Chiral unitary theory: Application to nuclear problems

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Abstract. In this talk we briefly describe some basic elements of chiral perturbation theory, χPT, and how the implementation of unitarity and other novel elements lead to a better expansion of the T-matrix for meson–meson and meson–baryon interactions. Applications are then done to the $\pi\pi$ interaction in nuclear matter in the scalar and vector channels, antikaons in nuclei and K^- atoms, and how the ϕ meson properties are changed in a nuclear medium.

1. Introduction

Nowadays it is commonly accepted that QCD is the theory of the strong interactions, with the quarks as building blocks for baryons and mesons, and the gluons as the mediators of the interaction. However, at low energies typical of the nuclear phenomena, perturbative calculations with the QCD Lagrangian are not possible and one has to resort to other techniques to use the information of the QCD Lagrangian. One of the most fruitful approaches has been the use of chiral perturbation theory, χPT. The theory introduces effective Lagrangians which involve only observable particles, mesons and baryons, respects the basic symmetries of the original QCD Lagrangian, particularly chiral symmetry, and organizes these effective Lagrangians according to the number of derivatives of the meson and baryon fields.

The lowest order chiral Lagrangian for the meson–meson interaction, invariant under Lorentz transformations, parity and charge conjugation with only two derivatives and linear in the quark masses is [1]

$$\mathcal{L}_2 = \frac{f^2}{4} \left\langle D_\mu U^\dagger D^\mu U + U^\dagger \mathcal{M} + \mathcal{M}^\dagger \mathcal{U} \right\rangle, \tag{1}$$

where $\langle \ \rangle$ means SU(3)-flavour trace, with $U(\Phi) = \exp(i\sqrt{2}\Phi/f)$, $D^\mu U$ the covariant derivative of U (normal derivative in the absence of external fields), and

$$\Phi = \frac{\vec{\lambda}}{\sqrt{2}} \vec{\phi} = \begin{pmatrix} \frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta_8 & \pi^+ & K^+ \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta_8 & K^0 \\ K^- & \bar{K}^0 & -\frac{2}{\sqrt{6}}\eta_8 \end{pmatrix}. \tag{2}$$

The mass matrix $\mathcal{M}$ is given in terms of the meson masses and in the limit of equal up and down quark masses reads

$$\begin{pmatrix} m_\pi^2 & 0 & 0 \\ 0 & m_\pi^2 & 0 \\ 0 & 0 & 2m_K^2 - m_\pi^2 \end{pmatrix}. \tag{3}$$

The meaning of the constant f can be appreciated when calculating from the lowest order Lagrangian the axial current. Then f becomes the pion decay constant in the chiral limit, which is about 93 MeV.

The next to leading order Lagrangian, $\mathcal{L}_4$, is constructed with the same building blocks as $\mathcal{L}_2$, preserving Lorentz invariance, parity and charge conjugation and explicit formulae can be seen in [1]. They contain four derivatives on the meson fields or meson masses to the fourth power and for the purpose of meson–meson interaction contain eight L_i free parameters which are adjusted to the data or alternatively derived in some models.

Chiral perturbation theory up to fourth order (in the number of derivatives) consists in the perturbative field theoretical calculation from the lowest order Lagrangian, which involves loop diagrams and divergences. These are cured by the introduction of the fourth order Lagrangian which also leaves some residual finite contribution to the amplitudes. Review papers on this issue can be seen in [2–4].

The inclusion of baryons in the chiral formalism is done in a similar way than the one used for the mesons. We consider here the octet of baryons

$$B = \begin{pmatrix} \frac{1}{\sqrt{2}}\Sigma^0 + \frac{1}{\sqrt{6}}\Lambda^0 & \Sigma^+ & p \\ \Sigma^- & -\frac{1}{\sqrt{2}}\Sigma^0 + \frac{1}{\sqrt{6}}\Lambda^0 & n \\ \Xi^- & \Xi^0 & -\frac{2}{\sqrt{6}}\Lambda^0 \end{pmatrix}. \tag{4}$$

The lowest order baryon–meson Lagrangian with at most two baryons can be written as:

$$\mathcal{L}_1 = \langle \bar{B} i\gamma^\mu \nabla_\mu B \rangle - M_B \langle \bar{B}B \rangle + \frac{1}{2}D \langle \bar{B}\gamma^\mu\gamma_5 \{u_\mu, B\} \rangle$$

$$+ \frac{1}{2}F \langle \bar{B}\gamma^\mu\gamma_5 [u_\mu, B] \rangle, \tag{5}$$

where

$$\nabla_\mu B = \partial_\mu B + [\Gamma_\mu, B], \tag{6}$$

with Γ_μ defined below and $D + F = g_A = 1.257$ and $D - F = 0.33$ [2]. We have

$$\Gamma_\mu = \frac{1}{2} \left\{ u^\dagger \partial_\mu u + u \partial_\mu u^\dagger \right\} \tag{7}$$

with u such that $u^2 = U$ and

$$u_\mu = i u^\dagger D_\mu U u^\dagger = u_\mu^\dagger. \tag{8}$$

2. Chiral unitary theory

The perturbative series in powers of p^2 of the meson momenta is slowly convergent and in the case of $\pi\pi$ interaction only works up to values of the centre of mass energy of the two meson system of the order of 400 MeV. This range is such that one misses the meson resonances like the $f_0(980)$, $a_0(980)$, $\sigma(500)$, $\rho(770)$, $K^*(900)$, etc. On the other hand, the intrinsic nature of the perturbative scheme makes it impossible to obtain the infinite value for the T matrix which characterizes the appearance of a resonance. Nonperturbative schemes are hence necessary to obtain these singularities. Several methods have been proposed recently in the context of the chiral Lagrangians [5–8]. Although they are technically different, they share many things in common and the essence of the method can be qualitatively understood by looking at how the inverse amplitude method in coupled channels works in [6]. Unitarity in coupled channels reads in our normalization:

$$\operatorname{Im} T_{\mathrm{if}} = -T_{\mathrm{in}}\, \rho_{nn}\, T_{nf}^*, \tag{9}$$

where ρ is a real diagonal matrix whose elements account for the phase space of the two meson intermediate states n which are physically accessible. With our normalization, ρ is given by

$$\rho_{nn}(s) = \frac{k_n}{8\pi\sqrt{s}} \theta(s - (m_{1n} + m_{2n})^2), \tag{10}$$

where k_n is the on shell center mass (CM) momentum of the meson in the intermediate state n and m_{1n}, m_{2n} are the masses of the two mesons in this state.

Isolating ρ from eq. (9) one has

$$\begin{aligned}
\rho &= -T^{-1} \cdot \operatorname{Im} T \cdot T^{*-1} \\
&= -\frac{1}{2i} T^{-1} \cdot (T - T^*) \cdot T^{*-1} \\
&= -\frac{1}{2i} (T^{-1*} - T^{-1}) = \operatorname{Im} T^{-1}.
\end{aligned} \tag{11}$$

From eq. (11) we can write

$$T^{-1} = \operatorname{Re} T^{-1} + i\rho. \tag{12}$$

Once we are at this point we realize that unitarity gives $\operatorname{Im} T^{-1}$ for free independently of the dynamics. Then it also looks most natural to expand T^{-1}, instead of T, in powers of p^2. Indeed, whenever T has a pole T^{-1} will have a zero, but an expansion in a power series around zero is not a problem while we cannot do it around a singularity. These are hence the basic elements introduced in these nonperturbative methods:

1. Unitarity is satisfied exactly.
2. Coupled channels are used.
3. The chiral expansion is made in the inverse of the T-matrix instead of T.

These ingredients are sufficient to allow one to reproduce all basic features of the meson–meson data up to 1.2 GeV, using the same Lagrangians, up to order four, used in chiral perturbation theory, as shown recently in [9]. The method used in [5] is identical, except that only one channel was used. In that method one can obtain the $\sigma(500)$ in the $L = I = 0$ sector but not the $f_0(980)$ and $a_0(980)$ which require the explicit inclusion of the $\bar{K}K$ channel. In [7] the Bethe Salpeter equation is used to unitarize in a single channel and provides a justification of the method followed in [10], where using the Bethe Salpeter equation and the on shell amplitudes at order $O(p^2)$ a good agreement with the data in the scalar sector could be obtained by means of only one regularizing parameter. In the work [8] a different idea is followed since the input is now based on the lowest order chiral Lagrangians plus the explicit fields (essentially vector mesons) which would correspond to preexisting states ($\bar{q}q$) of QCD. These are states which would remain in the absence of the meson–meson scattering generated by the residual interaction of the lowest order chiral Lagrangian. This follows the idea of [11] that the fourth order chiral Lagrangian is nothing but the reflection of the exchange between the mesons of these preexisting resonances.

In the $\bar{K}N$ system the use of the chiral Lagrangians in the meson–baryon sector together with the unitarization in coupled channels also proves very efficient and lead to excellent results for all the K^-p reactions at low energies using the Bethe Salpeter equation and the lowest order Lagrangian. Only a cut off introduced to regularize the loop functions is needed and one also obtains the $\Lambda(1405)$ resonance in $L = 0$ and $I = 0$ as a meson–baryon resonance, generated dynamically through the multiple scattering of the meson–baryon states induced by the Bethe Salpeter equation [12]. Similar findings were first obtained in [13] using however a more restricted space of coupled channels but introducing higher order Lagrangians to cope for those missing channels.

A recent review of these methods with many results can be seen in [14].

3. Application to nuclear problems

3.1 *The $\pi\pi$ interaction in the nuclear medium in the scalar sector*

The $\pi\pi$ interaction in a nuclear medium in the $L = I = 0$ channel (σ channel) has stimulated much theoretical work lately. It was realized that the attractive P-wave interaction of the pions with the nucleus led to a shift of strength of the $\pi\pi$ system to low energies and eventually produced a bound state of the two pions around $2m_\pi - 10$ MeV [15]. This state would behave like a $\pi\pi$ Cooper pair in the medium, with repercussions in several observable magnitudes in nuclear reactions [15]. The possibility that such effects could have already been observed in some unexpected enhancement in the $(\pi, 2\pi)$ reaction in nuclei [16] was also noticed there. More recent experiments where the enhancement is seen in the $\pi^+\pi^-$ channel but not in the $\pi^+\pi^+$ channel [17] have added more attraction to that conjecture.

The advent of the chiral unitary methods has added new interest in the subject and has allowed one to focus on the implications of the chiral constraints which had been known

to be relevant in this kind of studies [18]. In [19] the $\pi\pi$ interaction in a nuclear medium was studied following the lines of [10], renormalizing the pion propagators in the medium and introducing vertex corrections for consistency. The diagrams considered are depicted in figures 1, 2 and 3. The results for the imaginary part of the $\pi\pi$ amplitude in $L = I = 0$ are shown in figure 4. One can appreciate that there is an accumulation of strength in the region of small invariant masses of the two pion system, with qualitative results similar to those found in [15], but we do not get poles in that energy region. The accumulation of strength at these small invariant masses could raise hopes that the enhancement of strength at small invariant masses found in the $(\pi, 2\pi)$ reactions in nuclei in [17] could be explained. However, according to a recent study [20], the small nuclear densities involved in this reaction, which is rather peripheral, make the changes found in [19] insufficient to explain the experimental data.

The work of [19] includes only the $\pi\pi$ channel, since one is only concerned about the low energy region. This work has been generalized to coupled channels in [21] in order to make predictions for the modification of the $f_0(980)$ and $a_0(980)$ resonances in a nuclear medium. One finds there that both these resonances become wider in the medium as the nuclear density increases, with the $a_0(980)$ eventually melting into a background for densities close to normal nuclear matter density. The $f_0(980)$ resonance, which in the free space is narrower than the $a_0(980)$, still would keep its identity at these high densities but with a width as large as 100 MeV or more. How to produce these resonances in a nucleus in order to check the predictions of these studies is a present experimental challenge.

3.2 *Isovector $\pi\pi$ scattering and the ρ meson in the nuclear medium*

The modification of the $\pi\pi$ amplitude in the $L = I = 1$ sector in the nuclear medium is equivalent to addressing the modification of the ρ properties in the medium. Once again this topic has received much attention. A good review of the current situation can be found in [22]. Once again we have looked at the problem from the chiral unitary point of view [23]. In a first step a combined study of the pion electromagnetic form factor and the
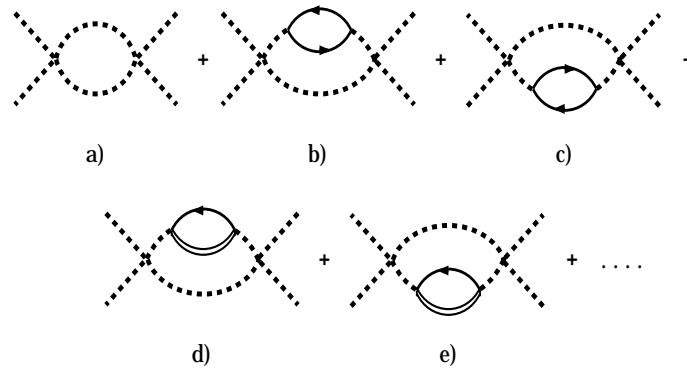

Figure 1. Terms appearing in the scattering matrix allowing the pions to excite ph and Δh components.
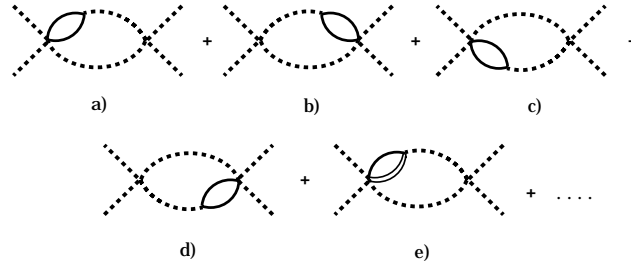
Figure 2. Terms of the $\pi\pi$ scattering series in the nuclear medium related to three meson baryon contact terms from the Lagrangian of eq. (5).
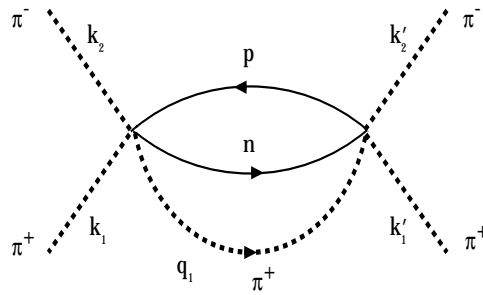

Figure 3. Diagram involving the three meson baryon contact terms of figure 2 in each of the vertices.

$\pi\pi$ scattering in the vector sector has been accomplished in [24] using the chiral unitary method with explicit resonances of [8]. The nuclear corrections are generated similarly to those discussed above in the scalar sector, introducing both selfenergy correction in the pion propagators as well as vertex corrections which are generated by the chiral Lagrangians and are requested by the gauge invariance of the vector mesons. The results obtained can be summarized in figure 5, where the real and imaginary parts of the $\pi\pi$ amplitude are plotted for normal nuclear matter density. The different curves correspond to different choices of a regularization parameter, or cut off, which lead to basically the same results in free space once the couplings and bare mass of the ρ are changed simultaneously. One can see that in the medium the results are also rather stable.

We can observe that the ρ becomes much broader in the nuclear medium but there is also a change in the position of the peak, implying a shift of the ρ mass to lower energies by about 50 MeV at $\rho = \rho_0$. The renormalization of the ρ still gets contributions from other sources, particularly from the excitation of the $N^*(1520)h$ by the ρ [25], since the decay of the $N^*(1520)$ resonance into $N\rho$ is one of the important channels, only reduced in practice because of the limited phase space for the decay. The inclusion of this channel produces some extra strength at lower energies, around 250 MeV below the ρ peak.
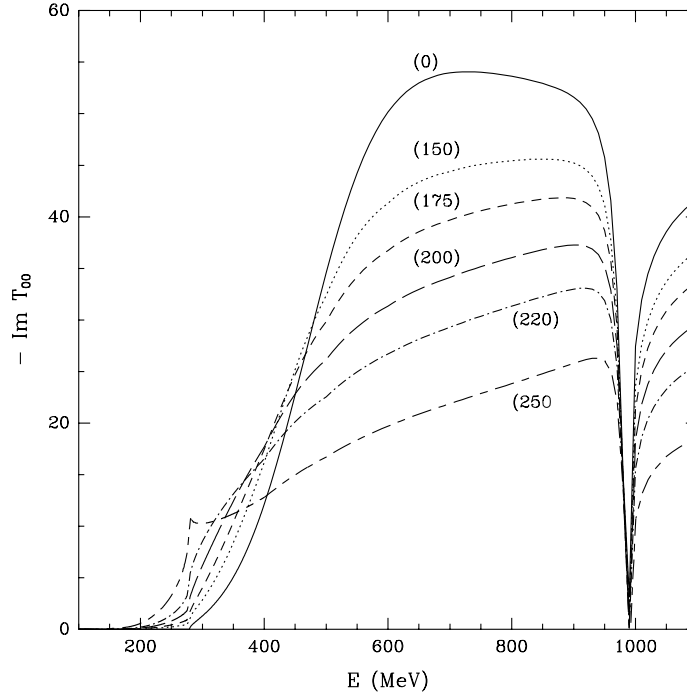
Figure 4. Im T_{22} for $\pi\pi \rightarrow \pi\pi$ scattering in $J = I = 0$ (T_{00} in the figure) in the nuclear medium for different values of k_F versus the CM energy of the pion pair. The labels correspond to the values of k_F in MeV.

In [25] the selfconsistent consideration of the $N^*(1520)$ resonance in the medium, together with the widening of the ρ width, leads to a broad peak for the spectral function below the ρ meson peak resulting in a considerable widening of the ρ strength which would make a shift of 50 MeV irrelevant when compared to an in medium width of more than 300 MeV. Present studies at CERN [26,27] and at lower energies at Bevalac [28] seem to be consistent with a large broadening of the ρ and more experimental studies are under way at GSI (HADES collaboration) [29,30].

4. Kaons in a nuclear medium

4.1 K^- *deuteron scattering length*

The $\bar{K}N$ interaction is quite strong and when one studies the interaction of $\bar{K}$ in nuclei important renormalization effects take place. This is already visible in the interaction of K^- with the deuteron which has been the subject of much study in the past [31–33]. It was already known that the evaluation of the K^- deuteron scattering length required the consideration of multiple scattering of the K^- which was done using Faddeev equations. We have also made some contribution to the field [34], by using again Faddeev equations in

the fixed scatterer approximation which is known to be rather reliable [35], but using input from the chiral unitary approach of [12] for the elementary amplitudes. The results are summarized in the table 1, where we show the results with the impulse approximation (IA), the IA plus double rescattering, idem plus triple rescattering, the effect of including the charge exchange ($\bar{K}^0$ exchanged between the nucleons), etc. We can see in the table that the full result of the Faddeev calculation is quite different from any of the approximations, and one can also see that the multiple scattering series does not converge, which forces the solution by means of the coupled Faddeev equations. We can also see there that the use of isospin symmetry leads to somewhat inaccurate results and finally, we also show for comparison the results of [35] which are also quite different than those obtained here, as is also the case for the results obtained in [36]. The main reason lies in the different elementary amplitudes $\bar{K}N$ provided by the chiral approach with respect to those used as input in the previous approaches. Experiments to measure this scattering length are under way in Frascati [37] and we hope they can be accomplished in the near future, hence providing further constraints to test the predictions of these chiral models.
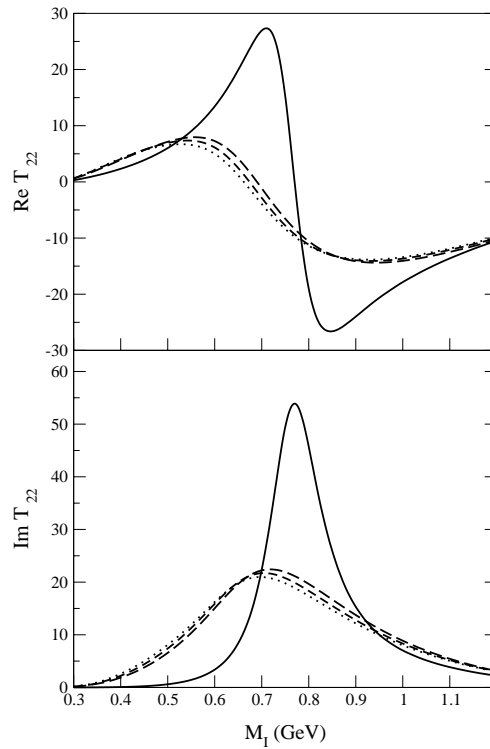

Figure 5. Real and imaginary part of the $\pi\pi$ amplitude in $L = I = 1$ for $\rho = \rho_0$ and several values of q_{max}. Short dashed lines stand for $q_{max} = 1$ GeV, long dashed and dotted lines being for $q_{max} = 0.9, 1.1$ GeV respectively.

Table 1. K^--deuteron scattering length (in fm) calculated using different approximations.

Approximations	Physical basis	Isospin basis	Isospin basis, ref. [19]
IA	$-0.260 + i\,1.872$	$-0.318 + i\,2.013$	$-0.364 + i\,1.826$
IA + double resc.	$-2.735 + i\,2.895$	$-3.168 + i\,3.717$	$-2.380 + i\,1.485$
IA + double + triple resc.	$-3.849 + i\,2.963$	$-5.195 + i\,4.935$	$-2.858 + i\,0.089$
A_{Kd} (only el. resc.)	$-1.161 + i\,1.336$	$-1.255 + i\,1.518$	$-0.997 + i\,1.212$
A_{Kd} (charge exch.)	$-0.454 + i\,0.573$	$-0.654 + i\,0.937$	$-0.539 + i\,0.079$
A_{Kd} (total)	$-1.615 + i\,1.909$	$-1.909 + i\,2.455$	$-1.536 + i\,1.291$

4.2 $\bar{K}$ in nuclei

Next we address the properties of the $\bar{K}$ in the nuclear medium which have been studied in [38]. The work is based on the elementary $\bar{K}N$ interaction which was obtained in [12] using a coupled channel unitary approach with chiral Lagrangians.

The coupled channel formalism requires to evaluate the transition amplitudes between the different channels that can be built from the meson and baryon octets. For K^-p scattering there are 10 such channels, namely K^-p, $\bar{K}^0n$, $\pi^0\Lambda$, $\pi^0\Sigma^0$, $\pi^+\Sigma^-$, $\pi^-\Sigma^+$, $\eta\Lambda$, $\eta\Sigma^0$, $K^+\Xi^-$ and $K^0\Xi^0$. In the case of K^-n scattering the coupled channels are: K^-n, $\pi^0\Sigma^-$, $\pi^-\Sigma^0$, $\pi^-\Lambda$, $\eta\Sigma^-$ and $K^0\Xi^-$.

At low energies the transition amplitudes can be written as

$$V_{ij} = -C_{ij} \frac{1}{4f^2} (k_j^0 + k_i^0) \,, \tag{13}$$

where $k_{i,j}^0$ are the energies of the mesons and the explicit values of the coefficients C_{ij} can be found in [12]. The coupled-channel Bethe–Salpeter equation with the kernel (potential) V_{ij} was used in [12] in order to obtain the elastic and transition matrix elements in the K^-N reactions. The diagrammatic expression of this series can be seen in figure 6. The Bethe Salpeter equations in the center of mass frame read

$$T_{ij} = V_{ij} + V_{il}\, G_l\, T_{lj}, \tag{14}$$

where the indices i, l, j run over all possible channels and G_l stands for the loop function of a meson and a baryon propagators.

In order to evaluate the $\bar{K}$ self-energy, one needs first to include the medium modifications in the $\bar{K}N$ amplitude, T_{eff}^α ($\alpha = \bar{K}p, \bar{K}n$), and then perform the integral over the nucleons in the Fermi sea:

$$\Pi_{\bar{K}}^s(q^0, \vec{q}, \rho) = 2 \int \frac{\mathrm{d}^3p}{(2\pi)^3} n(\vec{p}) \left[T_{\text{eff}}^{\bar{K}p}(P^0, \vec{P}, \rho) + T_{\text{eff}}^{\bar{K}n}(P^0, \vec{P}, \rho) \right]. \tag{15}$$

The values $(q^0, \vec{q})$ stand now for the energy and momentum of the $\bar{K}$ in the lab frame, $P^0 = q^0 + \varepsilon_N(\vec{p})$, $\vec{P} = \vec{q} + \vec{p}$ and ρ is the nuclear matter density.
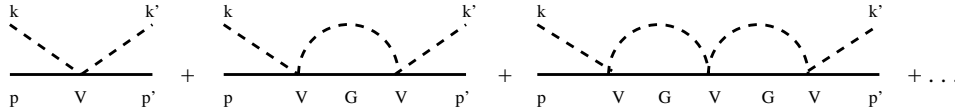
Figure 6. Diagrammatic representation of the Bethe–Salpeter equation.

We also include a p-wave contribution to the $\bar{K}$ self-energy coming from the coupling of the $\bar{K}$ meson to hyperon-nucleon hole (YN^{-1}) excitations, with $Y = \Lambda, \Sigma, \Sigma^{*}(1385)$. The vertices MBB' are easily derived from the D and F terms of eq. (5). The explicit expressions can be seen in [38]. At this point it is interesting to recall three different approaches to the question of the $\bar{K}$ self-energy in the nuclear medium. The first interesting realization was the one in [40–42], where the Pauli blocking in the intermediate nucleon states in figure 6 induced a shift of the $\Lambda(1405)$ resonance to higher energies and a subsequent attractive $\bar{K}$ self-energy. The work of [43] introduced a novel and interesting aspect, the self-consistency. Pauli blocking required a higher energy to produce the resonance, but having a smaller kaon mass led to an opposite effect, and as a consequence the position of the resonance was brought back to the free position. Yet, a moderate attraction on the kaons still resulted, but weaker than anticipated from the former work. The work of [38] introduces some novelties. It incorporates the self-consistent treatment of the kaons done in [43] and in addition it also includes the self-energy of the pions, which are let to excite ph and Δh components. It also includes the mean field potentials of the baryons. The obvious consequence of the work of [38] is that the spectral function of the kaons gets much wider than in the two former approaches because one is including new decay channels for the $\bar{K}$ in nuclei. This can be seen in figure 7. The work of [38] leads to an attractive potential around nuclear matter density and for kaons close to threshold of about 40 MeV and a width of about 100 MeV.

5. Kaonic atoms

In the work of [44] the kaon self-energy discussed above has been used for the case of kaonic atoms, where there are abundant data to test the theoretical predictions. One uses the Klein Gordon equation and obtains two families of states. One family corresponds to the atomic states, some of which are those already measured, and which have energies around or below 1 MeV and widths of about a few hundred KeV or smaller. The other family corresponds to states which are nuclear deeply bound states, with energies of 10 or more MeV and widths around 100 MeV. In figure 8 we can see the results obtained for shifts and widths for a large set of nuclei around the periodic table. The agreement with data is sufficiently good to endorse the fairness of the theoretical potential. A best fit with a strength of the potential slightly modified around the theoretical values can lead to even better agreement as shown in [45] and serves to quantify the level of accuracy of the theoretical potential, which is set there at the level of 20–30 per cent as an average. The curious thing is that there are good fits to the data using potentials with a strength at $\rho = \rho_0$ of the order of 200 MeV [46]. As shown in [44], the results obtained there and
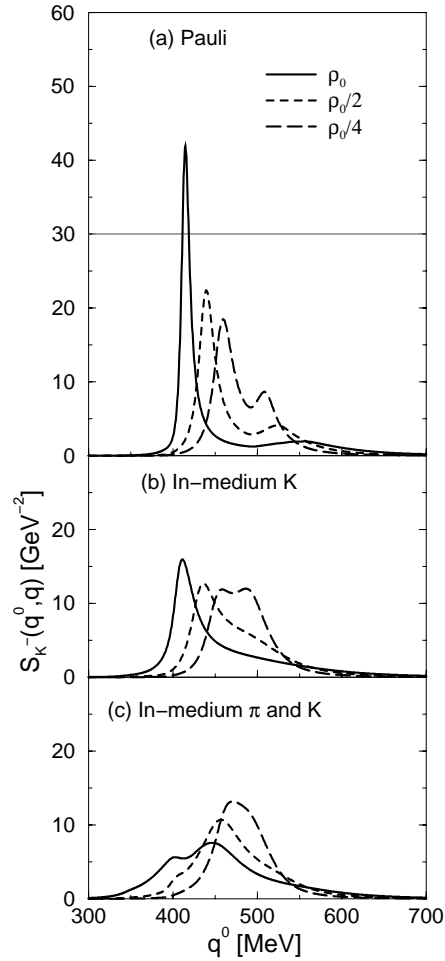
Figure 7. Kaon spectral function at several densities.

those obtained using the potential of [46] are in excellent agreement for the atomic states. The differences in the two potentials appear in the deeply bound nuclear states. The deep potential provides extra states bound by about 200 MeV, while the potential of [38] binds states at most by 40 MeV. This remarkable finding can be interpreted as saying that the extra bound states, forcing the atomic states to be orthogonal to them, introduces extra nodes in the wave function and pushes the atomic states more to the surface of the nucleus, acting effectively as a repulsion which counterbalances the extra attraction of the potential. This observation also tells that pure fits to the K^- atoms are not sufficient to determine the strength of the K^- nucleus potential. Other solutions with even more attraction at $\rho = \rho_0$ are in principle possible, provided they introduce new states of the deeply bound nuclear family. On the other hand, the work of [45] also tells us that at least an attraction as the one provided by the theoretical potential is needed.
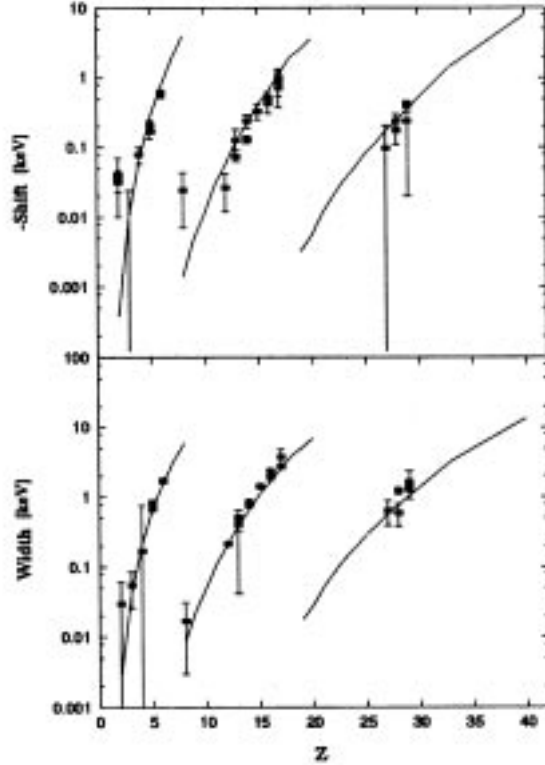
Figure 8. Shifts and widths of kaonic atoms.

The potential used for the K^- atoms has used only the s-wave potential, ignoring a possible contribution of the p-wave potential and some nonlocalities associated to the energy and momentum dependence of the s-wave part of the optical potential. In a recent paper [47] these corrections have been evaluated and they have been reported as providing a contribution to the potential as large or even larger than the local s-wave part of the potential used in [44]. In view of the conflict of such a finding with all previous works on the topic which have systematically relied on the local s-wave potential, we have reconducted a thorough study of these nonlocalities together with the contribution of the genuine p-wave part of the potential stemming from the elementary p-wave amplitude. The work is done in [48] and the results are quite different of those obtained in [47], the main reason for the differences being the inconsistent treatment of the low density limit in [47]. The findings of [48] lead to corrections from all sources of nonlocalities which are rather small, smaller than the experimental errors, and hence this justifies the neglect of such terms in [44] and in all previous studies of kaonic atoms.

6. ϕ decay in nuclei

Finally let us say a few words about the ϕ decay in nuclei. The work reported here [50] follows closely the lines of [51,52], however, it uses the updated $\bar{K}$ selfenergies of [38]. In

the present case the ϕ decays primarily in $K\bar{K}$, but these kaons can now interact with the medium as discussed previously. For the self-energy of the K, since the KN interaction is not too strong and there are no resonances, the $t\rho$ approximation is sufficient. In figure 9 we show the results for the ϕ width at $\rho = \rho_0$ as a function of the mass of the ϕ, separating the contribution from the different channels. What we observe is that the consideration of the s-wave $\bar{K}$-self-energy is responsible for a sizeable increase of the width in the medium, but the p-wave is also relevant, particularly the Λh excitation and the $\Sigma^* h$ excitation. It is also interesting to note that the vertex corrections [53] (Yh loops attached to the ϕ decay vertex) are now present and do not cancel off shell contributions like in the case of the scalar mesons. Their contribution is also shown in the figure and has about the same strength as the other p-wave contributions. The total width of the ϕ that we obtain is about 22 MeV at $\rho = \rho_0$, about a factor two smaller than the one obtained in [51,52], yet, the important message is the nearly one order of magnitude increase of the width with respect to the free one. We are hopeful that in the near future one can measure the width of the ϕ in the medium, from heavy ion reactions or particle nucleus interactions, although it will require careful analyses as shown in [54] for the case of $K\bar{K}$ production in heavy ion collisions, where consideration of the possibility that the observed kaons come from ϕ decay outside the nucleus leads to nuclear ϕ widths considerably larger than the directly observed ones.

In order to facilitate the experimental search of these medium modifications we have recently proposed a method based on the ϕ photoproduction in nuclei [55] producing ϕ with small momenta (around 150 MeV/c) which would be forced to decay inside the nucleus. These slow ϕ's would come from the elementary ϕ photoproduction on a nucleon in the backward direction in the CM frame. A recent measurement of this quantity at Jefferson Lab [56] gives great hopes that the photoproduction in nuclei with these small ϕ momenta could be feasible, since the cross sections obtained, [55], are of the same order of magnitude as those measured in [56]. The results of [55] indicate that, once the corrections for the final state interaction of the kaons from the ϕ decay are done, one can reconstruct an invariant mass of the ϕ in the medium at not too high densities but big enough to produce a ϕ width about twice as big as the free one.

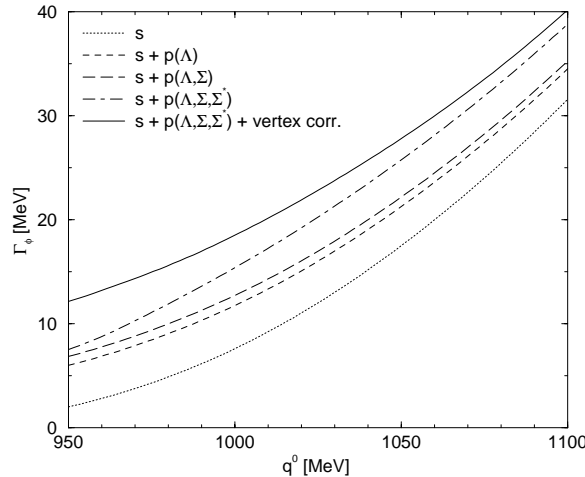**Figure 9.** ϕ width at $\rho = \rho_0$.

7. Summary

In summary, we have reported here on recent work which involves the propagation of kaons in the nuclear medium. All of them together provide a test of consistency of the theoretical ideas and results previously developed and reported here. If we gain confidence in those theoretical methods one can proceed to higher densities and investigate the possibility of kaon condensates in neutron stars [57]. The weak strength of our $\bar{K}$ potential would make however the phenomenon highly unlikely.

On the other hand, we can also extract some conclusions concerning the general chiral framework: (1) The chiral Lagrangians have much information in store. (2) Chiral perturbation theory allows one to extract some of this information. (3) The chiral unitary approach allows one to extract much more information. (4) These unitary methods combined with the use of standard many body techniques are opening the door to the investigation of new nuclear problems in a more accurate and systematic way, giving rise to a new field which could be rightly called 'chiral nuclear physics'. As chiral theory becomes gradually a more accepted tool to deal with strong interactions at intermediate energies, chiral nuclear physics is bound to follow analogously in the interpretation of old and new phenomena in nuclei.

Acknowledgments

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