Resonances in chiral unitary approaches

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Abstract. The extension of chiral theories to the description of resonances, via the incorporation of unitarity in coupled channels, has provided us with a new theoretical perspective on the nature of some of the observed excited hadrons. In this contribution some of the early achievements in the field of baryonic resonances are reviewed, the recent evidence of the two-pole nature of the Λ(1405) is discussed and results on charmed baryon resonances are presented.

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1 Introduction

The chiral effective lagrangians describing the interactions of mesons and baryons respect the basic symmetries of the fundamental QCD theory. The low energy processes are well described within the so called chiral perturbation theory (χPT), which is an expansion of these lagrangians according to the number of derivatives of the meson and baryon fields. In recent years, the introduction of unitarity constraints in coupled channels has led to extensions of the chiral theories that can be applied at much higher energies. An interesting consequence of these non-perturbative methods has been the prediction of many low-lying baryon and meson resonances that are generated dynamically through the multiple scattering of their meson or baryon components. The nature of these states is intrinsically different from the usual q̅q or qqq structure of meson and baryons in the sense that they are not preexistent states that remain in the large \( N_c \) limit where the multiple scattering is suppressed.

In this contribution we give an overview of some of the last findings in the field of dynamically generated baryon resonances paying an special attention to the recently confirmed two-pole nature of the Λ(1405). We will also present some predictions of charm resonances obtained from an extension of the model to SU(4).

2 Chiral unitary model

A compilation of chiral unitary methods can be found in Ref. [1]. The search for dynamically generated resonances proceeds by first constructing the meson-baryon coupled states from the octet of ground state positive-parity baryons (\( B \)) and the octet pseudoscalar mesons (\( \Phi \)) for a given strangeness channel. The interaction between mesons and baryons is built from the effective chiral lagrangian which, at lowest order in momentum, reads

\[
L^{(B)}_1 = \bar{B}i\gamma^\mu \frac{1}{4f^2} [\Phi \partial_\mu \Phi - \partial_\mu \Phi \Phi] B - B(\Phi \partial_\mu \Phi - \partial_\mu \Phi \Phi) \tag{1}
\]

and produces the following driving interaction in \( S \)-wave

\[
V_{ij} = -C_{ij} \frac{1}{4f^2} (2\sqrt{s} - M_i - M_j) \sqrt{\frac{M_i + E_i}{2M_i}} \sqrt{\frac{M_j + E_j}{2M_j}} , \tag{2}
\]

where \( i, j \) are indices standing for the different meson-baryon states in the coupled channel problem, the constants \( C_{ij} \) are SU(3) coefficients encoded in the chiral lagrangian and \( f \) is the meson decay constant. The scattering matrix amplitudes between the various meson-baryon states are then obtained by solving the coupled channel equation

\[
T_{ij} = V_{ij} + V_{id} G_i T_{ij} , \tag{3}
\]

where the \( V_{id} \) and \( T_{ij} \) amplitudes are taken on-shell. This is a particular case of the N/D unitarization method when the unphysical cuts are ignored [23]. Under these conditions the diagonal matrix \( G_i \) is simply built from the convolution of a meson and a baryon propagator and can be regularized either by a cut-off \( (q_{\text{max}}^2) \) or, alternatively, by dimensional regularization depending on a subtraction constant \( (a_i) \) coming from a subtracted dispersion relation.
3 The \( \Lambda(1405) \) and its two-pole nature

The history of the \( \Lambda(1405) \) as a dynamical resonance generated from the interaction of meson baryon components in coupled channels is long [4], but it has experienced a boost within the context of unitary extensions of chiral perturbation theory [5, 6, 7, 8, 9, 10], where the lowest order chiral Lagrangian and unitarity in coupled channels generates the \( \Lambda(1405) \) and leads to good agreement with the \( K^-p \) reactions. The models have been fine tuned recently by including the additional terms of the next-to-leading order lagrangian [11, 2, 13, 14]. The pioneer work of Ref. [5] included the channels closest to the \( \Lambda(1405) \), namely \( \pi\Sigma \) and \( KN \) in \( I = 0 \) and \( \pi\Lambda, \pi\Sigma \) and \( KN \) in \( I = 1 \), obtaining a good reproduction of the scattering observables only if the next-to-leading order lagrangian was also considered. In the work of Ref. [8] all the channels that can be built form the lowest lying pseudoscalar mesons and the octet of ground state baryons were employed, namely \( K^0n, \pi^0\Lambda, \pi^0\Sigma^0, \pi^+\Sigma^- \), \( \pi^-\Sigma^+ \), \( \eta\Lambda, \eta\Sigma^0, \eta^+\Xi^- \), \( \eta^0\Xi^0 \) in the case of \( K^-p \) scattering. It was noted that, even if the additional channels included are well above the \( K^-p \) threshold, there are important interferences between the real parts of the amplitudes and these channels, especially the \( \eta\Lambda \) and \( \eta\Sigma^0 \) ones, were extremely important to reproduce the threshold branching ratios taking simply the chiral lagrangian at lowest order. This is clearly illustrated in Fig. 1 were the elastic and inelastic \( K^-p \) cross section are compared with the results of the model of Ref. [8].

![Fig. 1. \( K^-p \) scattering cross sections as functions of the \( K^- \) momentum in the lab frame: with the full basis of physical states (solid line), omitting the \( \eta \) channels (long-dashed line) and with the isospin-basis (short-dashed line). Taken from Ref. [8].](image1)

The surprise, however, came with the realization that there are two poles in the neighborhood of the \( \Lambda(1405) \) both contributing to the final experimental invariant mass distribution [7, 8, 9, 10, 11, 15, 16, 17, 18, 19]. The properties of these two states are quite different, one has a mass around 1390 MeV, a large width of about 130 MeV and couples mostly to \( \pi\Sigma \), while the second one has a mass around 1425 MeV, a narrow width of about 30 MeV and couples mostly to \( KN \). The two states are populated with different weights in different reactions and, hence, their superposition can lead to different distribution shapes. Since the \( \Lambda(1405) \) resonance is always seen from the invariant mass of its only strong \( \pi\Sigma \) decay channel, hopes to see the second pole are tied to having a reaction where the \( \Lambda(1405) \) is formed from the \( KN \) channel. In this sense a calculation of the \( K^-p \to \gamma\Lambda(1405) \) reaction [20], prior to the knowledge of the existence of the two \( \Lambda(1405) \) poles, showed a narrow structure at about 1420 MeV. With the present perspective this is clearly interpreted as the reaction proceeding through the emission of the photon followed by the generation of the resonance from \( K^-p \), thus receiving a large contribution from the second narrower state at higher energy. The recently measured reaction \( K^-p \to \pi^0\pi^0\Sigma^0 \) [21] allows us to test already the two-pole nature of the \( \Lambda(1405) \). This process shows a strong similarity with the reaction \( K^-p \to \gamma\Lambda(1405) \), where the photon is replaced by a \( \pi^0 \).

A model for the reaction \( K^-p \to \pi^0\pi^0\Sigma^0 \) in the energy region of \( p_K = 514 \) to 750 MeV/c, as in the experiment [21], has been studied in Ref. [22]. The mechanisms considered are those in which a \( \pi^0 \) loses the necessary energy to allow the remaining \( \pi^0\Sigma^0 \) pair to be on top of the \( \Lambda(1405) \) resonance. It was found that the dominant contribution was the nucleon pole term, in which the \( \pi^0 \) is emitted from the proton and, consequently, the resonance is initiated by a \( K^-p \) state. The \( \Lambda(1405) \) thus obtained comes mainly from the \( K^-p \to \pi^0\Sigma^0 \) amplitude which, as mentioned above, gives the largest possible weight to the second (narrower) state.

![Fig. 2. The \( \pi^0\Sigma^0 \) invariant mass distribution for three different initial kaon momenta. Taken from Ref. [22].](image2)

In Fig. 2 the \( \pi^0\Sigma^0 \) invariant mass distribution obtained in Ref. [22] for three different energies of the incoming \( K^- \) are compared to the experimental data. Symmetrization of the amplitudes produces a sizable amount of background. At a kaon laboratory momentum of \( p_K = 581 \) MeV/c this background distorts the \( \Lambda(1405) \) shape producing cross section in the lower part of \( M_I \), while at \( p_K = 714 \) MeV/c the strength of this background is shifted toward the higher \( M_I \) region. An ideal situation is found for momenta around 687 MeV/c, where the background sits below the \( \Lambda(1405) \) peak distorting its shape minimally. The peak of the resonance shows up at \( M_I^2 = 2.02 \) GeV$^2$. 
which corresponds to $M_I = 1420$ MeV, larger than the nominal $\Lambda(1405)$, and in agreement with the predictions of Ref. [10] for the location of the peak when the process is dominated by the $f_{KN \to \pi \Sigma}$ amplitude. The apparent width of experiment is about 40 – 45 MeV, but a precise determination would require to remove the background mostly coming from the “wrong” $\pi^0\Sigma^0$ pairs due to the indistinguishability of the two pions. The theoretical analysis of Ref. [22] permits extracting the pure resonant part by not symmetrizing the amplitude, finding a width of $\Gamma \approx 38$ MeV, which is smaller than the nominal $\Lambda(1405)$ width of $50 \pm 2$ MeV [22], obtained from the average of several experiments, and much narrower than the apparent width of about 60 MeV that one sees in the $p \to K^0\pi\Sigma$ experiment [23], which also produces a distribution peaked at 1395 MeV. In order to illustrate the difference between the $\Lambda(1405)$ resonance seen in this latter reaction and in the present one, the two experimental distributions are compared in Fig. 3. We recall that the shape of the $\Lambda(1405)$ in the $p \to K^0\pi\Sigma$ reaction was shown in Ref. [23] to be largely built from the $\pi\Sigma \to \pi\Sigma$ amplitude, which is dominated by the lower, lower energy state. We have thus seen that the $K^- p \to \pi^0\pi^0\Sigma^0$ reaction is particularly suited to study the features of the second pole of the $\Lambda(1405)$ resonance, since it is largely dominated by a mechanism in which a $\pi^0$ is emitted prior to the $K^- p \to \pi^0\Sigma^0$ amplitude, which is the one giving the largest weight to the second narrower state at higher energy. The model of Ref. [22] has proved to be accurate in reproducing both the invariant mass distributions and integrated cross sections seen in a recent experiment [21]. The study of the $K^- p \to \pi^0\pi^0\Sigma^0$ reaction, complemental to the one of Ref. [23] for the $p \to K^0\pi\Sigma$ reaction, has shown that the quite different shapes of the $\Lambda(1405)$ resonance seen in these experiments can be interpreted in favour of the existence of two poles with the corresponding states having the characteristics predicted by the chiral theoretical calculations. Besides demonstrating once more the great predictive power of the chiral unitary theories, this combined study of the two reactions gives the first clear evidence of the two-pole nature of the $\Lambda(1405)$.

Various other reactions testing the two-pole nature of the $\Lambda(1405)$ have been presented at this workshop [20].

4 Extensions within the flavour SU(3) sector

The strangeness $S = 0$ channel was also investigated using the Lipmann–Schwinger equation and coupled channels in Ref. [27,28]. The $N^*(1535)$ resonance was found to be generated dynamically within this approach. Subsequently, work was done in this sector [29,30], and the $N^*(1535)$ resonance, as well as the low-energy scattering observables, were well reproduced, with the exception of the isospin 3/2 channel. Ref. [31] continued and further improved work along these lines by introducing the $\pi N \to \pi NN$ channels, which proved essential in reproducing the isospin 3/2 part of the $\pi N$ amplitude.

The $S = -2$ sector was investigated within a chiral unitary theory in the work of Ref. [32]. By taking the parameters of the model to be of natural size, a pole in the amplitude was found at 1605 + i65 MeV, which coupled strongly to $\pi \Xi$ and $K \Lambda$ states and very mildly to the $K \Sigma$ and $\eta \Xi$ channels. While the width of this resonance might appear as grossly overestimating that of the experimentally observed $\Xi(1620)$ and $\Xi(1690)$ resonances, of around 50 MeV, the calculated $\pi \Xi$ invariant mass distribution shows a smaller apparent width compared to the one obtained at the pole position, due to the fact that the resonance appears just below the threshold of the $K \Lambda$ channel to which it couples very strongly (Flatté effect). The generated resonance was identified with the $\Xi(1620)$, because the $\Xi(1680)$ has branching ratios for $K \Sigma$ to $K \Lambda$ of around 3 and for $\pi \Xi$ to $K \Sigma$ of less than 0.09, values that are incompatible with the properties obtained in the model of Ref. [32]. Therefore, it was possible to assign the unknown spin and parity of this resonance to $J^P = 1/2^-$, which are the values of the states produced with the meson-baryon $S$-wave lagrangians. These findings were confirmed in Ref. [17], where a second $S = -2$ pole that could be identified with the $\Xi(1680)$ resonance was also found.

Recently, a good amount of work has been devoted to the study of $J^P = 3/2^-$ states that can be dynamically generated through the interaction of pseudoscalar 0$^-$ mesons with decuplet 3/2$^+$ baryons [33,34,35,36,37,38,39], and including also the vector meson degrees of freedom within spin-flavor SU(6) symmetry [10].

The $S = -1$ and $I = 0$ sector forms a two-coupled $\pi \Sigma^*$ and $K \Xi^*$ problem that generates a real pole close to the mass of the observed $\Lambda(1520)$. Obtaining a width is only possible if one also includes lower-lying channels formed by a pseudoscalar meson and a baryon of the 1/2$^+$ octet, such as $KN$ and $\pi \Sigma$, interacting in $D$-wave with couplings fitted to experimental amplitudes [34,35]. The cross sections of the model for the $K^- p \to \pi^0\pi^0\Lambda$ and $K^- p \to \pi^+\pi^-\Lambda$ reactions compare very satisfactory with data [41,42]. This has to be considered as a non-trivial accomplishment of the theory since the $KN \to \pi \Sigma^*$ amplitude has not been included in the fit. Other studies related

Fig. 3. Two experimental shapes of $\Lambda(1405)$ resonance. See text for more details. Taken from Ref. [22].
to the $\Lambda(1520)$ comprise the coupling of the resonance to $K^*N$ states \cite{36} and its radiative decays \cite{37}.

In the $S = 0 I = 3/2$ sector, the $\Delta(1700)$ is another of the resonances that are obtained from the interaction of the octet of mesons with the baryon decuplet \cite{38,39}. The coupling of this resonance to its $\pi\Delta$, $K\Sigma^*$ and $\eta\Delta$ components are found to be very different than those obtained if the $\Delta(1700)$ was assumed to belong to a SU(3) decuplet, as suggested by the PDG \cite{23}. Assuming one or the other picture would give rise to very different predictions for cross sections in reactions involving the excitation of the $\Delta(1700)$. Therefore, the agreement of the theory with data for pion and photon induced reactions \cite{29} must again be considered as a success of the chiral unitary models.

5 Dynamical resonances with charm

The physics in the charm $C = 1$ sector bears a strong resemblance with the phenomenology seen in $KN$ dynamics, once the $s$ quark is replaced by a $c$ quark. This is reinforced by an apparent correspondence between the two $I = 0$ $S$-wave $\Lambda(1405)$ and $\Lambda_c(2593)$ resonances. The mass of the former lies in between the thresholds of the $\pi\Sigma$ and $KN$ channels, to which it couples strongly. The $\Lambda_c(2593)$ lies below the $DN$ threshold and just slightly above the $\pi\Sigma_c$ one. It is therefore logical to explore whether the $\Lambda_c(2593)$ can also be dynamically generated from the interaction of $DN$ pairs and its related channels. The study of the $DN$ interaction is also interesting by itself, in connection to the behavior of $D$ mesons in a nuclear medium for the implications it may have on the phenomenon of $J/\Psi$ suppression, which is connected with the formation of a quark-ghon plasma in relativistic heavy-ion collisions, and on the production of open-charm mesons that will be explored in the CBM experiment at the future FAIR facility \cite{43}.

The in-medium modifications of $D$ mesons from a coupled channel perspective was first explored in Ref. \cite{44}, where free space amplitudes were constructed from a set of separable coupled-channel interactions obtained with chirally motivated lagrangians upon replacing the $s$ quark by the $c$ quark, to reproduce the $I = 0 \Lambda_c(2593)$ as a $DN$ $S$-wave hadronic molecule state of binding energy $\approx 200$ MeV with a width of $\approx 3$ MeV, sitting very close to the $\pi\Sigma_c$ threshold. This work represented the first indication that the $\Lambda_c(2593)$ could have a dynamical origin but, by ignoring the strangeness degree of freedom, the role of the $D_s\Lambda$, $D_s\Sigma$, $K\Xi_c$, and $K\Xi'_c$ states was excluded from this model. In addition, both $\pi$ and $K$ Goldstone mesons should be considered on an equal footing within chiral symmetry, while, on the other hand, the $D$ mesons are much heavier and obey heavy-quark symmetry. A blind $s \to c$ replacement breaks both of those symmetries.

A different approach, which respected the proper symmetries, was attempted in Ref. \cite{45}. There, charmed baryon resonances were generated dynamically from the scattering of Goldstone bosons off ground-state $J^P = \frac{3}{2}^+$ charmed baryons. 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, came out to be more than twenty times larger than the experimental width of about 4 MeV. The problem of this model is that, apart from ignoring the $D_s\bar{Y}$ states, it does not account either for the coupling to the $DN$ channel, which contributes strongly to generating the narrow $\Lambda_c(2593)$ according to Ref. \cite{41}.

A substantial improvement came in a recent work \cite{46} in which 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 by a $t$-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 Tomozawa-Weinberg ($T$-$W$) vector type. The model generates a narrow $C = 1$, $I = 0$ resonance that it is identified with the $\Lambda_c(2593)$, together with an almost degenerate $S$-wave resonance at 2620 MeV in the $C = 1$, $I = 1$ channel, which is not seen experimentally. An application of this model to a preliminary study of $D$ and $D_s$ mesons in nuclear matter may be found in Ref. \cite{47}.

The work of Ref. \cite{48} implemented some modifications over the model developed in Ref. \cite{46}. In the first place, the $t$-dependence of the interaction kernel was eliminated and only the leading order zero-range terms were considered. This was motivated from the observation \cite{48} that the $t \to 0$ limit implemented in the energy denominator of the kernel compensated to a good extent the term $k_\rho k_{\rho'}/m^2_\chi$ in the numerator, which was retained in Ref. \cite{46}. The resulting interaction was of the form shown in Eq. 2 but with an additional factor $1/4$ for charm-exchange transitions, accounting approximately for the ratio $(m_N/m_\chi)^2$ between the squared masses of uncharmed and charmed vector mesons, such as the $\rho$ and $D^*$, respectively. Another novel feature of the model of Ref. \cite{48} is that the $T$-$W$ interaction is supplemented with a scalar-isoscalar attraction characterized by a $\Sigma_{DN}$ term, which, due to the large difference between the $c$ and $s$ quark masses, might be more relevant in the charm sector than in the strange one. Mean-field approaches have shown that this term has a tremendous influence on the in-medium $D$ and $\bar{D}$ properties \cite{49,50,51}. Finally, in view of the application of the model to determining the $D$ meson properties in nuclear matter, a cut-off regularization scheme was implemented in Ref. \cite{48}. The value of the cut-off was adjusted to reproduce the position and width of the $\Lambda_c(2593)$ resonance, allowing also for a fine-tuning of the strength of the Goldstone boson decay constant $f$, as in Ref. \cite{46}. Two models were explored, the difference being the inclusion or not of the $\Sigma_{DN}$ term. The corresponding parameters are, model A: $f = 1.15 f_\pi$, $\Sigma = \Sigma_{DN}/f^2_\rho = 0.05$ MeV$^{-1}$, $\Lambda = 727$ MeV, and model B: $f = 1.15 f_\pi$, $\Sigma = \Sigma_{DN}/f^2_\rho = 0$ MeV$^{-1}$, $\Lambda = 787$ MeV, where $\Lambda$ is the ultra-violet cut-off value for the integration in the loop $G_l$.

As seen on the left panel of Fig. 4 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. In the $I = 1$
sector, shown in the right panel of the figure, both models generate a resonance close to 2800 MeV, with a width of around 30 MeV. We note that the Belle Collaboration has recently measured in this energy range an isotriplet of excited charmed baryons decaying into $A^+_1 \pi^-$, $A^{++} \pi^0$ and $A^{++} \pi^+$. 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 predicted width $\Gamma \sim 15$ MeV [54] is substantially smaller than the observed one. Actually, the fits performed in [52] were not too sensitive to varying the signal parameterization using $S$-wave or $P$-wave Breit-Wigner functions, hence this state could still qualify as $S$-wave type and, therefore, according to results of the charm-extended chiral unitary model of Ref. [48], it could be interpreted as being a dynamically generated resonance.

Models have already been extended to the study of $D$-wave charm resonances through the interaction of Goldstone mesons with decuplet baryons [54], or considering the additional role of the interaction of vector mesons with octet and decuplet baryons [54].

We finalize this section by showing the properties of $D$ mesons in a hot nuclear medium, as obtained in Ref. [57], using the above described dynamical model for the $DN$ interaction [48]. The in-medium solution at finite temperature incorporates Pauli blocking effects, mean-field binding on all the baryons involved, and $\pi$ and open-charm meson self-energies in a self-consistent manner.

The obtained evolution of the $D$-meson spectral function with temperature is seen in Fig. 4 for two densities, $\rho_0$ and $2\rho_0$, and two momenta, $q = 0$ MeV/c and $q = 450$ MeV/c, for model A. Two peaks appear in the spectral density at $T = 0$. The lower one corresponds to the $A_cN^{-1}$ excitations, whereas the higher one is the quasi-particle peak mixed with the $\Sigma_cN^{-1}$ mode. The position of the quasi-particle peak lies not too far away from the free mass, a feature which is in strong contrast with the large attractive shifts found in mean-field models [49, 50, 51]. The $A_cN^{-1}$ and $\Sigma_cN^{-1}$ structures dissolve with increasing temperature, while the quasi-particle peak becomes narrower and moves closer to its free value position.

This is due to the fact that the self-energy receives contributions from $DN$ pairs at higher momentum where the interaction is weaker. Similar features were observed in an earlier study [55], which used the three-flavour dynamical model developed in Ref. [44].

The widening of the quasi-particle peak for larger nuclear density may be understood as due to enhancement of collision and absorption processes. The $A_cN^{-1}$ mode moves down in energy with increasing density due to the lowering in the position of the $A_c$ resonance induced by the more attractive $\Sigma_{DN}$ term in model A.

The properties of the $D$ meson in hot nuclear matter were also investigated in the work of Ref. [57], where it was found that its mass develops a mild repulsive shift and that, despite the absence of resonant structures in the $DN$ interaction, the low-density approximation to the repulsive $D$ self-energy was not reliable even at subsaturation densities.

Taking into account these results, one can already state that, while $D^-$-mesic atoms will always be bound by the Coulomb interaction, no strongly bound nuclear states or even bound $D^0$ nuclear systems are expected due to the repulsive in-medium mass shift obtained. In the charm $C = 1$ sector, an experimental observation of bound $D$ nuclear states is ruled out by the moderate attraction and large width found for the $D$ meson. With respect to the implications on $J/\Psi$ suppression in hadronic reactions, we note that the process $J/\Psi \to DD$ cannot proceed spontaneously in free space due to a mass difference of $\approx 650$ MeV. According to the self-consistent models in nuclear matter, the in-medium $D$ mass is seen to increase by about 10–20 MeV whereas the tail of the quasi-particle peak of the $D$ spectral function extends to lower energies due to the thermally spread $Y_cN^{-1}$ configurations. Nevertheless, it is very unlikely that this lower tail extends as far down by more than 600 MeV with sufficient strength to give rise to a direct disappearance of $J/\Psi$'s. Alternatively, $J/\Psi$ suppression in a hadronic environment could also proceed by

![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 with model A (solid lines) and model B (dotted lines). Taken from Ref. [48].](image1.png)

![Fig. 5. The evolution of the $D$ meson spectral function with temperature for $\rho = \rho_0$ (upper panels) and $\rho = 2\rho_0$ (lower panels). Taken from Ref. [57].](image2.png)
cutting its supply from the excited charmonia: $\chi_{cJ}(1P)$ or $\psi'$. These states will be strongly absorbed in the medium by multi-nucleon processes, which will be enhanced by the low-energy tail of the spread out $D$-meson spectral function.

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