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$\Lambda(1520)$ and $\Sigma(1385)$ in the nuclear medium

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Recent studies of the $\Lambda(1520)$ resonance within chiral unitary theory with coupled channels find the resonance as a dynamically generated state from the interaction of the decuplet of baryons and the octet of mesons, essentially a quasibound state of $\pi\Sigma^*(1385)$ in this case, although the coupling of the $\Lambda(1520)$ to the $\bar{K}N$ and $\pi\Sigma$ makes this picture only approximate. The $\pi\Sigma^*(1385)$ decay channel of the $\Lambda(1520)$ is forbidden in free space for the nominal mass of the $\Sigma^*(1385)$, but the coupling of the π to *ph* components in the nuclear medium opens new decay channels of the $\Lambda(1520)$ in the nucleus and produces a much larger width. Together with medium modifications of the $\bar{K}N$ and $\pi\Sigma$ decay channels, the final width of the $\Lambda(1520)$ at nuclear matter density is more than five times bigger than the free one. We perform the calculations by dressing simultaneously the $\Lambda(1520)$ and the $\Sigma^*(1385)$ resonances, finding moderate changes in the mass but substantial ones in the width of both resonances.

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I. INTRODUCTION

The $\Lambda(1520)$ is an intriguing resonance that has captured much attention and is easily produced in K^--induced reactions [1–4] or photon-induced reactions [5]. Two recent different initiatives have brought this resonance again into a focus of attention. First, this resonance appears in experiments that seek to identify the pentaquark Θ^+ [6] (see [7] for a detailed reference of papers on this issue and [8] for a recent review) and is an important reference. In fact, the large background appearing in the tail of the $\Lambda(1520)$ at energies higher than the nominal mass [9] is one of the issues to be clarified when trying to analyze the Θ^+. The issue of the large background has already been addressed in [10] and it is found to be associated to the large coupling of the resonance to the $\pi\Sigma^*(1385)$ channel. More detailed theoretical studies of the $\Lambda(1520)$ photoproduction have been done in [11–13].

The other initiative concerning the $\Lambda(1520)$ has been the study of the interaction of the decuplet of baryons with the octet of pseudoscalar mesons [14,15], which has demonstrated that many of the low lying $3/2^-$ resonances are dynamically generated from the interaction of the coupled channels of these two multiplets. In particular, the $\Lambda(1520)$ appears basically as a quasibound state of the $\pi\Sigma^*(1385)$ system. The small free width of the $\Lambda(1520)$ of $\simeq$15 MeV comes from the decay into $\bar{K}N$ and $\pi\Sigma$, since the decay into $\pi\Sigma^*(1385)$ is forbidden for the nominal mass of the $\Sigma^*(1385)$. Of course, the coupling of the $\Lambda(1520)$ to $\bar{K}N$ and $\pi\Sigma$ makes the picture of the $\Lambda(1520)$ more elaborate, with $\pi\Sigma^*(1385)$ being a very important component but with also sizable admixtures of $\bar{K}N$ and $\pi\Sigma$ [16–18].

The change of resonance properties in the nuclear medium is also a field that captures permanent attention [19–27]. The decay of the $\Lambda(1520)$ in the nuclear medium bears resemblance to that of the $\Delta(1232)$ [28]. The Δ decays into πN and the π gets renormalized in the medium by exciting *ph* and Δh components, as a consequence of which the Δ is renormalized and its pion- (photon-) induced excitation in nuclei incorporates now the mechanisms of pion (photon) absorption in the medium. In the present case, the $\Lambda(1520)$ decay into $\pi\Sigma^*(1385)$, only allowed through the $\Sigma^*(1385)$ width, gets drastically modified when the π is allowed to excite *ph* and Δh components in the nucleus, since automatically the phase space for the decay into $ph\Sigma^*(1385)$ gets tremendously increased. This fact, together with the large coupling of the $\Lambda(1520)$ to the $\pi\Sigma^*(1385)$ channel predicted by the chiral theory, leads to a very large width of the $\Lambda(1520)$ in nuclei. Similar nuclear effects will modify the $\pi\Sigma$ decay channel and the $\bar{K}N$ will be analogously modified when $\bar{K}$ is allowed to excite hyperon-hole excitations. All these channels lead to a considerable increase of the width of the $\Lambda(1520)$ in the nucleus.

The medium corrections of the $\pi\Sigma^*(1385)$ channel already has a precedent in physics in the decay of an ordinary $\Lambda(1115)$ in nuclei. The free $\Lambda(1115)$ decay into πN through weak interactions, mesonic decay, is largely suppressed by Pauli blocking in nuclei. However, the pion can excite *ph* components in the medium, leading to a new $\Lambda(1115)$ decay mode $\Lambda(1115)\to Nph$, or equivalently $\Lambda(1115)N\to NN$, a nonmesonic decay. This new decay channel is far bigger than the mesonic decay in the nucleus and as large as the free one [29–34].

As we shall see, another large source of medium modification of the $\Lambda(1520)$ will come from the medium corrections to the $\pi\Sigma$ and $\bar{K}N$ decay channels. These channels are not present in the interaction of the decuplet of baryons with the octet of mesons with the chiral Lagrangians. They are included as additional channels, with empirical couplings to the other channels and among themselves, that are determined by fits to the $\bar{K}N\to\bar{K}N$ and $\bar{K}N\to\pi\Sigma$ *D*-wave scattering. Altogether, we find a large width of the $\Lambda(1520)$ in the medium, more than a factor of 5 times the free width.

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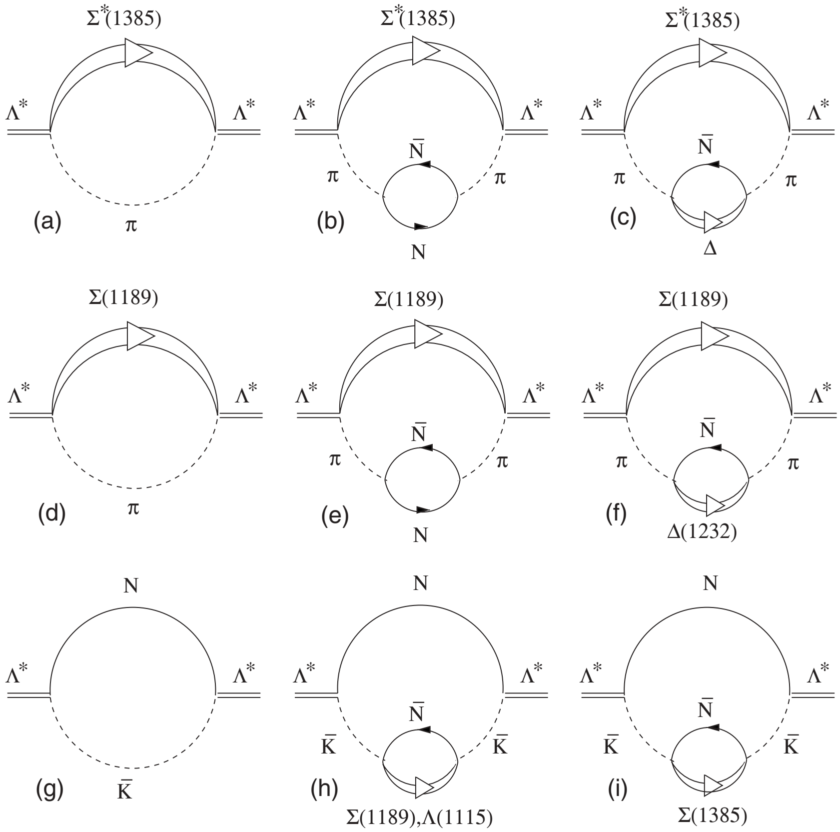

FIG. 1. The renormalization of the $\Lambda(1520)$ in the nuclear medium.

Clear indications that this might be the case can be seen in the analysis of $\Lambda(1520)$ production in heavy-ion reactions [35,36].

Using similar methods we shall also study the modification of the $\Sigma^*(1385)$ in the nuclear medium, by looking at the modification of the $\pi\Sigma$ and $\pi\Lambda$ decay channels, as well as modification of the virtual $\bar{K}N$ channel when the $\bar{K}$ is allowed to excite Λh, Σh, and $\Sigma^* h$ components in the nucleus.

We have organized the paper as follows. In Secs. II and III the model for the $\Lambda(1520)$ and $\Sigma^*(1385)$ self-energies in the nuclear medium is described. The result and discussions are presented in Secs. IV and V. General conclusions are given in Sec. VI.

II. RENORMALIZATION OF THE $\Lambda(1520)$

In this section we discuss the formalism used in the present work for the description of in-medium properties of the $\Lambda(1520)$ hyperon. Here we follow the standard approach where the nuclear medium is described by the noninteracting Fermi sea and the baryonic resonances get modified in the nuclear medium in the dressing procedure by coupling the mesons in the loops to baryon—hole and hyperon—hole excitations.

A. The S-wave decay of the $\Lambda(1520) \to \pi\Sigma^*(1385)$

In the study of Refs. [10,14,16] the $\Lambda(1520)$ is generated dynamically from the $\pi\Sigma^*(1385)$ and $K\Xi^*$ channels interacting in the S wave. In [16] the $\bar{K}N$ and $\pi\Sigma(1189)$ channels in the D wave are added to produce the proper width of the $\Lambda(1520)$. Of all these channels the most important one is the $\pi\Sigma^*(1385)$, but the width into this channel is very small owing to the largely reduced phase space for the $\Lambda(1520)$ decay into $\pi\Sigma^*(1385)$, which is only possible through the tail of the $\Sigma^*(1385)$ when its width is considered. But the relevance of the $\pi\Sigma^*(1385)$ channel and the fact that in the nuclear medium the phase space for decay into this and associated channels becomes very large make the renormalization of this channel very important in studying the $\Lambda(1520)$ in the nuclear medium.

The in-medium renormalization of $\Lambda(1520)$ in the $\pi\Sigma^*(1385)$ channel can be represented by the first three diagrams in Fig. 1. In these diagrams the $\Sigma^*(1385)$ arises as an intermediate state but will be also dressed in its relevant decay channels. Hence, in the present work we address simultaneously the dressing of the $\Sigma^*(1385)$ and $\Lambda(1520)$ in the nuclear matter.

In the following we specify the in-medium propagators of the hyperons $\tilde{D}_Y$ and pions $\tilde{D}_\pi$. The kaon propagation in the nuclear medium will be addressed separately. In terms of a dispersion relation representation we have

$$\tilde{D}_Y(K,\rho) = \int_0^\infty dW \frac{S_Y(W,\boldsymbol{K},\rho)}{K^0 - W + i0^+}, \tag{1}$$

$$\tilde{D}_\pi(k,\rho) = \int_0^\infty d\omega(2\omega) \frac{S_\pi(\omega,\boldsymbol{k},\rho)}{(k^0)^2 - \omega^2 + i0^+}. \tag{2}$$

Here $K(k)$ are the four-momenta, $S_{Y(\pi)}$ are the spectral functions of hyperons (pions), and ρ is the nuclear matter density. Equation (1) accounts for the positive energy part of the fermionic propagator only. For $S_{Y(\pi)}$ we have

$$S_{Y(\pi)} = -\frac{1}{\pi}\mathrm{Im}[\tilde{D}_{Y(\pi)}], \tag{3}$$

where

$$\mathrm{Im}[\tilde{D}_Y(W,\boldsymbol{K},\rho)] = \frac{M_Y}{E_Y(\boldsymbol{K})} \times \frac{\mathrm{Im}[\Sigma_Y(W,\boldsymbol{K},\rho)]}{[W - E_Y(\boldsymbol{K}) - \mathrm{Re}\Sigma_Y(W,\boldsymbol{K},\rho)]^2 + [\mathrm{Im}\Sigma_Y(W,\boldsymbol{K},\rho)]^2}, \tag{4}$$

$$\mathrm{Im}[\tilde{D}_\pi(\omega, \boldsymbol{k}, \rho)] = \frac{\mathrm{Im}[\Pi_\pi(\omega, \boldsymbol{k}, \rho)]}{[\omega^2 - \tilde{\omega}^2(\boldsymbol{k}) - \mathrm{Re}\Pi_\pi(\omega, \boldsymbol{k}, \rho)]^2 + [\mathrm{Im}\Pi_\pi(\omega, \boldsymbol{k}, \rho)]^2}. \tag{5}$$

In Eqs. (4) and (5) $E_Y(\boldsymbol{K}) = \sqrt{\boldsymbol{K}^2 + M_Y^2}$ and $\tilde{\omega}(\boldsymbol{k}) = \sqrt{\boldsymbol{k}^2 + m_\pi^2}$ are the on-mass-shell energies of hyperons and pions, respectively, and the in-medium self-energy Σ_Y is the subject of the present calculations. The P-wave pion polarization operator Π is given by

$$\Pi_\pi(k, \rho) = \left(\frac{D+F}{2f_\pi}\right)^2 \boldsymbol{k}^2 \mathcal{U}(k, \rho) \times \left[1 - \left(\frac{D+F}{2f_\pi}\right)^2 g' \mathcal{U}(k, \rho)\right]^{-1}, \tag{6}$$

where D and F are the axial vector coupling constants and $f_\pi = 93$ MeV is the pion decay constant. One finds $F \simeq 0.51$ and $D \simeq 0.76$ and the axial coupling constant used in the present calculations is $g_A = D + F \simeq 1.27$. Also in Eq. (5) $g' = 0.7$ is the Landau-Migdal parameter [37] and $\mathcal{U}(k, \rho) = \mathcal{U}^d(k, \rho) + \mathcal{U}^c(k, \rho)$ is the Lindhard function including the direct and crossed contributions of p-h and Δ-h excitations with the normalization of the appendix of Ref. [31]. These are the conventional definitions. Latter on we shall modify the formalism to account for the short-range correlations relevant for the in-medium dressing of the $\Sigma^*(1385)$.

The S-wave character of the $\Lambda(1520) \to \pi\Sigma^*(1385)$ decay requires the following transition amplitude:

$$-it_{\Lambda(1520)\to\pi\Sigma^*} = -ig_{\Lambda^*\pi\Sigma^*}, \tag{7}$$

where $g_{\Lambda^*\Sigma^*\pi}$ already accounts for the three charge states in the isospin $I = 0$ channel [see Eq. (8)], in the decay $\Lambda(1520) \to \pi\Sigma^*(1385)$. Alternatively, one can use the isospin decomposition (using the convention for phases from Ref. [16]

$$\Lambda(1520) \to |\pi\Sigma^*(1385); I = 0\rangle = \frac{1}{\sqrt{3}}[|\pi^-\Sigma^{*+}\rangle - |\pi^+\Sigma^{*-}\rangle - |\pi^0\Sigma^{*0}\rangle] \tag{8}$$

and multiply the coupling constant $g_{\Lambda^*\Sigma^*\pi}$ by Clebsch-Gordan coefficients $\pm\sqrt{1/3}$ in the corresponding vertices. The sum of partial decays width into these channels equals the $I = 0$ contribution provided by the single coupling of Eq. (7).

The value of the coupling constant $g_{\Lambda^*\Sigma^*\pi} = 1.57$ was determined in Ref. [16] by means of the residue of the scattering amplitude near the pole position of the $\Lambda(1520)$. A value of $g_{\Lambda^*\Sigma^*\pi} = 1.21$ was used in Ref. [10], where the $\bar{K}N$ and $\pi\Sigma$ channels are not considered. Recent studies that reproduce $\bar{K}N \to \bar{K}N$ and $\bar{K}N \to \pi\Sigma$ D-wave scattering data result in $g_{\Lambda^*\Sigma^*\pi} = 0.89$ [17,18]. The latter value is used in the present work. Actually in Refs. [17,18] there is a discussion of many reactions involving the $\Lambda(1520)$. However, for the purpose of setting the strength of $g_{\Lambda^*\Sigma^*\pi}$ it suffices to quote that this coupling is needed to interpret the $K^-p \to \pi^+\pi^-\Lambda$ reaction [38] using the formalism of Ref. [16] and the empirical width of 15 MeV for the $\Lambda(1520)$. (The $K^-p \to \pi^+\pi^-\Lambda$ cross section is twice as big as that of $K^-p \to \pi^0\pi^0\Lambda$ studied in [16].)

By using Eq. (7) the one-loop contribution to the $\Lambda(1520)$ self-energy from the $\pi\Sigma^*$ intermediate state takes the form

$$-i[\Sigma(P, \rho)]^{\Lambda(1520)}_{\pi\Sigma^*(1385)} = g^2_{\Lambda^*\Sigma^*\pi}\int \frac{d^4k}{(2\pi)^4}\tilde{D}_{\Sigma^*}(P-k, \rho)\tilde{D}_\pi(k, \rho), \tag{9}$$

where $\tilde{D}_{\Sigma^*}$ and $\tilde{D}_\pi$ are in-medium propagators of the $\Sigma^*(1385)$ and pion, respectively. The in-medium renormalization of the $\Sigma^*(1385)$ that enters Eq. (9) will be addressed in Sec. III.

B. The D-wave decay of the $\Lambda^*(1520) \to \pi\Sigma + \bar{K}N$

The conventional D-wave decay of the $\Lambda(1520)$ into the $\bar{K}N$ and $\pi\Sigma$ channels accounts for practically all of the $\Lambda(1520)$ free width [39]. The diagrams responsible for the renormalization of the $\Lambda(1520)$ in these decay channels are shown in Figs. 1(d)–1(i). The pertinent D-wave transition operator is given by

$$-it_{\pi\Sigma\to\Lambda^*(\bar{K}N\to\Lambda^*)} = -ig_{\Lambda^*\pi\Sigma(\Lambda^*\bar{K}N)}(\boldsymbol{S}^\dagger\cdot\boldsymbol{k})(\boldsymbol{\sigma}\cdot\boldsymbol{k}), \tag{10}$$

where $\boldsymbol{k}$ is the pion three momentum and $S_i^\dagger$ is the 2×4 transition operator from spin 1/2 to spin 3/2 fulfilling the relation $S_i S_j^\dagger = 2\delta_{ij}/3 - i\epsilon_{ijk}\sigma_k/3$. Both coupling constants $g_{\Lambda^*\pi\Sigma}$ and $g_{\Lambda^*\bar{K}N}$ are adjusted to reproduce the free decay branches of the $\Lambda(1520)$. By using Eq. (10) the partial decay width of the $\Lambda(1520)$ into the $\bar{K}N$ or $\pi\Sigma$ decay channels can be calculated using the formula

$$\Gamma_{\Lambda^*\to\Sigma(N)+\pi(\bar{K})}(s) = -2\mathrm{Im}[\Sigma(s)]^{\Lambda^*}_{\Sigma\pi(\bar{K}N)} = \frac{1}{3}\left(\frac{g^2_{\Lambda^*\pi\Sigma(\Lambda^*\bar{K}N)}}{2\pi}\right)\frac{M_{\Sigma(N)}}{\sqrt{s}}|\boldsymbol{k}_{\mathrm{c.m.}}|^5, \tag{11}$$

where $|\boldsymbol{k}_{CM}| = \lambda^{1/2}(s, M^2_{\Sigma(N)}, m^2_{\pi(\bar{K})})/(2\sqrt{s})$ with $s = P^2$ and λ is the Kallén function. At the nominal pole position we get $g_{\Lambda^*\pi\Sigma} = 10.75\mathrm{GeV}^{-2}$ and $g_{\Lambda^*\bar{K}N} = 16.01\mathrm{GeV}^{-2}$. In dimensionless units, comparable to $g_{\Lambda^*\Sigma^*\pi}$ in Eq. (7)

$$g_{\Lambda^*\pi\Sigma(\Lambda^*\bar{K}N)}\boldsymbol{k}^2_{\mathrm{c.m.}}/\sqrt{3} \equiv \tilde{g}_{\Lambda^*\pi Y(\Lambda^*\bar{K}N)}; \tag{12}$$

the value of these couplings at the pole position of the $\Lambda^*(1520)$ would be $\tilde{g}_{\Lambda^*\pi\Sigma} = 0.44$ and $\tilde{g}_{\Lambda^*\bar{K}N} = 0.54$, which shows that the coupling $g_{\Lambda^*\Sigma^*\pi} = 0.89$ is still the largest.

1. $\Lambda^*(1520) \to \pi\Sigma(1189)$ *channel*

In the nuclear medium there are peculiarities that force us to consider the $\pi\Sigma$ and $\bar{K}N$ channels separately. The $\pi\Sigma$ channel is a simplest one in the conventional decay of the $\Lambda^*(1520)$ hyperon. The self-energy loop integral in this decay channel reads

$$[\Sigma(P,\rho)]^{\Lambda^*}_{\pi\Sigma} = g^2_{\Lambda^*\pi\Sigma}\frac{i}{3}\int\frac{d^4k}{(2\pi)^4}\boldsymbol{k}^4\tilde{D}_\Sigma(P-k,\rho)\tilde{D}_\pi(k,\rho). \tag{13}$$

The imaginary part of Eq. (13) is meaningful by itself and can be obtained using the Cutkosky rules

$$\begin{aligned}\Sigma(P,\rho) &\to 2i\mathrm{Im}[\Sigma(P,\rho)],\\ \tilde{D}_Y(K,\rho) &\to 2i\mathrm{Im}[\tilde{D}_Y(K,\rho)]\cdot\theta(K^0),\\ \tilde{D}_\pi(k,\rho) &\to 2i\mathrm{Im}[\tilde{D}_\pi(k,\rho)]\cdot\theta(k^0).\end{aligned} \tag{14}$$

From this it is given by

$$\begin{aligned}\mathrm{Im}[\Sigma(P,\rho)]^{\Lambda^*}_{\pi\Sigma} &= \frac{g^2_{\Lambda^*\pi\Sigma}}{3}\int\frac{d^3\boldsymbol{k}}{(2\pi)^3}\frac{\boldsymbol{k}^4 M_\Sigma}{E_\Sigma(\boldsymbol{P}-\boldsymbol{k})}\\ &\quad\times\mathrm{Im}[\tilde{D}_\pi(\omega,\boldsymbol{k},\rho)]\cdot\theta(\omega)|_\omega\\ &= {}_{P_0-E_\Sigma(\boldsymbol{P}-\boldsymbol{k})-V_\Sigma(\rho)},\end{aligned} \tag{15}$$

where θ is the step function and $V_\Sigma(\rho)$ is the binding correction for the $\Sigma(1189)$ hyperon [see Eq. (36)].

2. $\Lambda^*(1520) \to \bar{K}N$ *channel*

The proper treatment of the $\bar{K}N$ channel is a more subtle problem. First, we consider the modification of the antikaon propagator in the nuclear medium. In its particle-antiparticle decomposition, the dispersion relation representation of the $\bar{K}$ propagator is given by

$$\tilde{D}_{\bar{K}}(k^0,\boldsymbol{k};\rho) = \int_0^\infty d\omega\frac{S_{\bar{K}}(\omega,\boldsymbol{k};\rho)}{k^0-\omega+i0^+} - \int_0^\infty d\omega\frac{S_K(\omega,\boldsymbol{k};\rho)}{k^0+\omega-i0^+}, \tag{16}$$

where $S_{\bar{K}(K)} = -\mathrm{Im}\tilde{D}_{\bar{K}(K)}/\pi$ is the spectral function of the $\bar{K}(K)$ meson, which depends on $\bar{K}(K)$ in-medium self-energy $\Pi_{\bar{K}(K)}$ as provided by Eqs. (3) and (5). As is well known, the interactions of $\bar{K}$ and K with the nucleons of the nuclear medium are rather different and, in principle, it is necessary to treat them separately.

In the calculation of the in-medium self-energy we need the in-medium nucleon propagator, which is given by

$$\begin{aligned}\tilde{D}_N(p,\rho) &= \frac{M_N}{E_N(\boldsymbol{p})}\frac{1-n(|\boldsymbol{p}|)}{p^0-E_N(\boldsymbol{p})-V_N(\rho)+i0^+}\\ &\quad+\frac{M_N}{E_N(\boldsymbol{p})}\frac{n(|\boldsymbol{p}|)}{p^0-E_N(\boldsymbol{p})-V_N(\rho)-i0^+},\end{aligned} \tag{17}$$

where $n(|\boldsymbol{p}|) = \theta(k_F - |\boldsymbol{p}|)$ and k_F is the Fermi momentum. The first term in Eq. (17) describes the Pauli-blocked propagation of nucleons and the second term is the hole propagator. The effect of the nucleon binding is accounted for by the mean-field potential $V_N(\rho) \simeq -60\rho/\rho_0$ MeV. The self-energy integral, after integration over the k^0 component, takes the following form:

$$\begin{aligned}[\Sigma(P,\rho)]^{\Lambda^*}_{\bar{K}N} &= g^2_{\Lambda^*\bar{K}N}\frac{i}{3}\int\frac{d^4k}{(2\pi)^4}\boldsymbol{k}^4\tilde{D}_N(P-k,\rho)\tilde{D}_{\bar{K}}(k,\rho)\\ &= -g^2_{\Lambda^*\bar{K}N}\frac{1}{3}\int_0^\infty d\omega\int\frac{d^3\boldsymbol{k}}{(2\pi)^3}\frac{M_N}{E_N(\boldsymbol{P}-\boldsymbol{k})}\\ &\quad\times\frac{\boldsymbol{k}^4[1-n(|\boldsymbol{P}-\boldsymbol{k}|)]S_{\bar{K}}(\omega,\boldsymbol{k};\rho)}{P^0-E_N(\boldsymbol{P}-\boldsymbol{k})-\omega-V_N(\rho)+i0^+}\\ &\quad-g^2_{\Lambda^*\bar{K}N}\frac{1}{3}\int_0^\infty d\omega\int\frac{d^3\boldsymbol{k}}{(2\pi)^3}\frac{M_N}{E_N(\boldsymbol{P}-\boldsymbol{k})}\\ &\quad\times\frac{\boldsymbol{k}^4 n(|\boldsymbol{P}-\boldsymbol{k}|)S_K(\omega,\boldsymbol{k};\rho)}{P^0-E_N(\boldsymbol{P}-\boldsymbol{k})+\omega-V_N(\rho)-i0^+}.\end{aligned} \tag{18}$$

In Eqs. (18) and (13) a static regulating form factor $\tilde{F}(\boldsymbol{k}^2)^2$ is introduced: $\tilde{F}(\boldsymbol{k}^2) = F(\boldsymbol{k}^2)/F(\boldsymbol{k}^2_{\text{on-shell}})$ with $F(\boldsymbol{k}^2) = \Lambda^2/(\Lambda^2+\boldsymbol{k}^2)$ and $|\boldsymbol{k}_{\text{on-shell}}|$ is the on-shell momentum in $\Lambda(1520) \to \bar{K}N(\pi\Sigma)$. This normalization of the form factor is needed to guarantee the free width from Eqs. (13) and (18). The choice of Λ is not arbitrary but is demanded by the study done of the $\Lambda(1520)$ in Refs. [17,18]. In these works, in addition to the $\pi\Sigma^*$, $\bar{K}\Xi^*$ channels considered in the baryon decuplet meson octet chiral Lagrangian [14,15], the channels $\pi\Sigma$ and $\bar{K}N$ in the D wave are included. One has to regularize the loops, and together with the unknown couplings of the $\pi\Sigma$ and $\bar{K}N$ channels to the dominant channel $\pi\Sigma^*$, the regularization scale was fitted to the data. The data fitted were the D-wave amplitudes $\bar{K}N \to \bar{K}N$ and $\bar{K}N \to \pi\Sigma$ in $I = 0$. Dimensional regularization was used in the analysis as in [40,41]. However, given the equivalence of dimensional regularization and cutoff regularization see appendix in [42] for values of the cutoff reasonably larger than the on-shell momentum, the same numerical results could be obtained with the use of a cutoff of about 500 MeV in the D-wave $\bar{K}N$ and $\pi\Sigma$ loops.

Note that the equivalence of the cutoff method and dimensional regularization has been proved in [42] for one loop. For two or more loops a chiral invariant cutoff regularization is more subtle [43,44], but this is a problem that we are not facing here since we only have one-loop diagrams. (The *ph* excitation loop is only a formal way of summing over occupied states, or alternatively, it is a finite function since one integrates over the momentum up to the Fermi momentum.) This is also the case in the underlying theory for the $\Lambda(1520)$ [17,18], where the Bethe-Salpeter equation, which resumes only sequential one-loop diagrams, is used.

The use of dimensional regularization in the medium poses technical problems and the use of the cutoff is preferable. But to obtain the free amplitudes in the limit of $\rho \to 0$ we must use the same cutoff as in free space. To estimate uncertainties we shall take Λ in the range $\Lambda = 450-500$ MeV. Note that given this approximate equivalence of the cutoff method and dimensional regularization, basic symmetries such as chiral symmetry (up to the mass-breaking terms) are respected as

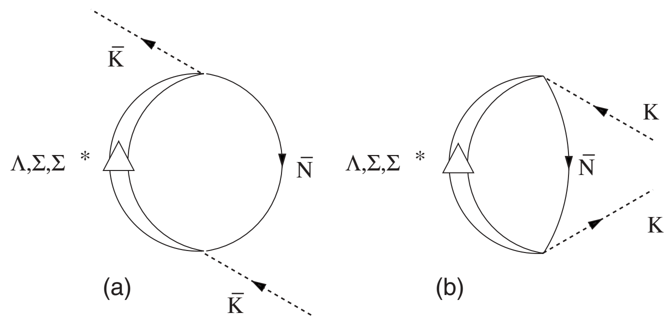

FIG. 2. Kaon P-wave self-energy diagrams: (a) $\bar{K}$ direct term; (b) K crossed term.

when using dimensional regularization [45]. However, this has only a symbolic value here from the moment that one diverts from a treatment using only chiral Lagrangians. Indeed, the important $\bar{K}N$ and $\pi\Sigma$ couplings to the $\pi\Sigma^*$ channel are introduced in our framework in a pure phenomenological way, simply demanding that the $\bar{K}N \to \bar{K}N$, $\bar{K}N \to \pi\Sigma$ D waves in $I = 0$ are reproduced.

Although the evaluation of Eq. (18) is straightforward, in the present work we neglect the contribution of the second term, which gives no contribution to the imaginary part and only a small one to the real part. So we consider the spectral function of antikaons only. The corresponding imaginary part of the $\Lambda^*(1520)$ self-energy reads

$$\begin{aligned}\mathrm{Im}[\Sigma(P,\rho)]^{\Lambda^*}_{\bar{K}N} &= -g^2_{\Lambda^*\bar{K}N}\frac{\pi}{3}\int\frac{d^3\boldsymbol{k}}{(2\pi)^3}\frac{\boldsymbol{k}^4 M_N}{E_N(\boldsymbol{P}-\boldsymbol{k})}\\ &\quad\times[1-n(|\boldsymbol{P}-\boldsymbol{k}|)]S_{\bar{K}}(\omega,\boldsymbol{k},\rho)\cdot\theta(\omega),\,|_{\omega}\\ &= {}_{P_0-E_N(\boldsymbol{P}-\boldsymbol{k})-V_N(\rho)}. \end{aligned}\tag{19}$$

In the following we shall briefly discuss how the antikaon spectral function is obtained. Here we follow closely Ref. [25].

The P-wave contribution to the $\bar{K}$ self-energy comes from the coupling of the $\bar{K}$ meson to hyperon particle-nucleon hole (YN^{-1}) excitations. The corresponding many-body mechanisms are shown in Fig. 2. Because of strangeness conservation, only direct terms [Fig. 2(a)] are permitted for the $\bar{K}$ excitations. Conversely, the K self-energy arises from the crossed terms [Fig. 2(b)]. The K^- meson can couple to $p\Lambda$, $p\Sigma^0$ or $n\Sigma^-$ and the $\bar{K}^0$ to $n\Lambda$, $n\Sigma^0$ or $p\Sigma^-$. The vertices $\bar{K}NY$ are derived from the D and F terms of the chiral Lagrangian given in the Appendix [see Eq. (A1)], expanding the unitary SU(3) matrix U up to one meson field. Using a nonrelativistic reduction of the $\gamma^\mu\gamma^5$ matrix, one finds

$$\begin{aligned}-it_{\bar{K}NY} &= C_{\bar{K}NY}(\boldsymbol{\sigma}\cdot\boldsymbol{k})\\ &= \left[\alpha_{\bar{K}NY}\frac{D+F}{2f}+\beta_{\bar{K}NY}\frac{D-F}{2f}\right](\boldsymbol{\sigma}\cdot\boldsymbol{k}),\end{aligned}\tag{20}$$

where $\boldsymbol{k}$ is the incoming $\bar{K}$ three-momentum, $f = 1.15 f_\pi$, and $\alpha_{\bar{K}NY}$, $\beta_{\bar{K}NY}$ are the SU(3) coefficients given in [25].

Following Ref. [46] we consider the $\bar{K}NY$ interaction in combination with the $\bar{K}N\Sigma^*(1385)$ transition, where the expression for the vertex function is given by

$$-it_{\bar{K}N\Sigma^*} = C_{\bar{K}N\Sigma^*}(\boldsymbol{S}^\dagger\cdot\boldsymbol{k}) = A_{\bar{K}N\Sigma^*}\frac{2\sqrt{6}}{5}\frac{D+F}{2f}(\boldsymbol{S}^\dagger\cdot\boldsymbol{k}).\tag{21}$$

The SU(3) coefficients $A_{\bar{K}N\Sigma^*}$ are given in Ref. [46]. These couplings were evaluated by first using the SU(6) quark model to relate the πNN coupling to the $\pi N\Delta$ one and then using SU(3) symmetry to relate the $\pi N\Delta$ coupling to the $\bar{K}N\Sigma^*$ one, since the $\Sigma^*(1385)$ belongs to the SU(3) decuplet of the Δ isobar. Alternatively, one can use the chiral Lagrangian that couples the octet of mesons and octet of baryons to the baryon decuplet given by [47]

$$\mathcal{L} = C\bar{T}_\mu A^\mu B + \mathrm{h.c.}\tag{22}$$

with the definition of the fields in [47]. In that reference it is also discussed that if SU(6) spin-flavor symmetry is further used to put together the octet and decuplet of baryons, the C coefficient and D, F of Eq. (A1) for the octet baryon-octet baryon-octet meson couplings can be related and one returns to the results of Ref. [46].

The P-wave $\bar{K}$ self-energy in symmetric nuclear matter can then be summarized as

$$\begin{aligned}\boldsymbol{k}^2\tilde{\Pi}^{(p)}_{\bar{K}}(\omega,\boldsymbol{k},\rho) &= \tfrac{1}{2}\left[C^2_{K^-p\Lambda}\mathcal{R}^2_\Lambda\right]\boldsymbol{k}^2\,\mathcal{U}_\Lambda(\omega,\boldsymbol{k},\rho)\\ &\quad+\tfrac{3}{2}\left[C^2_{K^-p\Sigma^0}\mathcal{R}^2_\Sigma\right]\boldsymbol{k}^2\,\mathcal{U}_\Sigma(\omega,\boldsymbol{k},\rho)\\ &\quad+\tfrac{1}{2}\left[C^2_{K^-p\Sigma^{*0}}\mathcal{R}^2_{\Sigma^*}\right]\boldsymbol{k}^2\mathcal{U}_{\Sigma^*}(\omega,\boldsymbol{k},\rho),\end{aligned}\tag{23}$$

where the Lindhard function $\mathcal{U}_Y(q)$ ($Y = \Lambda$, Σ or Σ^*) accounts for the direct term only [see Fig. 2(a)]. Its explicit expression can be found in Ref. [46]. In Eq. (23) $\mathcal{R}_{\Lambda(\Sigma)} = (1 - \omega/2M_{\Lambda(\Sigma)})$, $\mathcal{R}_{\Sigma^*} = (1 - \omega/M_{\Sigma^*})$ are the relativistic recoil vertex corrections [25]. In addition, we use the static form factors at the antikaon-baryon vertices of monopole type, $\Lambda^2/(\Lambda^2 + \boldsymbol{k}^2)$, with $\Lambda = 1$ GeV. We also take into account the short-range correlations in the hyperon-hole $YN^{-1} - YN^{-1}$ channels using a standard prescription

$$\Pi_{\bar{K}}(\omega,\boldsymbol{k},\rho) = \frac{\boldsymbol{k}^2\tilde{\Pi}^{(p)}_{\bar{K}}(\omega,\boldsymbol{k},\rho)}{1-g'\tilde{\Pi}^{(p)}_{\bar{K}}(\omega,\boldsymbol{k},\rho)},\tag{24}$$

where we assume the same value of the Landau-Migdal parameter $g' = 0.7$ as in Eq. (6).

III. RENORMALIZATION OF THE Σ*(1385)

As we have seen from Eq. (9), the problem of the in-medium renormalization of the $\Lambda(1520)$ in the $\pi\Sigma^*(1385)$ channel can be linked to the proper description of the properties of pions and $\Sigma^*(1385)$ in the nuclear medium. The renormalization scheme we employ for the $\Sigma^*(1385)$ is essentially the same as in the previous case except for the P-wave nature of the hadronic $\Sigma^*(1385)$ decay. This implies some peculiarities, for instance, the proper treatment of short-range correlations.

A. Σ*(1385) → πΛ + πΣ channel

We consider the following P-wave decay branches of the $\Sigma^*(1385)$ hyperon: $\Sigma^*(1385) \to \Lambda(1115) + \pi$ and $\Sigma^*(1385) \to \Sigma(1189) + \pi$. In what follows we shall further address the octet states $\Lambda(1115)$ and $\Sigma(1189)$ using the phenomenological optical potentials. The vertex functions describing the transitions are given by

$$-\,i t_{\pi Y\to\Sigma^*} = -C_{\pi Y\Sigma^*}(\boldsymbol{S}^\dagger\cdot\boldsymbol{k}), \tag{25}$$

where $Y = \Sigma(1189)$ or $\Lambda(1115)$. For coupling constants $C_{\pi Y\Sigma^*}$ we follow Ref. [46] and use the values

$$C_{\pi\Lambda\Sigma^*} = \frac{6}{5}\frac{D+F}{2f_\pi}, \quad C_{\pi\Sigma\Sigma^*} = -\frac{2\sqrt{3}}{5}\frac{D+F}{2f_\pi}. \tag{26}$$

From this, the explicit expressions for $\Sigma_{\pi Y}^{\Sigma^*}$ are given by

$$[\Sigma(P,\rho)]_{\pi\Lambda(1115)}^{\Sigma^*(1385)} = i\,(C_{\pi\Lambda\Sigma^*})^2\,\frac{1}{3}\int\frac{d^4k}{(2\pi)^4}\boldsymbol{k}^2\tilde{D}_\Lambda(P-k,\rho)\tilde{D}_\pi(k,\rho), \tag{27}$$

$$[\Sigma(P,\rho)]_{\pi\Sigma(1189)}^{\Sigma^*(1385)} = i\,(C_{\pi\Sigma\Sigma^*})^2\,\frac{2}{3}\int\frac{d^4k}{(2\pi)^4}\boldsymbol{k}^2\tilde{D}_\Sigma(P-k,\rho)\tilde{D}_\pi(k,\rho). \tag{28}$$

The additional factor of 2 in Eq. (28) comes from contribution of two possible charge states.

At this point we would like to mention that in a realistic calculation one would have to add strong repulsive forces at short distances and these would generate a short-range correlation. The correlations of this type of the interaction would effectively modulate the in-medium π exchange interaction [48,49], introducing the correlation parameter g'. The denominator in Eq. (6) takes into account this effect between P-wave bubbles in the diagrams of Fig. 3, but not between the external hyperon and the contiguous bubble. To account for this we make the separation between the longitudinal $\mathcal{V}_l$ and transverse $\mathcal{V}_t$ parts of the pion effective interaction [50,51] in the P-wave loop integrals

$$\left(\frac{D+F}{2f_\pi}\right)^2\frac{k_ik_j}{(k^0)^2-\boldsymbol{k}^2-m_\pi^2+i0^+} \longrightarrow \mathcal{V}_l(k)\hat{k}_i\hat{k}_j + \mathcal{V}_t(k)(\delta_{ij}-\hat{k}_i\hat{k}_j). \tag{29}$$

Here k_i is the Cartesian component of the unit vector $\hat{\boldsymbol{k}} = \boldsymbol{k}/|\boldsymbol{k}|$ and

$$\mathcal{V}_l(k) = \left(\frac{D+F}{2f_\pi}\right)^2\left[\frac{\boldsymbol{k}^2}{(k^0)^2-\boldsymbol{k}^2-m_\pi^2+i0^+}+g'\right]F(\boldsymbol{k})^2,$$
$$\mathcal{V}_t(q) = \left(\frac{D+F}{2f_\pi}\right)^2 g'F(\boldsymbol{k})^2. \tag{30}$$

This procedure can be represented by the last four diagrams in Fig. 3. In Eq. (30) $F(\boldsymbol{k})$ is a static form factor, $\Lambda^2/(\Lambda^2 + \boldsymbol{q}^2)$, with the cutoff scale $\Lambda = 1$ GeV. Hence, we make the following substitution in the self-energy of the $\Sigma^*(1385)$ [see Eqs. (27) and (28)]:

$$\left(\frac{D+F}{2f_\pi}\right)^2\boldsymbol{k}^2\tilde{D}_\pi(k,\rho) \to \mathcal{W}(k,\rho) = \frac{\mathcal{V}_l(k)}{1-\mathcal{U}(k,\rho)\mathcal{V}_l(k)} + \frac{2\mathcal{V}_t(k)}{1-\mathcal{U}(k,\rho)\mathcal{V}_t(k)}. \tag{31}$$

Using the Cutkosky rules [Eqs. (14)] supplemented by

$$\mathcal{W}(k,\rho) \to 2i\,\mathrm{Im}[\mathcal{W}(k,\rho)]\cdot\theta(k^0) \tag{32}$$

one may calculate the imaginary part of the loop integrals,

$$\mathrm{Im}[\Sigma(P,\rho)]_{\pi\Lambda}^{\Sigma^*} = \left(\tilde{C}_{\pi\Lambda\Sigma^*}\right)^2\frac{1}{3}\int\frac{d^3\boldsymbol{k}}{(2\pi)^3}\frac{M_\Lambda}{E_\Lambda(\boldsymbol{P}-\boldsymbol{k})} \times \mathrm{Im}[\mathcal{W}(k,\rho)]\cdot\theta(k^0)\big|_{k^0 = P^0 - E_\Lambda(\boldsymbol{P}-\boldsymbol{k})},$$

$$\mathrm{Im}[\Sigma(P,\rho)]_{\pi\Sigma}^{\Sigma^*} = \left(\tilde{C}_{\pi\Sigma\Sigma^*}\right)^2\frac{2}{3}\int\frac{d^3\boldsymbol{k}}{(2\pi)^3}\frac{M_\Sigma}{E_\Sigma(\boldsymbol{P}-\boldsymbol{k})} \times \mathrm{Im}[\mathcal{W}(k,\rho)]\cdot\theta(k^0)\big|_{k^0 = P^0 - E_\Sigma(\boldsymbol{P}-\boldsymbol{k})},$$

where $\tilde{C}_{\pi Y\Sigma^*}$ are the reduced coupling constants obtained from Eq. (26) by omitting the factor $(D+F)/2f_\pi$.

The vacuum-subtracted expression for $\mathcal{W}(k,\rho)$ reads

$$\delta\mathcal{W}(k,\rho) = \mathcal{W}(k,\rho) - \mathcal{W}(k,0) = \frac{\mathcal{U}(k,\rho)\mathcal{V}_l^2(k)}{1-\mathcal{U}(k,\rho)\mathcal{V}_l(k)} + \frac{2\mathcal{U}(k,\rho)\mathcal{V}_t^2(k)}{1-\mathcal{U}(k,\rho)\mathcal{V}_t(k)}. \tag{33}$$

From this the subtracted versions of the self-energy integrals [Eqs. (27) and (28)] take the form

$$[\delta\Sigma(P,\rho)]_{\pi\Lambda}^{\Sigma^*} = i\left(\tilde{C}_{\pi\Sigma\Sigma^*}\right)^2\frac{1}{3}\int\frac{d^4k}{(2\pi)^4}\tilde{D}_\Lambda(P-k,\rho)\delta\mathcal{W}(k,\rho), \tag{34}$$

$$[\delta\Sigma(P,\rho)]_{\pi\Sigma}^{\Sigma^*} = i\left(\tilde{C}_{\pi\Lambda\Sigma^*}\right)^2\frac{2}{3}\int\frac{d^4k}{(2\pi)^4}\tilde{D}_\Sigma(P-k,\rho)\delta\mathcal{W}(k,\rho). \tag{35}$$

Using Eqs. (34) and (35) we can evaluate the in-medium modification of both the real and imaginary parts. We also take into account the phenomenological binding corrections to Σ and Λ in nuclei to which the $\Sigma^*(1385)$ decays. These corrections are accounted for by [31]

$$V_\Lambda(\rho) = \mathrm{Re}\Sigma_\Lambda = V_\Sigma(\rho) = \mathrm{Re}\Sigma_\Sigma = -30\rho/\rho_0\,\mathrm{MeV}. \tag{36}$$

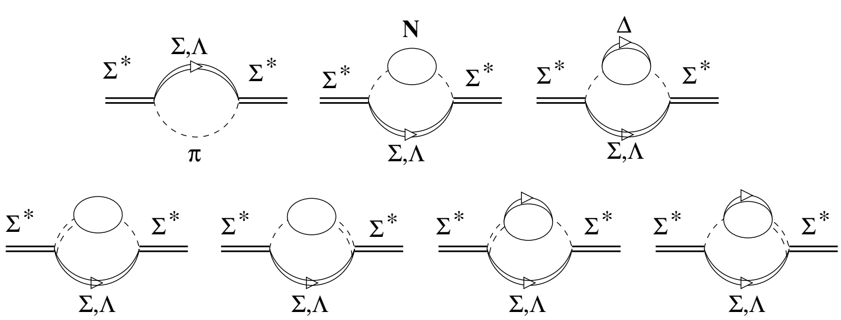

FIG. 3. In-medium renormalization of the $\Sigma^*(1385)$ in the $\pi\Sigma + \pi\Lambda$ channels. The last four diagrams account for the short-range correlations.

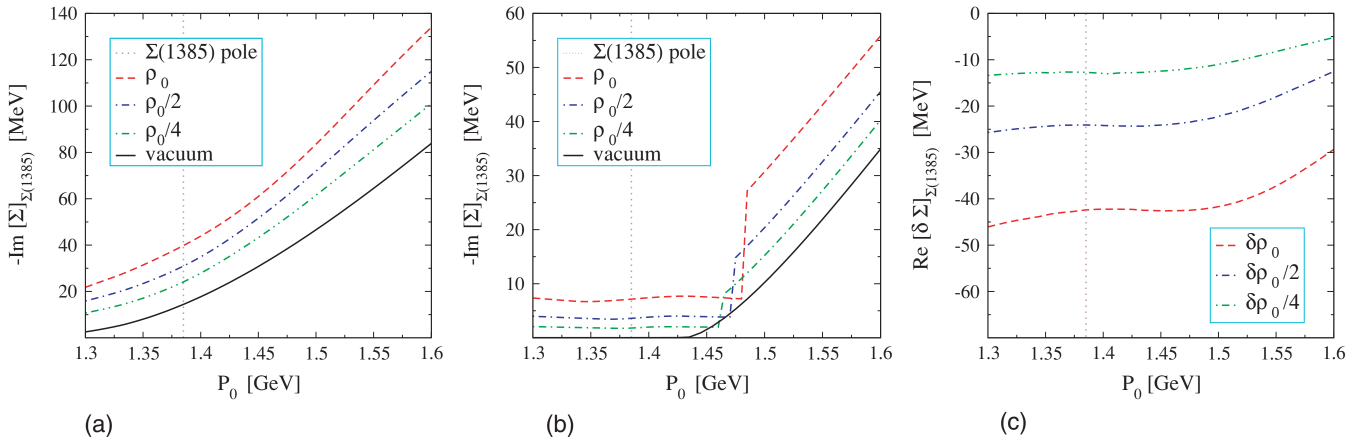

FIG. 4. (Color online) The imaginary parts (a) and (b) of the $\Sigma^*(1385)$ self-energy in the $(\pi\Sigma + \pi\Lambda)$ and $\bar{K}N$ channels, respectively, at several densities $\rho = \rho_0$ (dashed curve), $\rho_0/2$ (dot-dashed curve), and $\rho_0/4$ (dot-dot-dashed curve) as a function of $P^{\mu}_{\Sigma^*(1385)} = (P^0, 0)$. (c) The vacuum-subtracted real part of the $\Sigma^*(1385)$ self-energy. The dotted vertical line indicates the $\Sigma^*(1385)$ pole position.

B. $\Sigma^*(1385) \to \bar{K}N$ channel

The decay modes considered in the previous sections are allowed in free space. But there are channels such as $\Sigma^*(1385) \to \bar{K}N$ that may open up in the medium because of the additional phase space. By using the vertex function given by Eq. (21) the expression for the self-energy integral in this decay channel reads

$$[\Sigma(P,\rho)]^{\Sigma^*(1385)}_{\bar{K}N} = (C_{\bar{K}N\Sigma^*})^2 \frac{i}{3}\int_0^{\infty} d\omega \int \frac{d^4k}{(2\pi)^4}\boldsymbol{k}^2 \tilde{D}_N(P-k,\rho) \times \left[\frac{S_{\bar{K}}(\omega,\boldsymbol{k};\rho)}{k^0-\omega+i0^+} - \frac{S_K(\omega,\boldsymbol{k};\rho)}{k^0+\omega-i0^+}\right], \quad (37)$$

where $\tilde{D}_N$ is the in-medium nucleon propagator [Eq. (17)] and $\tilde{D}_{\bar{K}}$ is the $\bar{K}$-propagator, which is introduced according Eq. (16). The direct application of that result is not entirely correct because we deal with the P-wave kaons, which are also affected by the short-range correlations in this channel. For instance, in the expression for the imaginary part of the in-medium self-energy we obtain

$$\mathrm{Im}[\Sigma(P,\rho)]^{\Sigma^*}_{\bar{K}N} = (\tilde{C}_{\bar{K}N\Sigma^*})^2 \frac{1}{3}\int \frac{d^3\boldsymbol{k}}{(2\pi)^3}\frac{M_N}{E_N(\boldsymbol{P}-\boldsymbol{k})} \times [1-n(|\boldsymbol{P}-\boldsymbol{k}|)]\mathrm{Im}[\mathcal{W}_{\bar{K}}(\omega,\boldsymbol{k},\rho)]\cdot\theta(\omega)|_{\omega = P_0 - E_N(\boldsymbol{P}-\boldsymbol{k}) - V_N(\rho)}, \quad (38)$$

where $\tilde{C}_{\bar{K}N\Sigma^*} = C_{\bar{K}N\Sigma^*}(2f_\pi)/(D+F)$ and

$$\mathcal{W}_{\bar{K}} = \frac{\mathcal{V}_{l,\bar{K}}(k)}{1-\mathcal{U}_{\bar{K}}(k,\rho)\mathcal{V}_{l,\bar{K}}(k)} + \frac{2\mathcal{V}_{t,\bar{K}}(k)}{1-\mathcal{U}_{\bar{K}}(k,\rho)\mathcal{V}_{t,\bar{K}}(k)}, \quad (39)$$

with $\mathcal{V}_l$ and $\mathcal{V}_t$ given by Eq. (31) replacing m_π by $m_{\bar{K}}$ and using the same g'. In addition, in Eq. (39) $\mathcal{U}_{\bar{K}} = \tilde{\Pi}^{(p)}_{\bar{K}}/(\frac{D+F}{2f_\pi})^2$.

IV. RESULTS FOR THE $\Sigma^*(1385)$

Our results for the imaginary part of the $\Sigma^*(1385)$, or equivalently the width $\Gamma_{\Sigma^*} = -2\mathrm{Im}\Sigma_{\Sigma^*}$, in the $\pi\Sigma + \pi\Lambda$ decay channels and in the reference frame where $P = (P_0, 0)$ are shown in Fig. 4(a) for several densities. The vacuum value at the nominal pole position, $\Gamma_{\Sigma^*} \simeq 30$ MeV, is in good agreement with its empirical value $\Gamma_{\Sigma^*} = 35 \pm 4$ MeV [39]. Also the vacuum branching ratio $\Gamma_\Lambda/\Gamma_\Sigma \simeq 7.7$ compares well with the experimental value of 7.5 ± 0.5. By increasing the density we observe the broadening of the $\Sigma^*(1385)$ hyperon, and, as a result, at normal nuclear matter densities $\rho_0 = 0.17$ fm^{-3} the width becomes $\simeq$76 MeV. Note that the short-range correlations play an important role here. We find that, without short-range correlations, the effect of the medium on $\Sigma^*(1385)$ is unrealistically large and produces the increase of the width relative to the vacuum value by a factor of 5. This situation is similar to the behavior of the $\Delta(1232)$ isobar at finite density where the short-range correlations play an important role and strongly moderate the change of mass and width of the Δ isobar [51].

The results for another related channel, namely $\Sigma^*(1385) \to \bar{K}N$, is shown in Fig. 4(b). Because of Pauli blocking and the relatively weak interaction of $\bar{K}$ with the nuclear medium compared with pions the impact of this channel is small and adds an additional contribution $\simeq$7 MeV to the imaginary part. This is much smaller than that obtained in Ref. [52] where only the $\bar{K}N$ channel is considered and an on-shell approximation of the P-wave $\bar{K}N$ amplitude is made. This approximation is invalid for nuclei since it ties the momentum to the energy of the $\bar{K}$, which becomes imaginary below the $\bar{K}N$ threshold. However, the k^0 and $\boldsymbol{k}$ variables in matter are independent variables and $\boldsymbol{k}$ is always physical.

For the real part of the $\Sigma^*(1385)$ self-energy—the in-medium mass shift—we find an attractive potential at normal nuclear matter density with a strength of about $\simeq$−45 MeV [see Fig. 4(c)].

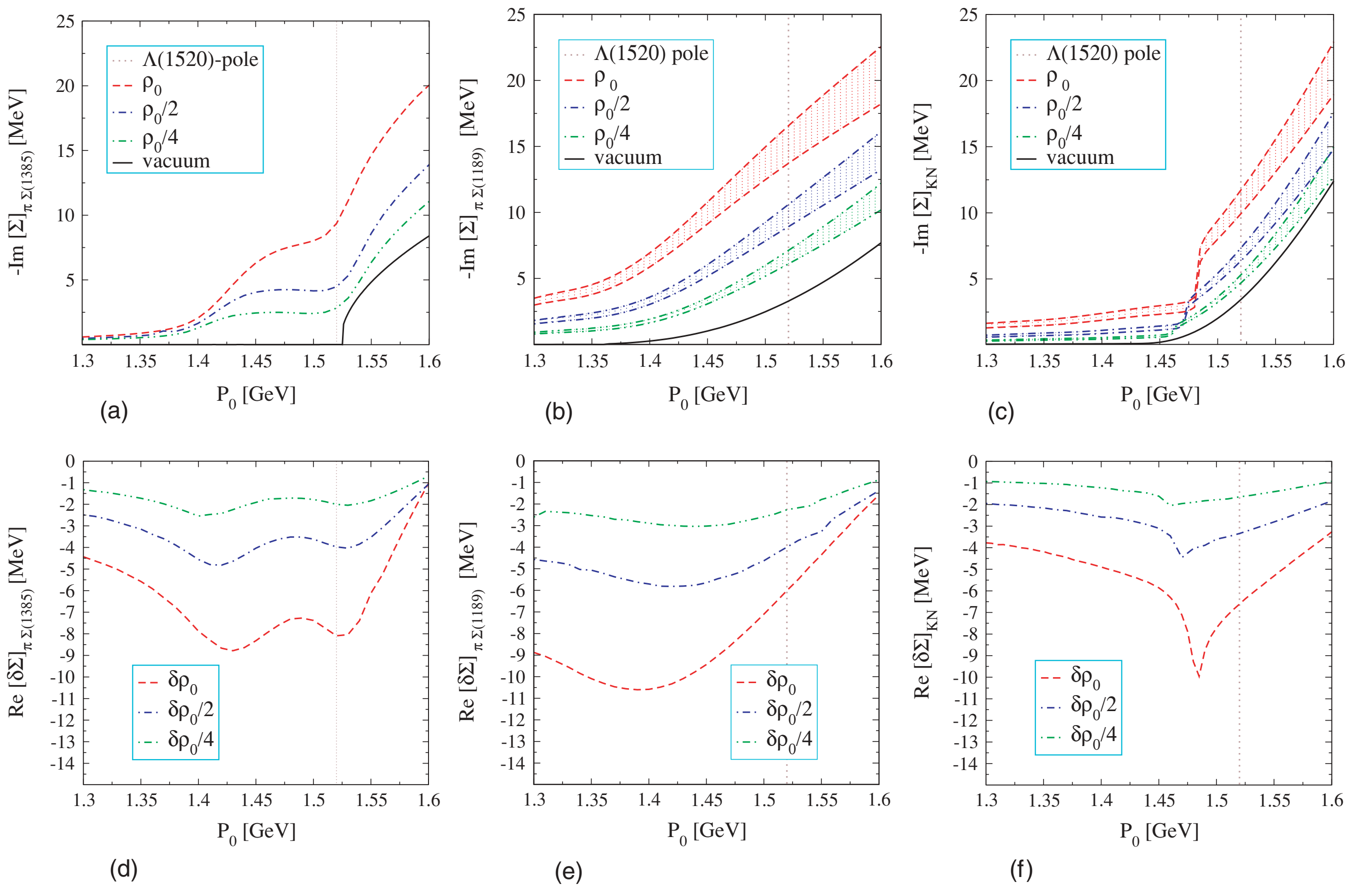

FIG. 5. (Color online) In-medium renormalization of the $\Lambda(1520)$ in the $\pi\Sigma^*(1385)$ (a,d), $\pi\Sigma$ (b,e), and $\bar{K}N$ (c,f) channels. The top and bottom panels are the imaginary and the vacuum-subtracted real parts of the self-energies, respectively. The vertical line indicates the $\Lambda(1520)$ pole position.

V. RESULTS FOR THE $\Lambda^*(1520)$

Using the $\Sigma^*(1385)$ discussed in the previous section we calculate the self-energy of the $\Lambda^*(1520)$ given by Eq. (9) in the $\pi\Sigma^*(1385)$ decay channel. Our results for the in-medium mass and width of the $\Lambda(1520)$ in this novel $\pi\Sigma^*$ channel are shown in Figs. 5 (a) and (d). As we have already noted for the nominal masses of hyperons involved there is an energy gap $\simeq$5 MeV, which makes the decay $\Lambda(1520) \to \pi\Sigma^*(1385)$ impossible. In the nuclear medium the pions decay to p-h and Δ-h and this effect opens up the $\Lambda(1520) \to (ph)\Sigma^*(1385)$ decay channel, which has a large available phase space. In Fig. 5 the dashed, dot-dashed, and dot-dot-dashed curves correspond to ρ_0, $\rho_0/2$, and $\rho_0/4$, respectively. In this channel we get a width for the $\Lambda(1520)$ of $\Gamma_{\Lambda(1520)} \simeq 18$ MeV at the nominal pole position and normal nuclear matter density. This value is even bigger than the free width of the $\Lambda(1520)$, which is $\Gamma_{\Lambda(1520)} \simeq 15$ MeV. The corresponding results for the vacuum-subtracted real part of the self-energy are shown in Fig. 5(d). Here we find a relatively weak attractive potential of about $\simeq -8$ MeV at the resonance pole position and normal nuclear matter density.

The renormalization of the $\Lambda^*(1520)$ in the $\pi\Sigma(1189)$ channel is shown in Figs. 5(b) and 5(e). These curves correspond to the regularization of the self-energy integral [Eq. (13)] when using the form factor $[F(\boldsymbol{k}^2)]^2 = [\Lambda^2/(\Lambda^2 + \boldsymbol{k}^2)]^2$ with the cutoff scale $\Lambda = 450-500$ MeV as we discussed previously. We use both limiting values and consider them as a sort of theoretical uncertainties. The results are presented as a band where the upper limit correspond to $\Lambda = 500$ MeV and the lower limit to $\Lambda = 450$ MeV. For instance, for $\Lambda = 500$ MeV we get a width of $\simeq$32 MeV for the $\Lambda^*(1520)$ at normal nuclear matter density ρ_0 and for $\Lambda = 450$ MeV we get $\simeq$26 MeV. One should compare this result with the free decay in this channel (solid curve) where the corresponding value is $\simeq$7 MeV only. The vacuum-subtracted real part of the self-energy is shown in Fig. 5(e) and was calculated with cutoff $\Lambda = 500$ MeV. The changes are moderate and for $\rho = \rho_0$ we get the attraction $\simeq -6$ MeV at energies near the pole position. In Figs. 5(c) and 5(f) we show our results for the renormalization of the $\Lambda(1520)$ in the $\bar{K}N$ channel. Here the results are quantitatively similar to those for the $\pi\Sigma$ channel. At normal density we get a width of $\simeq$20 MeV and additional attraction $\simeq -7$ MeV.

With present uncertainties we give a band of values for the width of the $\Lambda(1520)$ in the nuclear medium including now the free width and the in-medium renormalization from the $\pi\Sigma^*(1385)$, $\bar{K}N$, and $\pi\Sigma$ related channels. We show these

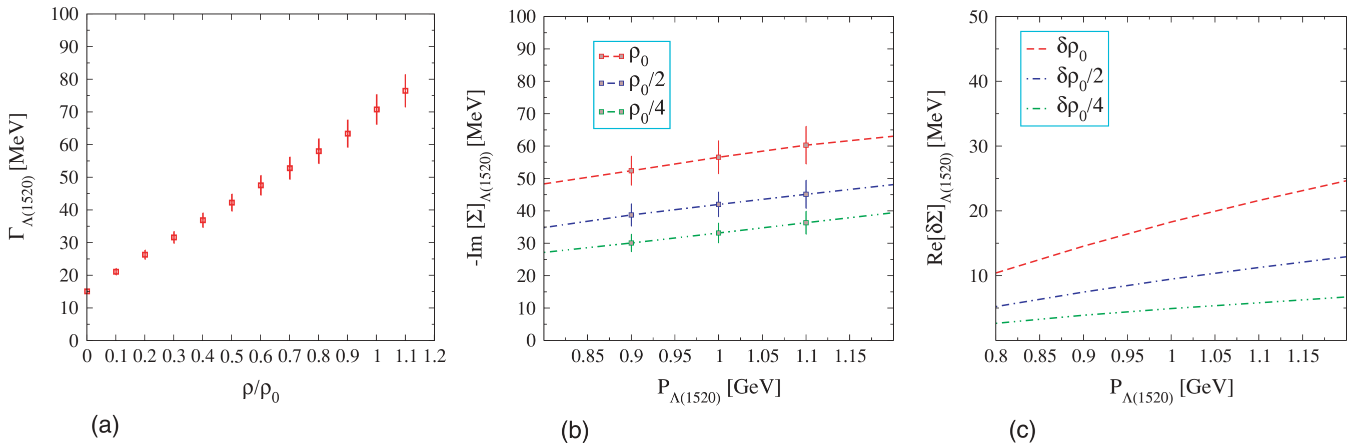

FIG. 6. (Color online) Values with theoretical uncertainties for the width of the Λ(1520) at rest in the medium, including the free width, as function of the nuclear matter density ρ/ρ_0 (a). The imaginary (b) and the vacuum-subtracted real (c) parts of the Λ(1520) self-energy as a function of a three-momentum $|\boldsymbol{P}_\Lambda|$.

results in Fig. 6(a) as a function of ρ/ρ_0. As one can see, at $\rho = \rho_0$ we get a Λ(1520) width of about ≃70−80 MeV, which is about five times the free width. These results are of the same order of magnitude as those obtained in Ref. [52] with the $\bar{K}N$ channel alone. The previous comments about the approximations done in [52] hold also in the present case. In addition, as mentioned in [47] the value of C in Eq. (22) has uncertainties of about 10%. If we allow C to change within this range and add the errors of the results in quadrature to those in Fig. 6(a) the results do not change appreciably.

Finally, we extend the discussion to some particular kinematics relevant for possible application of the presented formalism to reactions such as $\gamma p \to K^+\Lambda(1520)$ in nuclei where the Λ(1520) hyperon is produced with large momentum. In Fig. 6 the imaginary (b) and real (c) parts of the self-energy are shown for the Λ(1520) moving in nuclear matter as a function of a three-momentum $|\boldsymbol{P}_\Lambda|$. We can see that the imaginary part of the Λ(1520) self-energy is not changed much from its value at zero momentum. However, the real part changes sign from $|\boldsymbol{P}_\Lambda| = 0$ to $|\boldsymbol{P}_\Lambda| \simeq 1000$ MeV. But in both cases these changes are relatively small. Even if the experiment quoted here were more suited to determine the Λ(1520) width in the nucleus, there is already experimental information that allows us to get some hint on its size. The study performed in heavy-ion collisions [35,36] indicates that a better agreement of theory with experimental data is obtained by assuming that about half of the Λ(1520) produced are absorbed in the nucleus. Such a reduction can only be obtained with an in-medium width of tens of MeV as one can guess from comparison to studies done in ϕ production in the pA reaction [53].

When reaching this point we must face a problem of a different kind. Are the results that we obtain necessarily linked to the nature of the Λ(1520) as a coupled channel resonance as we maintain, or are they independent of its nature? The answer is not clear, but imagine how we would proceed with this problem if we did not know anything about the fact that the Λ(1520) can be dynamically generated through the interaction of the $\pi\Sigma^*$, $\bar{K}\Xi^*$, $\pi\Sigma$, and $\bar{K}N$ channels. From the partial decay channels of the Λ(1520) into $\pi\Sigma$ and $\bar{K}N$ we can obtain the coupling of the resonance to these channels. The coupling to $\pi\Sigma^*$ would be more problematic given the lack of phase space for this decay. The generation of the Λ(1520) within the coupled channels has as an output a rather precise value of the coupling of the Λ(1520) to $\pi\Sigma^*$ and this allows us to evaluate the $\pi\Sigma^*$-related decay channels in the medium. Hence, this contribution can be fairly accounted for within the framework of the Λ(1520) as dynamically generated from the interaction of the channels mentioned here. Without this knowledge we would have a large uncertainty from this source. It should however be mentioned that the in-medium width of the Λ(1520) coming from this source is only one-third of the total and with the other two-thirds coming from the $\bar{K}N$- and $\pi\Sigma$-related channels. For the latter ones we have the couplings to Λ(1520) independently of the nature of this resonance from the Λ(1520) partial decay widths into $\bar{K}N$ and $\pi\Sigma$ channels. With this knowledge we would proceed as we have done but would find that the results would depend sensibly on the cutoff that we use for the loops. The empirical hadron form factors from the *MBB* vertices help in making results converge, but the *D*-wave character of the interaction introduces an appreciable cutoff dependence for the extra cutoff in the loops. This means we could obtain the Λ(1520) self-energy in the medium but with a considerable uncertainty, not smaller than a factor of 2 with respect to the results that we obtain.

Can we conclude that a measurement of the Λ(1520) width in the medium that would agree with our predictions would confirm the nature of the Λ(1520) that we assume? We think the answer is no, because given the large uncertainties that a conventional many-body calculation not tied to the specific nature of the Λ(1520) would have, such a result would not be incompatible with other possible structures of the Λ(1520). However, we can think in a reverse way. What we have seen is that within our assumptions for the nature of the Λ(1520) resonance, we can provide the Λ(1520) width in the medium with relatively good precision. Hence, if an experiment found that this renormalization in the medium is much smaller or sensibly larger than we predict, this would pose serious problems to our model of the free $\Lambda(1520)$.

VI. CONCLUSIONS

We have addressed the problem of the self-energy of the $\Lambda(1520)$ and $\Sigma^*(1385)$ resonances in a nuclear medium and we have found relevant changes in the medium in both the real and imaginary parts, particularly in the latter. Considering the coupled channel character of the $\Lambda(1520)$ resonance, where the $\pi\Sigma^*$, $\bar{K}N$, and $\pi\Sigma$ channels play a very important role, particularly the $\pi\Sigma^*$, one finds pronounced changes in the width in the nuclear medium when the π component is allowed to become a p-h excitation and the $\bar{K}$ component a Λ-h, Σ-h or Σ^*-h excitation. When these p-h or Y-h excitations are allowed, all three channels increase the width by about 15–20 MeV, a larger amount than the free width; as a consequence one obtains at the end a $\Lambda(1520)$ width at $\rho = \rho_0$ of about four to five times the free width.

Such a spectacular change should be in principle easily observable experimentally. It will be interesting to look at suitable reactions to measure this. Recently, there have been interesting developments in this direction and experiments to measure changes in the ϕ width in the medium have been conducted [54] or are being proposed [55] by looking at the A dependence of the ϕ production in nuclei. Theoretical calculations of this A dependence show indeed that the method is suited for such investigations [53,56] and offers advantages over other methods. The study of these medium effects would be a novel and interesting enterprise to learn more about the $\Lambda(1520)$ along the lines discussed in the previous section. A calculation for the specific case of the $\Lambda(1520)$ production in the γ- and p-induced reactions from nuclei is already available, showing an A dependence of the nuclear cross sections that could be easily tested in existing experimental facilities [57].

ACKNOWLEDGMENTS

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APPENDIX: EFFECTIVE SU(3) LAGRANGIAN

The lowest order chiral Lagrangian, coupling the octet of pseudoscalar mesons to the octet of $1/2^+$ baryons, is

$$\begin{aligned} L_1^{(B)} &= \langle \bar{B} i\gamma^\mu \nabla_\mu B\rangle - M_B \langle \bar{B}B\rangle \\ &\quad + \tfrac{1}{2} D \langle \bar{B}\gamma^\mu \gamma_5 \{u_\mu, B\}\rangle + \tfrac{1}{2} F \langle \bar{B}\gamma^\mu\gamma_5 [u_\mu, B]\rangle, \end{aligned} \tag{A1}$$

where $\langle\rangle$ denotes the trace of SU(3) matrices and

$$\begin{aligned} \nabla_\mu B &= \partial_\mu B + [\Gamma_\mu, B], \\ \Gamma_\mu &= \tfrac{1}{2}(u^\dagger \partial_\mu u + u\partial_\mu u^\dagger), \\ U &= u^2 = \exp(i\sqrt{2}\Phi/f), \\ u_\mu &= i u^\dagger \partial_\mu U u^\dagger. \end{aligned} \tag{A2}$$

The SU(3) matrices for the mesons and the baryons are the following:

$$\Phi = \begin{pmatrix} \frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & \pi^+ & K^+ \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & K^0 \\ K^- & \bar{K}^0 & -\frac{2}{\sqrt{6}}\eta \end{pmatrix}, \tag{A3}$$

$$B = \begin{pmatrix} \frac{1}{\sqrt{2}}\Sigma^0 + \frac{1}{\sqrt{6}}\Lambda & \Sigma^+ & p \\ \Sigma^- & -\frac{1}{\sqrt{2}}\Sigma^0 + \frac{1}{\sqrt{6}}\Lambda & n \\ \Xi^- & \Xi^0 & -\frac{2}{\sqrt{6}}\Lambda \end{pmatrix}. \tag{A4}$$

In Eq. (A1) D and F are the axial vector coupling constants and f is the pseudoscalar meson decay constant.

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