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Abstract

The present status of the in-medium invariant $\pi - \pi$ mass distribution in the sigma-meson channel around the threshold is reviewed from the theory side and contrasted with recent experiments on $(\pi, 2\pi)$ knock-out reactions. A preliminary investigation indicates that the strongly target-mass dependent invariant mass enhancement of the two pions can be explained theoretically. A more refined reaction theory is needed to confirm this result. In the theoretical description, based on the linear sigma model, emphasis is put on constraints from chiral symmetry.

In-medium pion-pion correlations have recently attracted much attention both on the theoretical and experimental sides. For instance, in the $J = I = 0$ (sigma-meson)-channel the interest is quite obvious, since it is common belief that some kind of effective sigma meson is responsible for the midrange attraction of the nucleon-nucleon potential [1]. It is therefore important to know how such a meson is modified in a nuclear medium, a question which is also intensively studied for other mesons. In the past we have developed phenomenological models in which a 'bare' sigma meson is coupled to the decay channel into two pions. The coupling constants were adjusted such that the experimental $\pi\pi$ phase shifts in the $J = I = 0$ are reproduced. The same procedure was adopted simultaneously for the rho meson in the $I = J = 1$ channel of the two pions with a perfect reproduction of the position and the width of the rho meson [2]. Certainly, because of the s-wave nature of the interaction, the coupling of the sigma meson to the pions is much stronger than in the rho-meson channel such that the sigma meson becomes completely hybridised with two pions and the mass distribution only shows a very broad peak, roughly $500 MeV$ wide. We then accounted for medium effects by coupling the pions to Δ-h and p-h excitations in the usual way and found as a function of density attractive downward shifts of portions of the sigma-meson mass distribution, even far below the $2m_\pi$ threshold [3]. This is driving the system into an instability at densities close to saturation. In-medium vertex corrections which could give some repulsion in the correlated pion-pair system were investigated in ref. [4] and found not to be sufficient to hinder this phenomenon. In fact, it was recognised later that it is absolutely necessary to fulfil constraints from chiral symmetry in order to generate the needed repulsion below threshold which can prevent the invasion of the sigma-meson mass distribution into this region. Chiral symmetry was implemented phenomenologically in [5] and it was indeed found that the build up of strength below threshold was strongly reduced. However, in ref. [5] we did not project on the sigma-meson channel but rather showed the imaginary part of the $\pi - \pi$ T-matrix Because of form factors in the T-matrix, the strength distribution below threshold is very much suppressed. In the following, we will show that the sigma-meson mass distribution proper is still substantially shifted downwards in the nuclear medium. However, this will only serve for demonstration purposes. For real processes such as $eg.$ the nucleon-nucleon force, mediated by a sigma meson-exchange, one has to be careful to also include the exchange of a correlated $\pi - \pi$ pair such that chiral symmetry is conserved [6].

As a model for the vacuum $\pi\pi$ s-wave correlations we consider first the leading-order contributions of the $1/N$-expansion using the two flavour linear σ-model [7]. This, if treated correctly, leads to a symmetry-conserving approach and hence fulfils Ward identities and all chiral symmetry constraints [7]. It also preserves the unitarity of the S-matrix, since it is based on an RPA-type equation. Therefore, the $\pi\pi$ T-matrix is non-perturbative which is absolutely needed if one wants to describe such features as accumulation of strength and eventually resonances at low energies.

Explicitly we have for the T-matrix to leading order

$$T(E, \vec{p}) = \frac{V_{\pi\pi}(E, \vec{p})}{1 - \frac{1}{2}V_{\pi\pi}(E, \vec{p})\Sigma_{\pi\pi}(E, \vec{p})} \ , \tag{1}$$

where $\Sigma_{\pi\pi}(p)$ is the usual $\pi\pi$ bubble

$$\Sigma_{\pi\pi}(p^2) = -i \int \frac{d^4q}{(2\pi)^4} D_\pi(q) D_\pi(p-q) , \qquad (2)$$

and $V_{\pi\pi}$ the tree-level $\pi\pi$-scattering amplitude given by

$$V_{\pi\pi}(E, \vec{p}) = N \frac{\mathcal{E}_\sigma^2(\vec{0}) - \mathcal{E}_\pi^2(\vec{0})}{f_\pi^2} \frac{E^2 - \mathcal{E}_\pi^2(\vec{p})}{E^2 - \mathcal{E}_\sigma^2(\vec{p})} , \qquad (3)$$

where $\mathcal{E}_\pi^2(\vec{q}) = \sqrt{m_\pi^2 + \vec{q}^2}$ and $\mathcal{E}_\sigma^2(\vec{q}) = \sqrt{\mathcal{E}_\sigma^2 + \vec{q}^2}$. Here m_π is the pion mass and $\mathcal{E}_\sigma$ is the quasi-sigma mass. To this order, the pion decay constant f_π is related to the sigma-condensate $\langle\sigma\rangle$ via: $f_\pi^2 = N\langle\sigma\rangle^2$, where N stands for the number of pion charges.

To achieve an acceptable description of the s-wave $\pi\pi$ phase shifts we supplement the $V_{\pi\pi}$ quasi-potential by form factors such that

$$V_{\pi\pi}(E, \vec{p}; \vec{q}, \vec{q}') \to v(\vec{q}) \, V_{\pi\pi}(E, \vec{p}) \, v(\vec{q}') \qquad with \qquad v(\vec{q}) = g \left(1 + \frac{q^2}{q_d^2}\right)^{-\alpha} \qquad (4)$$

This modification does not destroy in any way the properties of chiral symmetry as is easily verified. The parameters g, q_d and α are fixed through a fit to the data which yields the following values : $g = 0.9, \quad q_d = 1.GeV, \quad \alpha = 3$. The corresponding phase shifts are shown in Fig.1. The power $\alpha = 3$ of the Yukawa form factor is somewhat unusual but it accounts for our neglect of u- and t-channel exchange contributions and also for the omission of couplings to the $K\overline{K}$ channel.

It is interesting to note that the T-matrix in eq.(1) can also be recast in the following form

$$T_{ab,cd}(s) = \delta_{ab}\delta_{cd} \frac{D_\pi^{-1}(s) - D_\sigma^{-1}(s)}{N\langle\sigma\rangle^2} \frac{D_\sigma(s)}{D_\pi(s)} , \qquad (5)$$

where s is the Mandelstam variable. $D_\pi(s)$ and $D_\sigma(s)$ are respectively the full pion and sigma propagators, while $\langle\sigma\rangle$ is the sigma condensate. The expression in eq.(5) is in fact a Ward identity which links the $\pi\pi$ four-point function to the π and σ two-point functions as well as to the σ one-point function. To this order, the pion propagator and the sigma-condensate are obtained from the Hartree-Bogoliubov (HB) approximation [7]. In terms of the pion-mass m_π and decay constant f_π, they are given by

$$D_\pi(s) = \frac{1}{s - m_\pi^2}, \qquad \langle\sigma\rangle = \frac{1}{\sqrt{N}} f_\pi. \qquad (6)$$

The sigma meson, on the other hand, is obtained from a Random Phase Approximation (RPA) involving
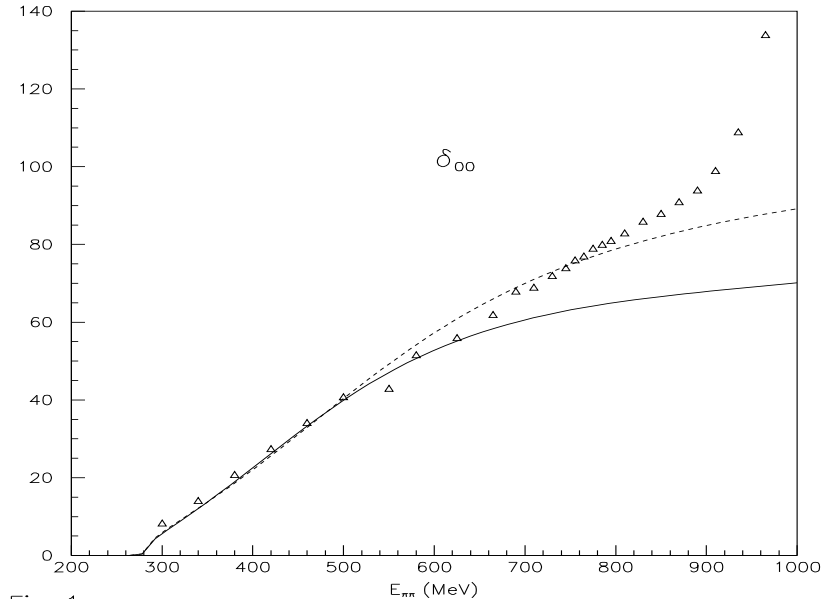
Fig. 1.

Figure 1: The s-wave phase shifts for the $\pi\pi$ scattering. Besides the data points, the full line denotes the leading-order result while the dashed line includes t- and u-channel corrections.

$\pi - \pi$ scattering [7]. It reads

$$D_\sigma(s) = \left[s - \mathcal{E}_\sigma^2 - \frac{2\lambda^4 \langle \sigma \rangle^2 \, \Sigma_{\pi\pi}(s)}{1 - \lambda^2 \Sigma_{\pi\pi}(s)} \right]^{-1}, \tag{7}$$

where λ^2 is the bare coupling and $\mathcal{E}_\sigma$ is the mean-field sigma mass (mass of the quasi-sigma) given in terms of the coupling constant, the condensate and the pion mass by:

$$\mathcal{E}_\sigma^2 = m_\pi^2 + 2\lambda^2 \langle \sigma \rangle^2.$$

It is clear from what was said above that the σ-meson propagator in this approach is correctly defined since it satisfies a hierarchy of Ward identities.

We now proceed to put the sigma meson in cold nuclear matter. We stress again that chiral symmetry is fully preserved. The pion is coupled to Δ-h and p-h channels including nuclear short-range correlations simulated by the Migdal parameters g' as well as in-elasticities coming for instance from the coupling to $2p - 2h$ states (see [3]). The result for various densities is shown in Fig.2. We see that, as density increases, a strong downward shift of the sigma-mass distribution occurs. However, contrary to earlier phenomenological models with no repulsion below threshold, the invasion of strength below the $E < 2m_\pi$ threshold region is still strong but saturates at around $1.5m_\pi$ as density increases. On the other hand, we also see that the corresponding imaginary part of the T-matrix is less modified.

Before discussing the possible relevance of this result for the strongly target-mass dependent threshold

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enhancement of the $\pi\pi$ invariant mass spectrum in recent $\pi, 2\pi$ experiments off nuclei, let us briefly discuss some effects which could go in the opposite direction. So far, we have only considered self-energy corrections to the pions. To be consistent we should consider on the same footing vertex corrections which usually go in the opposite direction to self-energy effects. Also more care should be given to the Pauli blocking in matter.
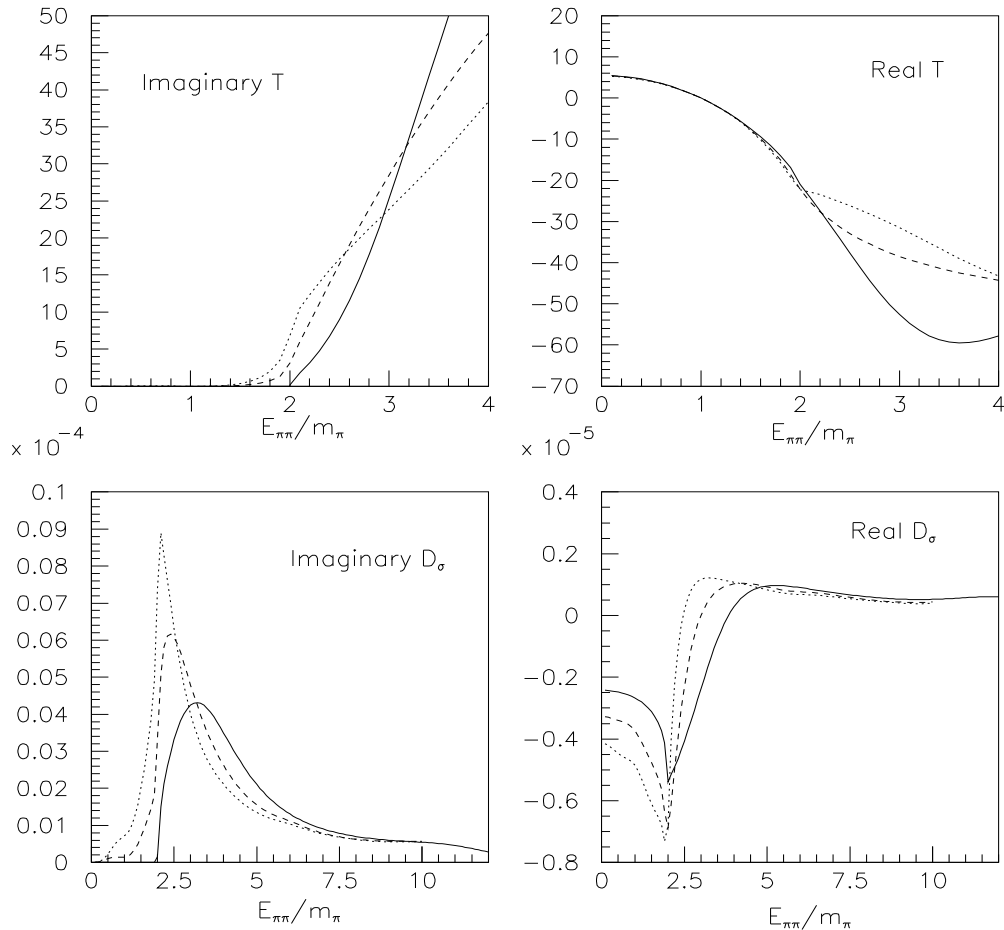

Figure 2: Results for the leading-order dynamics. The upper left (right) curves denote the imaginary (real) part of the $\pi\pi$-T-matrix. The lower left (right) curves denote the imaginary (real) parts of the sigma propagator. In all cases, the solid-line, the dashed-line and the dotted-line curves are respectively for the vacuum case, the medium at nuclear density $0.5\rho_0$ and at normal density.

Indeed, even though the in-medium renormalisation of the single pion is well established through the extensive analysis of the π-nucleus optical potential, the in-medium renormalisation of correlated pion-pairs requires an extra piece which is the Pauli exchange contributions to the pion p-h and Δ-h self energies. However, we do not think that these vertex corrections will completely cancel the effect of accumulation

of strength in the threshold region as matter density increases.

Again we want to mention that in the nuclear medium we have to be careful not to consider the sigma isolated from the correlated pions for instance in their interaction with nucleons. This, in connection with chiral symmetry, will be the subject of a separate study. We only want to demonstrate here that, for the sigma-meson mass distribution alone, the in-medium renormalisation is very strong. On the contrary in the $\pi - \pi$ T-matrix the effect at threshold is very much suppressed because of form factors which reflect the fact that the on-shell T-matrix at threshold has to go to zero in the chiral limit(see eq. (1)). In spite of this there exists, even for the T-matrix, a strong reshaping and in particular there is considerable strength invading the region below the $2m_\pi$ threshold. It is therefore indeed quite tempting to associate our finding with the strong strength accumulation found in recent $\pi, 2\pi$ knock-out experiments off nuclei by Grion et al.[8, 9]. However in these experiments the sigma-meson mass distribution is not measured directly but rather the imaginary part of the $\pi\pi$ T-matrix.

It is clear, from the phase-shift slope, that the s-wave $\pi\pi$ scattering is attractive. The Weinberg scattering length are known to be $a_0^0 = \frac{7}{32\pi}\frac{m_\pi}{f_\pi^2}$. However, the leading-order contributions of the $1/N$-expansion, as described above, leads to too much attraction. The scattering lengths, in the case of the three physical charges ($N = 3$), are given by

$$a_0^0 = \frac{9}{32\pi}\frac{m_\pi}{f_\pi^2}.$$

Therefore, the t- and u-channel contributions are ultimately needed. They are known to yield repulsion at threshold and enter as next-to-leading-order corrections. Before this has been worked out consistently we adopt, for the time being, a more phenomenological quasi-potential picture and use a Lippmann-Schwinger equation as a unitarization procedure of the full tree-level $\pi\pi$ scattering amplitude

$$
\begin{aligned}
V_{ab,cd}(s,t,u) &= \delta_{ab}\delta_{cd}A(s) + \delta_{ac}\delta_{bd}A(t) + \delta_{ad}\delta_{bc}A(u) \ , \\
A(s) &= \frac{m_\sigma^2 - m_\pi^2}{f_\pi^2}\frac{s - m_\pi^2}{s - m_\sigma^2} \ .
\end{aligned}
\tag{8}
$$

The functions $A(t)$ and $A(u)$ are obtained from $A(s)$ by substituting respectively t and u for s. This, indeed, corrects for the threshold and low-energy region and also allows for a reasonable description of the phase shifts. For the three-dimensional reduction of the Mandelstam variables t and u we choose an on-shell one : $t = (\omega_q - \omega_{q'})^2 - (\vec{q} - \vec{q}')^2$ and $u = (\omega_q - \omega_{q'})^2 - (\vec{q} + \vec{q}')^2$. Furthermore the t and u variables will be neglected (in comparison to the sigma mass) in the denominator of $A(t)$ and $A(u)$, respectively This leads to an analytically solvable integral equation which will account for the full leading-order contribution of eq. (1) and will include a piece of the vertex renormalisation coming from the t-and u-channels. In fact, by this phenomenological procedure, the $A(t)$ and $A(u)$ pieces play a role, similar to the Migdal

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parameters for the pion self-energy in the nuclear medium, which account for the repulsion present in the vertex corrections. This modification still preserves the symmetry conserving properties of the iterated T-matrix, as is easily verified. As a regularisation of the divergent integrals, we supplement the T-matrix in eq. (5) with a monopole-like form factor with an adjustable cutoff parameters Q_d. For a bare sigma mass, $m_\sigma = 1.GeV$, and a cutoff, $Q_d = 8m_\pi$, one gets a reasonable description of the phase shifts up to $800 MeV$ (Fig.1). The corresponding imaginary part of the T-matrix is shown in Fig.3.
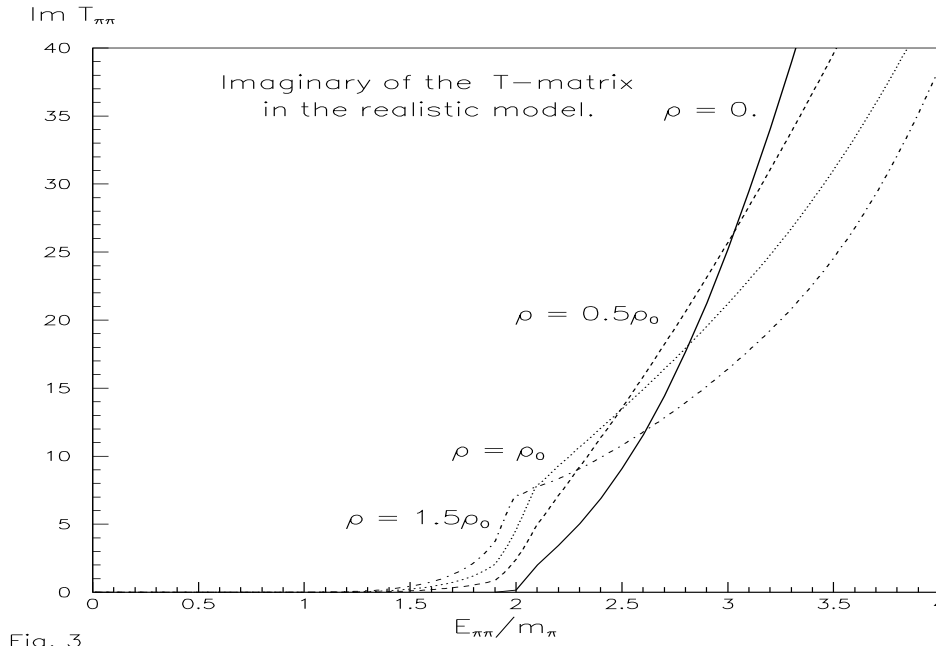

Figure 3: The imaginary part of the T-matrix with t- and u-channel corrections. The in-medium modifications are accounted for by means of a realistic model.

As a first attempt to explain with our theory the findings of Bonutti et. al. [9] we consider the scenario of quasi-free scattering. That this process dominates the pion knock-out reaction has been demonstrated experimentally in measuring a proton in coincidence with the pions [10]. The elementary process has actually been calculated by Oset and Vicente-Vacas [11] using a perturbative treatment at the tree level. For the invariant $\pi\pi$ mass , this theory explains reasonably well the ^2H data (see [11]). However, the mass dependence of the cross section and notably the resonant structure at threshold, growing with target mass, remains totally unexplained. As an attempt to include in-medium effects in this process, the authors of ref. [11] considered the in-medium modification of the one-pion-exchange pole-graph which was found to be negligible. It is clear from what was said above that the crucial region where one should look for possibly important in-medium effects is the $2m_\pi$ threshold. On the other hand this region seems to be singled out by the in-medium $\pi\pi$ correlation processes. Therefore it is rather natural to consider those as a final state interaction effects in describing the $(\pi, 2\pi)$ knock-out data.

Therefore, as a first step, we propose to include these final state interactions using our model for the in-medium $\pi\pi$ correlations, in the analysis performed by Oset[*] et al. in ref.[11]. This will be done in a very schematic way in order to obtain some insight as to whether or not our theory can at all explain the features found in these pion knock-out reactions. We thus replace in the analysis of ref.[11] the tree-level four-pion vertex by our medium-modified $\pi\pi$-T-matrix which, as we have said, gives rise to the σ-meson mass distribution shown earlier. Furthermore, to keep the numerical calculation within reasonable limits, we use a rather crude model for the in-medium renormalisation of the pions. Since the latter are subject to particle-hole and delta-hole couplings with p-wave dominance, one can show by comparison with more realistic studies that, to a good approximation, the single-pion dispersion relation in matter can be modelled in the following way

$$\omega_\pi(\vec{q}) = \sqrt{m_\pi^2 + \vec{q}^2} \quad \rightarrow \quad \omega_\pi^\gamma(\vec{q}) = \sqrt{m_\pi^2 + \gamma\vec{q}^2} \ ,$$

where γ takes values from 0.8 to 0.4 depending on density. The results for the in-medium strength distribution in this toy-model approach is shown in Fig. 4 for typical values of the parameter γ.

By comparing the strength distributions from the full (realistic) model of Fig. 3 and those from the toy model in Fig.4 one clearly sees that the toy model is overestimating the effects near threshold. On the other hand, the toy model misses completely the subthreshold strength. In fact, in the full model, the presence of subthreshold p-h and 2p-2h cuts renders the peaks at threshold much broader than in the toy model. Therefore the strength which is accumulated at threshold in the case of the toy model is spread into the subthreshold region in the case of the full model.

At first glance, the use of the toy model to assess the in-medium effects on the $(\pi, 2\pi)$ knock-out reaction off nuclei, may seem rather erroneous due to this overestimation of the strength at threshold. However, one should keep in mind that a reaction theory for the knock-out process has to take into account the finite three-momenta of the emitted pion-pairs, whereas the calculation presented in Figs. and 4 are all done in a back-to-back kinematics (c.m frame). Finite three-momenta $\vec{P}$ of the pair induce in fact an interesting feature which we want to discuss now.

Consider two pions detected outside the nucleus at three-momenta $\vec{q}_1$, $\vec{q}_2$ (with $\vec{P} = \vec{q}_1 + \vec{q}_2$) and energies $\omega_\pi(\vec{q}_1) = \sqrt{m_\pi^2 + \vec{q}_1^2}$, $\omega_\pi(\vec{q}_2) = \sqrt{m_\pi^2 + \vec{q}_2^2}$, respectively. Inside the nucleus, on the other hand, this pair of pions have three-momenta $\vec{k}_1$, $\vec{k}_2$ and energies $\omega_\pi^\gamma(\vec{k}_1)$, $\omega_\pi^\gamma(\vec{k}_2)$ related by the in-medium dispersion

$$\omega_\pi^\gamma(\vec{k}_1) = \sqrt{m_\pi^2 + \gamma\vec{k}_1^2} \ , \quad \omega_\pi^\gamma(\vec{k}_2) = \sqrt{m_\pi^2 + \gamma\vec{k}_2^2} \ .$$

Knowing that the energy of each particle, in either the vacuum or the medium, has to be the same: $\omega_\pi(\vec{q}_1) = \omega_\pi^\gamma(\vec{k}_1)$ and $\omega_\pi(\vec{q}_2) = \omega_\pi^\gamma(\vec{k}_2)$, one can one can extract a scaling relation between the three-momenta of each particle inside and outside the nucleus

$$\vec{q}_1 = \sqrt{\gamma}\,\vec{k}_1 \ , \quad \vec{q}_2 = \sqrt{\gamma}\,\vec{k}_2 \ .$$

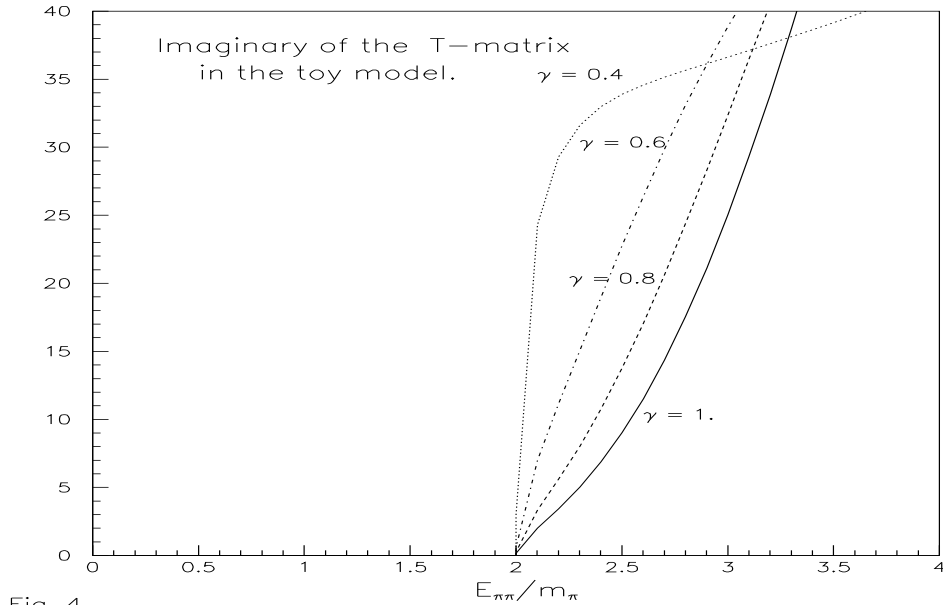[*]We are grateful to E. Oset and M. J. Vicente Vacas for providing to us with their code.

Fig. 4

Figure 4: The imaginary part of the T-matrix with t- and u-channel corrections. In-medium modifications are accounted for by means of a toy model.

Hence, the invariant mass $\tilde{M}_{\pi\pi}$ of a pion-pair inside the nucleus which is given by

$$\tilde{M}_{\pi\pi} = (\omega_\pi^\gamma(\vec{k}_1) + \omega_\pi^\gamma(\vec{k}_1))^2 - (\vec{k}_1 + \vec{k}_2)^2 ,$$

is related to the invariant mass, $M_{\pi\pi} = (\omega_\pi(\vec{q}_1) + \omega_\pi(\vec{q}_1))^2 - (\vec{q}_1 + \vec{q}_2)^2$, of the same pion-pair outside the nucleus through the relation :

$$\tilde{M}_{\pi\pi} = M_{\pi\pi} - \left(\frac{1}{\gamma} - 1\right) \vec{P}^2 .$$

One can see that, for the typical values taken by the γ parameter ($\gamma \leq 1$), the pion-pair, observed at the invariant mass $M_{\pi\pi}$, has actually a lower in-medium invariant mass $\tilde{M}_{\pi\pi} \leq M_{\pi\pi}$. This also suggests that, for finite three-momenta pion-pairs ($\vec{P} \neq \vec{0}$), one is in fact able to probe the strength of the $\pi\pi$ mass distribution at an in-medium invariant mass even below the $2m_\pi$ threshold. This interesting feature is not possible for the case of back-to-back pions even though each individual pion has a modified in medium dispersion relation owing to the factor γ. Therefore, since we presently do not possess a consistent reaction-theory at finite total three-momentum, it may be advantageous not to use the full realistic model of Fig. 3 in which the strength is totally spread out into the region below threshold but instead use the toy model which keeps the total strength concentrated at threshold. This option may eventually compensate for the inadequate back-to-back kinematics inherent in our model for in-medium $\pi\pi$ correlations. Of course, this argument certainly needs better support from a quantitative analysis at finite total three-momentum. This study is presently underway.

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Before concluding this short note, we present a numerical calculation going along the lines sketched above. To the theoretical analysis by Oset and Vicente Vacas of the $(\pi, 2\pi)$ knock-out reaction off nuclei, we add $\pi\pi$ final state interactions. This is done by replacing the tree-level $\pi\pi$-vertex considered in [11] by an in-medium renormalised $\pi\pi$ T-matrix. As indicated above this is only done in the framework of the economical toy model.
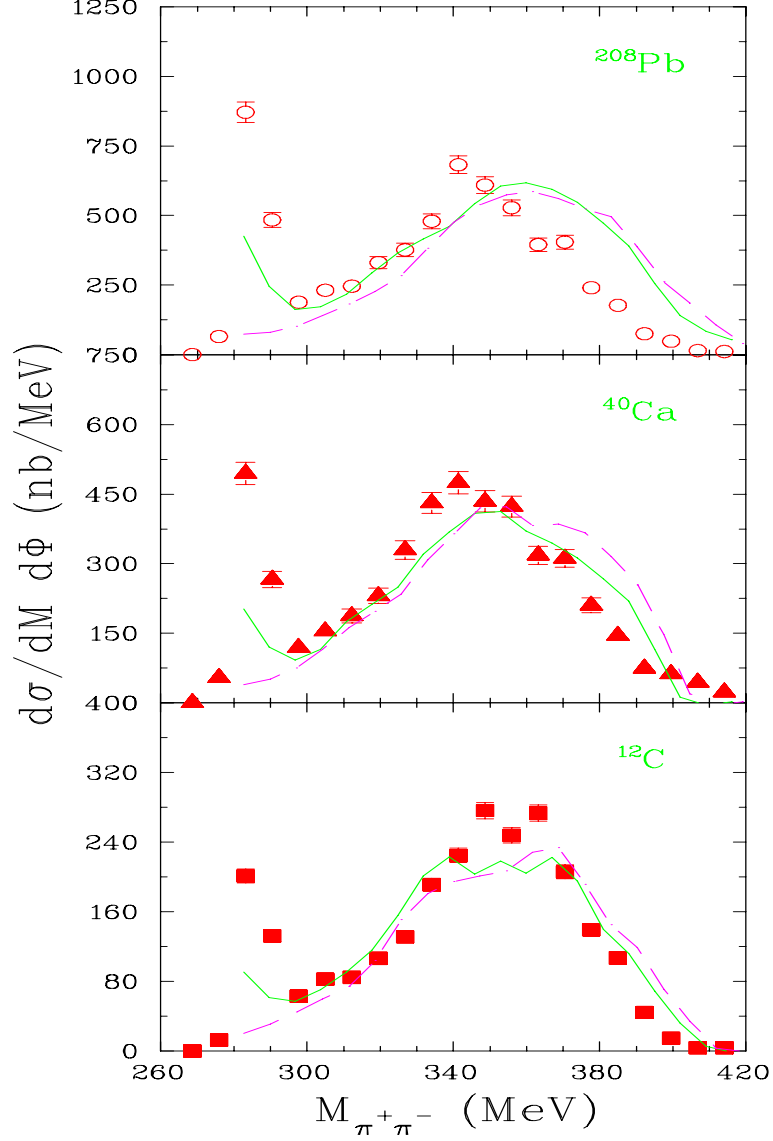

Figure 5: The $\pi^+\pi^-$ invariant-mass distribution in the $(\pi, 2\pi)$ knock-out reaction off Carbon, Calcium and Lead. Along with the experimental points curves from the theoretical calculations of Oset and Vicente Vacas [11] are displayed as dashed-dotted lines. The full lines include the in-medium final-state interaction in the $\pi\pi$ vertex. The in-medium modifications are taken into account by means of the toy model as discussed in the text.

We can see from the theoretical curves that the experimental findings can roughly be described with values of the γ parameter in the range of 0.6 to 0.8 (see figure 5). As discussed above, this essentially stems

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from the fact that the in-medium renormalisation of the pions induces an important downward shift of the strength in the $\pi\pi$ T-matrix. This analysis was made to illustrate the basic idea. It will be important to look carefully into vertex corrections which will certainly have a moderating effect. This remains to be seen in future work.

At the end let us emphasise that the low invariant $\pi - \pi$ mass accumulation considered here is of different origin than a similar effect obtained from $A(p, 2\pi)X$ reactions at SATURNE [10, 12]. There the incident proton energy was $1.6 GeV$ such that the pions have much higher energies than in the TRIUMF experiment. Because of that the pions in the SATURNE experiment have a much shorter mean free path, i.e. they are strongly absorbed by the medium, leading to a pronounced shadowing effect. The collinearity of the pions leaving the nucleus on the nuclear-matter-free side induces a low invariant-mass enhancement of the $\pi\pi$ cross section which increases with the size of the target nucleus. This scenario was confirmed in [10] from numerical BUU simulation but also experimentally from the fact that the cross section for heavier targets is maximum for small opening angles of the pions. In contrast, at TRIUMF, the outgoing pions are of low energy and therefore have a long mean free path which is confirmed by the fact that the cross section is flat as a function of the opening angles (see reference [8]). This kind of scenario is also individually confirmed by the mass dependence of the total $(\gamma, 2\pi)$ production cross section on nuclei which, as a function of mass, is totally flat for low energy γ's $(400 MeV)$ whereas it decreases for high energy photons (1 GeV). We therefore think that the experiment analysed here has nothing to do with nuclear shadowing and that the observed mass enhancement really signals a collective effect resulting in a strong downward shift of part of the $\pi^+ - \pi^-$-and σ-meson strength distribution.

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