} &= \frac{1}{4}T_{DN \rightarrow DN}^{(I=0)} + \frac{3}{4}T_{DN \rightarrow DN}^{(I=1)}, \quad \text{or} \quad \bar{T} = T_{D_s^\pm N \rightarrow D_s^\pm N}^{(I=\frac{1}{2})}. \quad (8)
\end{aligned}$$

In order to solve the self consistent set of equations (4), (6), (8) it is convenient to rewrite the coupled-channel system. Given our assumptions the coupled-channel problem reduces without any further restrictions to a single channel problem:

$$\begin{aligned}
\mathcal{T}_{DN \rightarrow DN}^{(I)} &= T_{DN \rightarrow DN}^{(I)} + T_{DN \rightarrow DN}^{(I)} \cdot (\mathcal{G}_{DN} - G_{DN}) \cdot \mathcal{T}_{DN \rightarrow DN}^{(I)}, \quad (9)
\end{aligned}$$

where we use a compact notation. The self consistent set of equations (6), (8), (9) is now completely determined by the vacuum amplitudes $T_{DN \rightarrow DN}^{(I)}$. Rather than specifying the coupled-channel interaction kernel of [12] it is sufficient to recall the s-wave scattering matrix describing the $DN \rightarrow DN$ process. The one of [12] has the separable form

$$\begin{aligned}
T^{(I)}(w) &= \frac{1}{2} \left(\frac{\psi}{\sqrt{w^2}} + 1 \right) M^{(I)}(\sqrt{s}), \\
f^{(I)}(\sqrt{s}) &= \frac{1}{8\pi\sqrt{s}} \left(\frac{\sqrt{s}}{2} + \frac{m_N^2 - m_D^2}{2\sqrt{s}} + m_N \right) M^{(I)}(\sqrt{s}) \\
&= \frac{1}{2ip_{DN}} (\eta^{(I)}(\sqrt{s}) e^{2i\delta^{(I)}(\sqrt{s})} - 1), \\
\frac{p_{DN}^2}{s} &= \frac{1}{4} \left(1 - \frac{(m_N + m_D)^2}{s} \right) \left(1 - \frac{(m_N - m_D)^2}{s} \right), \quad (10)
\end{aligned}$$

where we recall the parametrization of the amplitudes in terms of phase shift δ and inelasticity η .

Using (10) as input in (9) the in-medium amplitude takes the form

$$\begin{aligned}
\mathcal{T}^{(I)}(w, u) &= \frac{1}{2} \left(\frac{\psi}{\sqrt{w^2}} + 1 \right) \mathcal{M}^{(I)}(w, u), \\
\mathcal{M}^{(I)}(w, u) &= [1 - M^{(I)}(\sqrt{s}) \Delta J(w, u)]^{-1} M^{(I)}(\sqrt{s}), \quad (11)
\end{aligned}$$

where we neglect contributions of higher partial-wave amplitudes. The p-wave and d-wave scattering amplitudes are not determined in [12]. The reduced loop function $\Delta J(w, u)$ acquires the generic form

$$\begin{aligned}
\Delta J(w, u) &= \int \frac{d^4 l}{(2\pi)^4} g(l; w, u) \left(m_N + \frac{l \cdot w}{\sqrt{w^2}} \right), \\
g(l; w, u) &= 2\pi \Theta(l \cdot u) \delta(l^2 - m_N^2) \\
&\quad \times \frac{\Theta(k_F^2 + m_N^2 - (u \cdot l)^2)}{(w - l)^2 - m_D^2 - \Pi(w - l, u)} \\
&\quad - \frac{i}{l^2 - m_N^2 + i\epsilon} \frac{1}{(w - l)^2 - m_D^2 - \Pi(w - l, u)} \\
&\quad + \frac{i}{l^2 - m_N^2 + i\epsilon} \frac{1}{(w - l)^2 - m_D^2 + i\epsilon}. \quad (12)
\end{aligned}$$

We observe that the loop function $\Delta J(w, u)$ is a scalar under Lorentz transformation and therefore depends only on w^2 and $w \cdot u$. Thus it can be evaluated in any convenient frame without loss of information. In practice we perform the loop integration in the rest frame of nuclear matter with $u_\mu = (1, \vec{0})$. It is straightforward to perform the energy and azimuthal angle integration in (12). The energy integration of the last two terms in (12) is performed by closing the complex contour in the lower complex half plane. One picks up two contributions: first the nucleon pole leading to $l_0 = \sqrt{m_N^2 + \vec{l}^2}$ in (12) and second the meson pole at typically $w_0 - l_0 = -\sqrt{m_D^2 + (\vec{l} - \vec{w})^2}$. Here we neglect the meson pole contribution. All together one is left with a two-dimensional integral which must be evaluated numerically. In our simulation we restrict $|\vec{l}| < 800$ MeV.

The meson self energy follows

$$\begin{aligned}
\Pi(q, u) &= -2 \int_0^{k_F} \frac{d^3 p}{(2\pi)^3} \left(\frac{m_N}{E_p} + \frac{w_0^2 - \vec{w}^2 - q \cdot w}{E_p \sqrt{w^2}} \right) \bar{\mathcal{M}}(w, u), \\
\bar{\mathcal{M}}(w, u) &= \frac{1}{4} \mathcal{M}^{(I=0)}(w, u) + \frac{3}{4} \mathcal{M}^{(I=1)}(w, u), \quad (13)
\end{aligned}$$

where we used $u_\mu = (1, \vec{0})$, $q_\mu = (\omega, \vec{q})$ and $w_\mu = (\omega + E_p, \vec{q} + \vec{p})$ with $E_p = (m_N^2 + \vec{p}^2)^{1/2}$. With (11), (12) and (13) a self consistent set of equations is defined. It is solved numerically by iteration. First one determines the leading meson self energy $\Pi(\omega, \vec{q})$ by (13) with $\mathcal{M}(w, u) = M(\sqrt{s})$. That leads via the loop functions (12) and the in-medium Bethe–Salpeter equation (11) to medium modified scattering amplitudes $\mathcal{M}(w_0, \vec{w})$. The latter are used to determine the meson self energy of the next iteration. This procedure typically converges after 3 to 4 iterations. The manifest covariant form of the self energy and scattering amplitudes are recovered with $\Pi(q^2, \omega) = \Pi(q^2, q \cdot u)$ and $\mathcal{M}(w^2, w_0) = \mathcal{M}(w^2, w \cdot u)$ in a straightforward manner if considered as functions of q^2 , ω and w^2 , w_0 respectively.

3. Results

We discuss the free-space scattering amplitudes derived in [12] together with their implications for the in-medium properties of the D mesons. The results are based on the self-consistent scheme [10] recalled in the previous section. The s -wave scattering amplitudes are shown in Fig. 1. The spectral functions of the D mesons are shown in Fig. 2 for meson momenta 0, 200 and 400 MeV and nuclear densities 0.17 fm^{-3} and 0.34 fm^{-3} .

The amplitudes reflect the presence of various resonances. Only the D^-N s -wave scattering process is not influenced by a resonance. In this case the amplitude is characterized to a large extent by the scattering length

$$a_{D^-N}^{(I=0)} \simeq -0.16 \text{ fm}, \quad a_{D^-N}^{(I=1)} \simeq -0.26 \text{ fm}. \quad (14)$$

The latter can be read of Fig. 1, being defined as the threshold value of the scattering amplitude $f(\sqrt{s})$ introduced in (10).

Given the isospin averaged scattering length the mass shift of a D^- meson in nuclear matter is fully determined by the low-density theorem [8,13]. It holds

$$\begin{aligned} \Delta m_{D^-} &= -\frac{\pi}{2} \left(\frac{1}{m_N} + \frac{1}{m_D} \right) (a_{D^-N}^{(I=0)} + 3a_{D^-N}^{(I=1)}) \rho + \mathcal{O}(\rho^{4/3}) \\ &\simeq 17 \text{ MeV} \frac{\rho}{\rho_0}. \end{aligned} \quad (15)$$

The low-density mass shift of 17 MeV is quite close to the self consistent result shown in Fig. 2. Self-consistency leads to a mass shift of 18 MeV at saturation density and 38 MeV at twice saturation density.

The positively charged D^+ meson interacts with a nucleon in a more complicated manner. The s -wave scattering amplitude

may be characterized roughly in terms of a resonance-pole term and a background term

$$M(\sqrt{s}) \simeq -\frac{|g|^2}{\sqrt{s} - M_R + i\Gamma_R/2} + b. \quad (16)$$

The isospin-zero amplitude couples strongly to the $\Lambda_c(2594)$ resonance, the isospin-one amplitude to a so far unobserved $\Sigma_c(2620)$ resonance:

$$\begin{aligned} a_{D^+N}^{(I=0)} &\simeq -0.43 \text{ fm}, & |g_{\Lambda_c(2594)}^{(DN)}| &\simeq 6.6, \\ a_{D^+N}^{(I=1)} &\simeq -0.41 \text{ fm}, & |g_{\Sigma_c(2620)}^{(DN)}| &\simeq 5.8. \end{aligned} \quad (17)$$

It is noteworthy that the imaginary part of the scattering length (17) is negligible. This is a consequence of the largely suppressed t -channel exchange of charm discussed before.

A self-consistent computation for the D^+ mass distribution is performed. Fig. 2 demonstrates that the spectral function has a two-mode structure, which is a consequence of important resonance-hole contributions. Note that the masses of the $\Lambda_c(2594)$ and $\Sigma_c(2620)$ are almost degenerate. At saturation density the main mode is pushed up by about 32 MeV as compared to the free-space meson. A self consistent computation where the isospin one amplitude is neglected results in a much reduced repulsion of about 6 MeV at saturation density only. The repulsive shifts are an immediate consequence of the negative scattering lengths (17) feeded into the low-density theorem (15). A scale that characterizes in part the breakdown of the leading order result for the main mode is estimated by

$$\mu_R^2 = \frac{m_N}{m_R} (m_D^2 - (m_R - m_N)^2) \simeq (490 \text{ MeV})^2, \quad (18)$$

where we identify the resonance mass with $m_R \simeq 2.62 \text{ GeV}$ [9]. For $k_F \gtrsim \mu_R$ the low-density result can be trusted no longer.

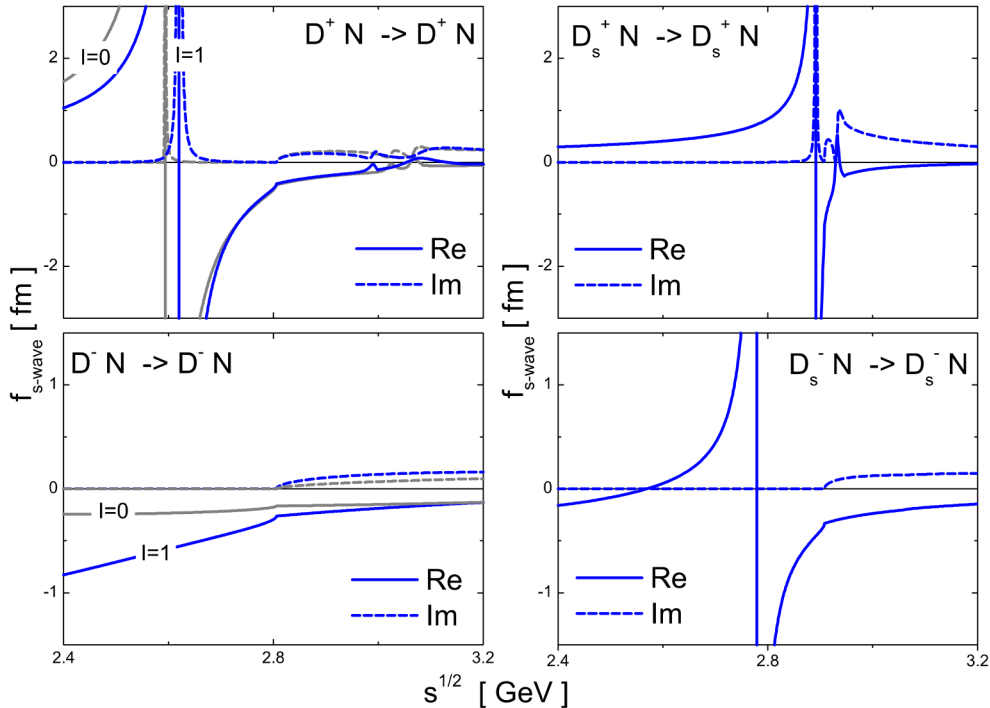


Fig. 1. s -wave scattering amplitudes of the D mesons off the nucleon. The amplitudes are taken from [12]. The normalization of the amplitudes is specified in (10).

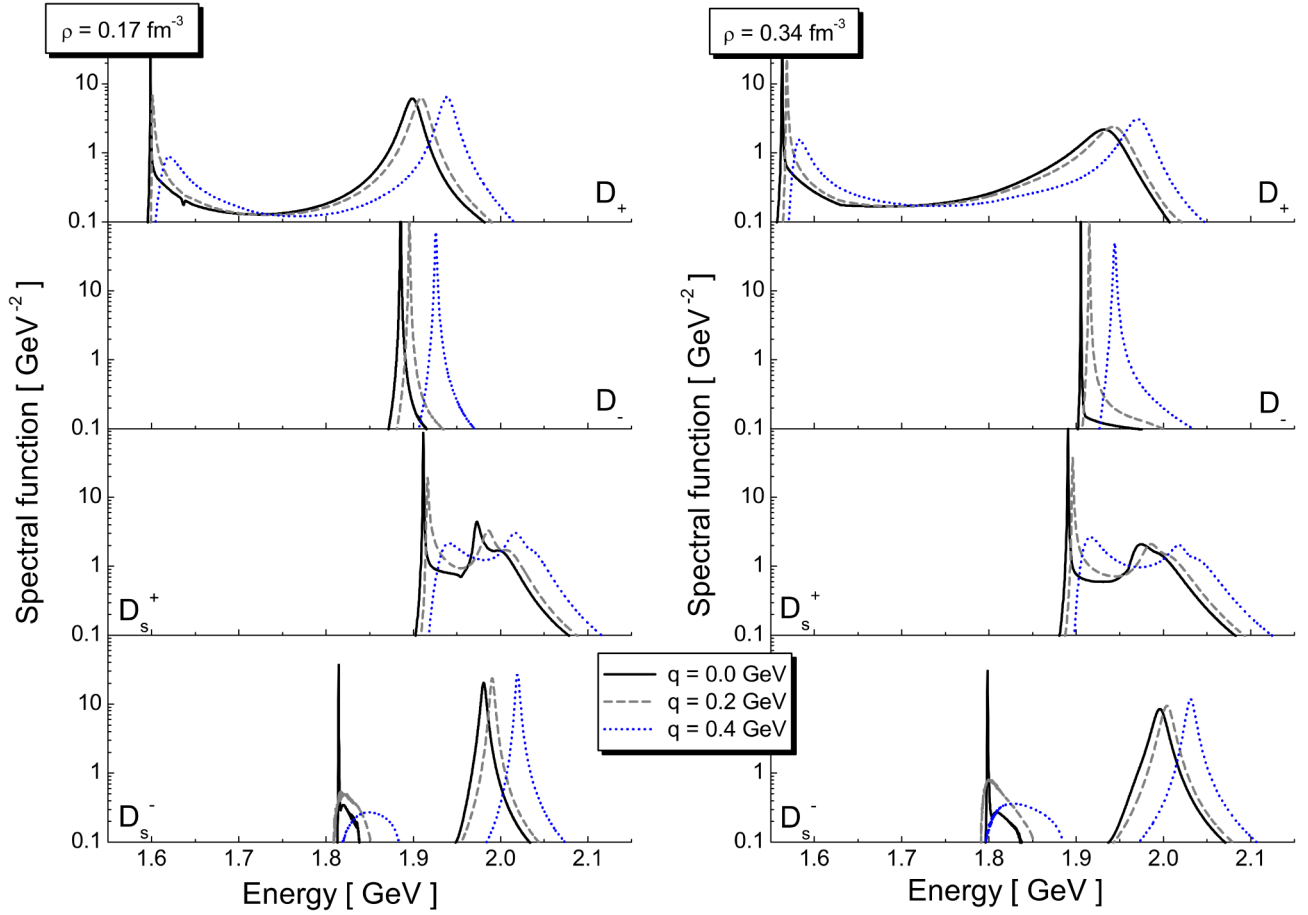


Fig. 2. Spectral distributions of the D^\pm and D_s^\pm mesons based on the scattering amplitudes of Fig. 1. Results are shown for meson momenta 0, 200 and 400 MeV and nuclear densities 0.17 fm^{-3} and 0.34 fm^{-3} . The self consistent many-body approach of [10] was applied.

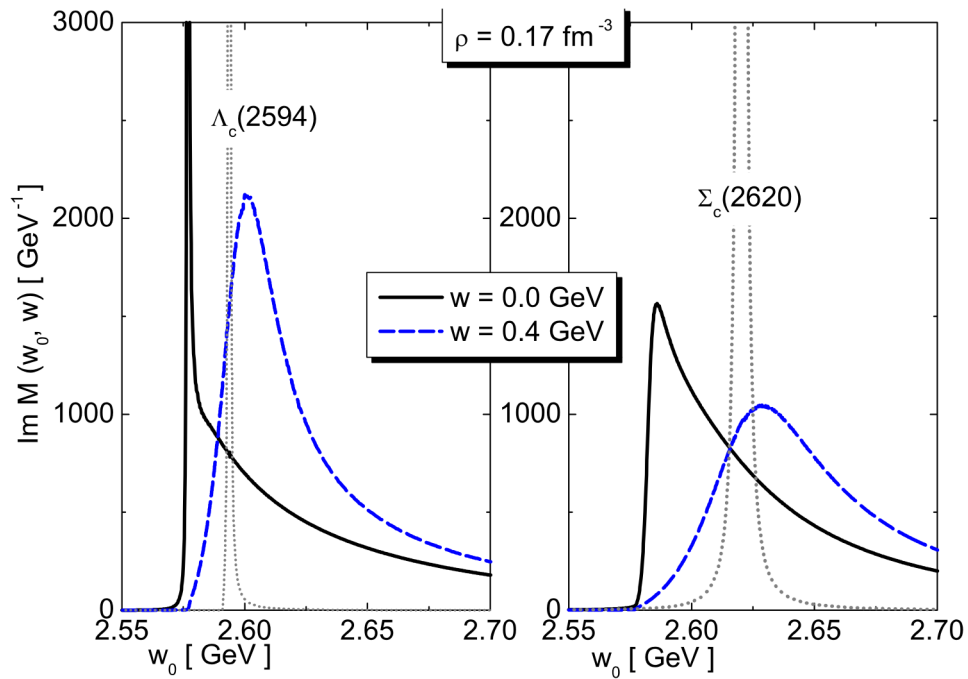


Fig. 3. Imaginary part of the isospin zero (l.h.p.) and isospin one (r.h.p.) D^+ -nucleon scattering amplitude at saturation density as compared to the free-space case (dotted lines). The amplitudes are shown for two values of the resonance three-momentum $w = 0 \text{ GeV}$ and $w = 0.4 \text{ GeV}$.

The breakdown scale is significantly larger for the D^+ meson as compared to the one of the K^- meson. In the latter case the breakdown scale ~ 140 MeV is determined by the $\Lambda(1405)$ resonance [9]. It is stressed that even for a repulsive shift of the main D^+ mode, due to the resonance-hole state one may still expect that the production of D^+ mesons is enhanced in heavy-ion reactions as compared to nucleon–nucleon collisions. Our results for the D^+ meson differ from the previous study [5] significantly. This is a consequence of the quite different interaction used in the two computations. In particular the work [5] did not predict the isospin one resonance $\Sigma_c(2620)$. The latter dominates the resonance-hole component in the spectral distribution of the present work. In Fig. 3 the spectral functions of the $\Lambda_c(2594)$ and $\Sigma_c(2620)$ resonances are shown at saturation density as compared to their free-space distributions. We observe small attractive shifts in their mass distributions and significant broadening, at least once the resonances move relative to the matter bulk. The shifts are the result of a subtle balance of the repulsive Pauli blocking effect and attraction implied by the coupling of the D^+ meson to resonance-hole modes.

We turn to the mesons with non-zero strangeness. The study of [12] predicts the existence of a coupled-channel molecule in the $(D_s^- N, \bar{D}\Lambda, \bar{D}\Sigma)$ system. The latter carries exotic quantum numbers that cannot be arranged by three quarks only. Exotic s-wave states with $C = -1$ were discussed first by Gignoux, Silvestre-Brac and Richard [19] and later by Lipkin [20]. The binding of about 190 MeV for the $N_{s\bar{c}}(2780)$ state predicted in [12] awaits experimental confirmation. If confirmed the s-wave $D_s^- N \rightarrow D_s^- N$ scattering amplitude must show a prominent pole structure at subthreshold energies. The amplitude in Fig. 1 may be characterized with

$$a_{D_s^- N} \simeq -0.33 \text{ fm}, \quad \left| g_{N_{s\bar{c}}(2780)}^{(D_s^- N)} \right| \simeq 3.1. \quad (19)$$

The presence of the $N_{s\bar{c}}(2780)$ leads to a well separated two-mode structure of the D_s^- spectral distribution. The main mode is pushed up by less than 10 MeV at nuclear saturation density.

Recently, it was pointed out that coupled-channel dynamics predicts attraction also for the $(K\Lambda_c, D_s^+ N, K\Sigma_c)$ system [11,12]. An exotic state $N_{\bar{c}s}(2892)$ about 75 MeV below

the $D_s^+ N$ threshold is predicted. From the $D_s^+ N \rightarrow D_s^+ N$ scattering amplitude of Fig. 1 we extract the values

$$a_{D_s^+ N} \simeq -1.21 + i0.01 \text{ fm}, \quad \left| g_{N_{\bar{c}s}(2892)}^{(D_s^+ N)} \right| \simeq 2.8. \quad (20)$$

The in-medium spectral distribution of the D_s^+ derived in the self consistent many body approach is most striking. The two modes expected from the possible existence of the $N_{\bar{c}s}(2892)$ state are almost merged into one broad structure, in particular at intermediate meson momenta 400 MeV. The results are quite analogous to the spectral distribution of the K^- where the $\Lambda(1405)$ nucleon-hole state gives rise to a broad distribution [9,10]. Clearly this result is an immediate consequence of the small binding energy of the $N_{\bar{c}s}(2892)$ state.

We close with the quest for more detailed studies of the D -meson nucleon scattering processes to more reliably predict the in-medium properties of the D mesons.

References

- [1] Conceptual design report, An international accelerator facility for beams of ions and antiprotons, GSI 2001.
- [2] K. Tsushima, et al., Phys. Rev. C 59 (1999) 2824.
- [3] A. Sibirtsev, K. Tsushima, A.W. Thomas, Eur. Phys. J. A 6 (1999) 351.
- [4] A. Hayashigaki, Phys. Lett. B 487 (2000) 96.
- [5] L. Tolos, J. Schaffner-Bielich, A. Mishra, Phys. Rev. C 70 (2004) 025203.
- [6] D.B. Kaplan, A.E. Nelson, Phys. Lett. B 175 (1986) 57.
- [7] M. Lutz, A. Steiner, W. Weise, Phys. Lett. B 278 (1992) 29.
- [8] M. Lutz, A. Steiner, W. Weise, Nucl. Phys. A 574 (1994) 755.
- [9] M. Lutz, Phys. Lett. B 426 (1998) 12.
- [10] M.F.M. Lutz, C.L. Korpa, Nucl. Phys. A 700 (2002) 309.
- [11] M.F.M. Lutz, E.E. Kolomeitsev, Nucl. Phys. A 730 (2004) 110.
- [12] J. Hofmann, M.F.M. Lutz, hep-ph/0507071.
- [13] C.D. Dover, J. Hüfner, R.H. Lemmer, Ann. Phys. 66 (1971) 248.
- [14] H.W. Wyld, Phys. Rev. 155 (1967) 1649.
- [15] R.H. Dalitz, T.C. Wong, G. Rajasekaran, Phys. Rev. 153 (1967) 1617.
- [16] C. García-Recio, M.F.M. Lutz, J. Nieves, Phys. Lett. B 582 (2004) 49.
- [17] E.E. Kolomeitsev, M.F.M. Lutz, Phys. Lett. B 585 (2004) 243.
- [18] C.L. Korpa, R. Malfliet, Phys. Rev. C 52 (1995) 2756.
- [19] G. Gignoux, B. Silvestre-Brac, J.M. Richard, Phys. Lett. B 193 (1987) 323.
- [20] H.J. Lipkin, Phys. Lett. B 195 (1987) 484.