The σ meson in a nuclear medium through two pion photoproduction

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Abstract

We present theoretical results for $(\gamma, \pi^0\pi^0)$ production on nucleons and nuclei in the kinematical region where the scalar isoscalar $\pi\pi$ amplitude is influenced by the σ pole. The final state interaction of the pions modified by the nuclear medium produces a spectacular shift of strength of the two pion invariant mass distribution induced by the moving of the σ pole to lower masses and widths as the nuclear density increases.

The nature of the σ meson has been the subject of continuous debate. Its nature as an ordinary $q\bar{q}$ meson or a $\pi\pi$ resonance has centered most of the discussion [1]. The advent of χPT has brought new light into this problem and soon it was suggested [2, 3] that the σ could not qualify as a genuine meson which would survive in the limit of large N_c. The reason is that the $\pi\pi$ interaction in s-wave in the isoscalar sector is strong enough to generate a resonance through multiple scattering of the pions. This seems to be the case, and even in models starting with a seed of $q\bar{q}$ states, the incorporation of the $\pi\pi$ channels in a unitary approach leads to a large dressing by a pion cloud which makes negligible the effects of the original $q\bar{q}$ seed [4]. This idea has been made more quantitative through the introduction of the unitary extensions of χPT ($U\chi PT$) [5, 6, 7, 8]. These works implement unitarity in coupled channels in an exact form and use the input of the lowest and second order chiral Lagrangians of [9]. The inverse amplitude method is used in [7] and an expansion in powers of $O(p^2)$ is done for the real part of the inverse of the scattering amplitude, while in [8] the dynamics of the second order chiral Lagrangian is introduced via the explicit use of the exchange of genuine mesons, following the lines of [10], and unitarizing with the N/D method. These works also justified the success in the scalar sector of the Bethe Salpeter approach used in [11, 12]. In all these cases the σ meson appears as a pole of the $\pi\pi$ scattering amplitude in the second Riemann sheet, even when the second order Lagrangian, which contains information of the exchange of genuine mesons according to [10], is neglected. These unitary models have been tested with success in many elementary reactions [13].

Another point of interest which can help us to understand the nature of the σ meson is the modification of its properties at finite nuclear density. The importance of the medium modification of the $\pi\pi$ interaction in the scalar sector was suggested in [14] where the $\pi\pi$ amplitude in the medium developed large peaks below the two pion threshold, somehow indicating that the σ pole had moved to much lower energies. The issue has been revised and the models have been polished incorporating chiral constraints [15, 16, 17] with the result that the peaks disappear at normal density, but still much strength is shifted to low energies.

Since present theoretical calculations agree on a sizeable modification in the nuclear medium of the $\pi\pi$ scattering in the σ region, our purpose here is to find out its possible experimental signature in a very suited process like the $(\gamma, \pi^0\pi^0)$ reaction in nuclei. Recent experiments at Mainz [18], where preliminary results show a very clear shift of strength of the invariant mass distribution of the two pions towards low invariant masses, seem to indicate that medium effects are indeed large.

In the present paper we shall show how the $\pi\pi$ interaction in the final state of the $(\gamma, \pi^0\pi^0)$ reaction on nucleons enhances the cross section of this reaction and how the medium corrections on the $\pi\pi$ interaction in nuclei lead to an appreciable shift of the strength of the invariant mass distribution towards lower invariant masses. Although a similar shift has been claimed in the $(\pi, \pi\pi)$ reaction in nuclei [19, 20, 21], the fact remains that there are still some discrepancies between these experiments and that the theoretical calculations [22, 23] do not reproduce the data. The reason is that the $(\pi, \pi\pi)$ reaction, involving initial pions which are much distorted in the nucleus, is quite peripheral and the effective densities tested are small. A possible way out to reconcile theory and experiment was suggested in [23], showing that the small cross section in $\pi^- p \to \pi^+\pi^- n$ at small invariant masses was abnormally small because of a subtle cancellation of large terms. A medium modification of these terms, through changes in the N^* properties and others not having to do with the $\pi\pi$ interaction, could offset that cancellation and lead to larger final results. The $(\gamma, \pi^0\pi^0)$ reaction is much better suited to investigate the modification of the $\pi\pi$ interaction in the medium because the photons are not distorted and one can test larger nuclear densities.

For the model of the elementary $(\gamma, \pi\pi)$ reaction we follow [24] which considers the coupling of the photons to mesons, nucleons, and the resonances $\Delta(1232)$, $N^*(1440)$, $N^*(1520)$ and $\Delta(1700)$. In the region of relevance to the present work, $E_\gamma = 400 - 460$ MeV, apart from some background terms, the Δ Kroll Ruderman term, diagram $i)$ of Fig. 1 in [24], is of importance and will be dealt with separately from the rest. The model of [24] relies upon tree level diagrams. Final state interaction of the πN system is accounted for by means of the explicit use of resonances with their widths. However, since we do not include explicitly the σ resonance, the final state interaction of the two pions has to be implemented to generate it.

The $\gamma N \to N\pi^0\pi^0$ amplitude can be decomposed in two parts: the one that has as final state the combination of pions with isospin I=0, first term of the RHS of Eq. (1), and the I=2 combination, last two terms of Eq. (1).

$$|\pi^0(1)\pi^0(2)>=\underbrace{\frac{1}{3}|\pi^0(1)\pi^0(2)+\pi^+(1)\pi^-(2)+\pi^-(1)\pi^+(2)>}_{\text{I=0 part}}$$
$$\underbrace{-\frac{1}{3}|\pi^0(1)\pi^0(2)+\pi^+(1)\pi^-(2)+\pi^-(1)\pi^+(2)>+|\pi^0(1)\pi^0(2)>}_{\text{I=2 part}} \quad (1)$$

The interaction of pions in I=2 in s-wave at these energies is very weak and hence we do not modify this part of the $\gamma N \to N\pi^0\pi^0$ amplitude due to the final state interaction of the pions. However, the I=0 part is strongly modified. We have also checked that at the low energies involved here the pions come essentially in s-wave.

In ref. [23] the renormalization of the $I = 0$ $(\pi, \pi\pi)$ amplitude was done by factorizing the on shell tree level $\pi N \to \pi\pi N$ and $\pi\pi \to \pi\pi$ amplitudes in the loop functions. This was justified in [11] for the $\pi\pi$ interaction. The same approach would lead in our case to

$$T^{I_{\pi\pi}=0}_{(\gamma,\pi^0\pi^0)} \to T^{I_{\pi\pi}=0}_{(\gamma,\pi^0\pi^0)}\left(1+G_{\pi\pi}t^{I=0}_{\pi\pi}(M_I)\right) \quad (2)$$

where $G_{\pi\pi}$ is the loop function of the two pion propagators, which appears in the Bethe Salpeter equation, and $t^{I=0}_{\pi\pi}$ is the $\pi\pi$ scattering matrix in isospin I=0. In order to show clearly how one is lead to this equation, we show in Fig. 1 the diagrams involved in the two pion production including their final state interaction.

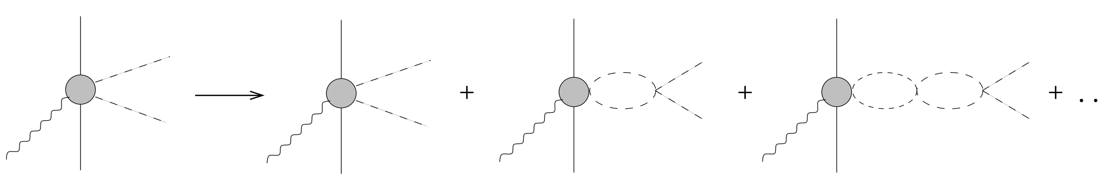

Figure 1: Diagrammatic series for pion final state interaction in I=0

The two final pions undergo multiple scattering which can be accounted for by means of the Bethe Salpeter equation,

$$t = V + VG_{\pi\pi}t \quad (3)$$

where, following [11], V is given by the lowest order chiral amplitude for $\pi\pi \to \pi\pi$ in $I = 0$ and $G_{\pi\pi}$, the loop function of the two pion propagators, is regularized by means of a cut off in [11], or alternatively with dimensional regularization in [8]. In both approaches it was shown that V factorizes with its on shell value in the Bethe-Salpeter equation of Fig. 1. Hence, in the Bethe-Salpeter equation the integral involving Vt and the product of the two pion propagators affects only these latter two, since V and t factorize outside the integral, thus leading to Eq. (3) where $VG_{\pi\pi}t$ is the algebraic product of V, the loop function of the two propagators, $G_{\pi\pi}$, and the t matrix. Coming back to Fig. 1 it is now clear that in case the vertex from where the two pions emerge is a contact term with a constant amplitude, the series implicit in Fig. 1 is summed as

$$\begin{aligned} &T + TG_{\pi\pi}V + TG_{\pi\pi}VG_{\pi\pi}V + \ldots = \\ &= T[1 + G_{\pi\pi}(V + VG_{\pi\pi}V + \ldots)] = T(1 + G_{\pi\pi}t_{\pi\pi}) \end{aligned} \tag{4}$$

which is Eq. (2). Now in the model for $(\gamma, 2\pi)$ of [24] there are indeed contact terms as implied before, as well as other terms involving intermediate nucleon states or resonances. In this latter case the first loop function in the diagrams of Fig. 1 is more complicated since it involves three propagators. Yet, if the intermediate baryon is far off shell, as is the case for most diagrams, then the baryon propagator does not change much in the loop function and the factorization of Eq. (2) still holds. There is, however, an exception in the Δ Kroll Ruderman term, since as we increase the photon energy we get closer to the Δ pole. For this reason, and because its weight is important at these energies in the $(\gamma, \pi^+\pi^-)$ amplitude which is needed in Eq. (1), we have singled out this term which we work out in detail below. The I=0 part of the amplitude requires the combination of the diagrams shown in Fig. 2 a) b)

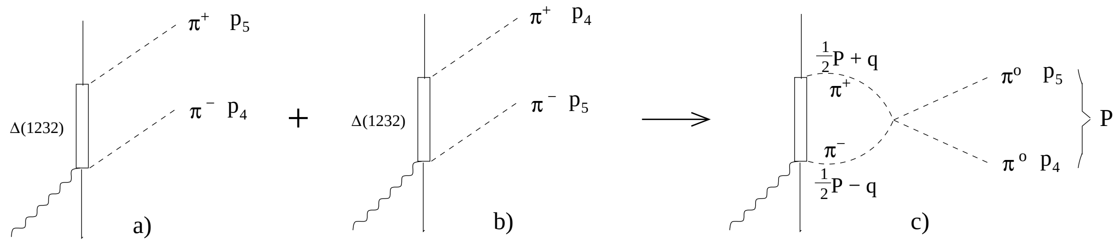

Figure 2: Diagrammatic series for the Δ Kroll Ruderman term with final pions in I=0.

and hence the vertex contribution is (assuming the Δ propagator, G_Δ, the same in both cases)

$$G_\Delta(\vec{S}\vec{p}_5\ \vec{S}^\dagger\vec{\epsilon} + \vec{S}\vec{p}_4\ \vec{S}^\dagger\vec{\epsilon}) = G_\Delta\vec{S}(\vec{p}_4 + \vec{p}_5)\ \vec{S}^\dagger\vec{\epsilon} = G_\Delta\vec{S}\vec{P}\ \vec{S}^\dagger\vec{\epsilon} \tag{5}$$

where $\vec{S}^\dagger$ is the spin transition operator from spin 1/2 to 3/2 and $\vec{\epsilon}$ is the photon polarization.

In the loop function originated from the sum of the terms a) and b) depicted in Fig. 2, assuming also the Δ propagator constant, we would have the contribution

$$i\int \frac{d^4q}{(2\pi)^4} 2\ \vec{S}(\frac{\vec{P}}{2}+\vec{q})\ \vec{S}^\dagger\vec{\epsilon}\ G_\Delta \frac{1}{(\frac{1}{2}P+q)^2-m_\pi^2+i\epsilon}\ \frac{1}{(\frac{1}{2}P-q)^2-m_\pi^2+i\epsilon}\ t_{\pi\pi} = $$
$$= G_\Delta\ \vec{S}\vec{P}\ \vec{S}^\dagger\vec{\epsilon}\ G_{\pi\pi}(s)\ t_{\pi\pi}(s) \tag{6}$$

since the term proportional to $\vec{S}\vec{q}$ vanishes because of parity reasons, and then we see explicitly the factorization of the tree level amplitude. If we keep explicitly the Δ propagator in the loop some corrections arise since now we have the loop function with the Δ propagator and two pion propagators. This loop integral is performed following the steps of [25], where the same loop with three propagators, albeit with only strong vertices, is evaluated. By performing explicitly the q^0 integration in Eq. (6) we obtain

$$\int \frac{d^3q}{(2\pi)^3} 2\ \vec{S}(\frac{\vec{P}}{2}+\vec{q})\ \vec{S}^\dagger\vec{\epsilon}\ \frac{M_\Delta}{E_\Delta}\frac{1}{2\omega\omega'}\ \frac{1}{P^0+\omega+\omega'}\ \frac{1}{P^0-\omega-\omega'+i\epsilon}$$
$$\cdot \frac{1}{\sqrt{s}-\omega'-E_\Delta+i\frac{\Gamma_\Delta(p_\Delta^2)}{2}}\ \frac{1}{\sqrt{s}-\omega-P^0-E_\Delta+i\frac{\Gamma_\Delta(p_\Delta^2)}{2}}$$
$$\cdot \Big[(\omega+\omega')(\sqrt{s}-E_\Delta-\omega-\omega')-\omega P^0\Big] \tag{7}$$

where $E_\Delta = \sqrt{{M_\Delta}^2+{\vec{p}_\Delta}^2}$, $\omega = \sqrt{m_\pi^2+(\frac{\vec{P}}{2}+\vec{q})^2}$, $\omega' = \sqrt{m_\pi^2+(\frac{\vec{P}}{2}-\vec{q})^2}$, $p_\Delta = \frac{\vec{P}}{2}-\vec{q}$, $p_\Delta^0 = \frac{P^0+\omega+\omega'}{2}$ and $\sqrt{s}$ is the CM energy of the initial photon and nucleon.

The improved loop calculation results in a 10 per cent reduction in the cross section. The pion pole term (diagram j of Fig. 1 of ref. [24]) is essential in the $(\gamma,\pi\pi)$ model to guarantee gauge invariance, but is numerically negligible at the energies which we have here since it is proportional to $\vec{P}^{\,2}$, while the Δ Kroll-Ruderman term i) is proportional to $\vec{P}$. Yet, inside the loops, the contribution can be bigger since the $\gamma\pi\pi$ vertex now is proportional to the loop momentum and not to the external pion momentum. In addition there is an extra p-wave $\pi N\Delta$ vertex which also goes like the loop momentum. However, there is also an extra pion propagator and one also finds cancellations from the poles of the different propagators. This was done explicitly in [26] where moderate effects from this pion pole term were obtained. We have performed the appropriate calculation and found that the loop diagram from the pion pole term is of the order of 10 per cent of the corresponding loop with the Δ Kroll-Ruderman term. Hence, we simply take care of it through the factorization approximation. We also take advantage of this numerical calculation to include two extra baryon form factors in the loop, as done in [25], to account for the $\pi N\Delta$ vertex correction. This is also done for the other diagrams since πNN vertices are also involved. The inclusion of the form factors leads to a further reduction of about 15 per cent in the final cross section. Hence, the actual

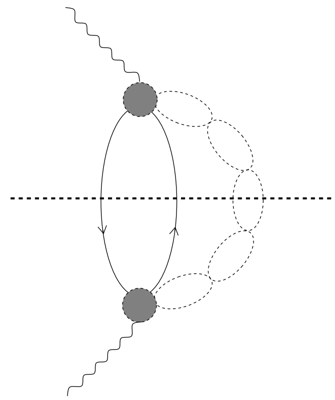

Figure 3: Selfenergy diagram for the evaluation of the cross section including the final pions rescattering.

modification of the amplitude should be

$$\begin{aligned} T \to \quad & T - T^{00} + \left(T^{00} - T^{00}_{KR}\right)(1 + G_{\pi\pi} t_{\pi\pi}) + T^{00}_{KR} + G_{KR} t_{\pi\pi} \\ & = T + \left(T^{00} - T^{00}_{KR}\right) G_{\pi\pi} t_{\pi\pi} + G_{KR} t_{\pi\pi} \end{aligned} \tag{8}$$

where T is the full $(\gamma, \pi^0\pi^0)$ amplitude, T^{00} is the $(I = 0, I_3 = 0)$ amplitude for $(\gamma, \pi^0\pi^0)$, T^{00}_{KR} is the same as T^{00} but calculated only with the Δ Kroll Ruderman term, $G_{\pi\pi}$ is the two pions loop including the strong $\pi BB'$ form factors, $t_{\pi\pi}$ is the $\pi\pi$ scattering matrix in isospin I=0, G_{KR} is the loop of Fig. 2 including in the integrand the two form factors and the T^{00}_{KR} that depends on the momentum in the loops. (This is, Eq. (7) including the form factors and the coupling constants and isospin coefficients of [24]).

After all this is done, we find it technically useful, in order to account for these elaborate loop corrections, to still apply the factorization of the $(\gamma, 2\pi)$ tree level amplitude of Eq. (2) but with a slightly modified form factor included in the $G_{\pi\pi}$ loop function. The cut off of the monopole form factor is changed from 1 GeV to 625 MeV to implement these changes, including also the small effects from the projection in the two pion s-wave channel. This procedure is quite accurate numerically and prevents the numerical task from blowing up when we perform the calculations in nuclei.

There is also a small technical detail. One of the terms in our approach contains the Roper excitation and its posterior decay into two pions in s-wave. The Lagrangian used is given in [23] and the coupling constant is renormalized, as in [23], in such a way that when the $\pi\pi$ final state interaction is taken into account the empirical N^* width is obtained.

The cross section for the nuclear process can be calculated using many body techniques in a similar way to [27], [28]. The method proceeds in two steps. In the first one the probability per unit length of a $(\gamma, 2\pi)$ process in nuclear matter is evaluated from the imaginary part of the photon selfenergy diagram of Fig. 3 corresponding to the cut of the horizontal line. By using the local density approximation and, hence, substituting the density of nuclear matter by the empirical density of the nucleus $\rho(\vec{r})$ at a point $\vec{r}$, this provides the reaction probability in this point of the nucleus. The second step requires to follow the individual particles produced through the nucleus in an inclusive process where the nucleus can be excited to any state. This second step is done by using semiclassical methods in which the pions produced follow classical trajectories and are allowed to undergo quasielastic collisions or pion absorption according to probabilities calculated previously using many body techniques. As shown in [28], this procedure was very successful in describing pion nucleus phenomenology at intermediate energies.

With these ingredients the nuclear cross section is given by

$$
\begin{aligned}
\sigma = \quad & \frac{\pi}{k}\int d^3\vec{r}\int \frac{d^3\vec{p}}{(2\pi)^3}\int \frac{d^3\vec{q_1}}{(2\pi)^3}\int \frac{d^3\vec{q_2}}{(2\pi)^3}\frac{1}{2\omega(\vec{q_1})}\frac{1}{2\omega(\vec{q_2})} \\
& \cdot \sum_{s_i,s_f}\overline{\sum_{pol}} \mid T\mid^2 n(\vec{p})\; [1-n(\vec{k}+\vec{p}-\vec{q_1}-\vec{q_2})] \\
& \cdot\, \delta(k^0+E(\vec{p})-\omega(\vec{q_1})-\omega(\vec{q_2})-E(\vec{k}+\vec{p}-\vec{q_1}-\vec{q_2})) \\
& \cdot\, F_1(\vec{r},\vec{q_1})\; F_2(\vec{r},\vec{q_2})
\end{aligned}
\tag{9}
$$

where $n(\vec{p})$ is the occupation number for a density $\rho(\vec{r})$. The factors $F_i(\vec{r},\vec{q_i})$ take into account the distortion of the final pions in their way out through the nucleus and are given by

$$
\begin{aligned}
F_i(\vec{r},\vec{q_i}) &= exp\left[\int_{\vec{r}}^{\infty} dl_i\;\frac{1}{q_i} Im\Pi(\vec{r_i})\right] \\
\vec{r_i} &= \vec{r}+l_i\;\vec{q_i}/\mid\vec{q_i}\mid
\end{aligned}
\tag{10}
$$

where Π is the pion selfenergy, taken from [29], where the interaction of low energy pions with nuclear matter was studied. This potential has been tested against the different pionic reaction cross sections, elastic, quasielastic and absorption. The imaginary part of the potential is split into a part that accounts for the probability of quasielastic collisions and another one which accounts for the pion absorption. As we shall see, the probability that there is loss of pion flux through pion absorption at low energies is larger than through quasielastic collisions. One of the reasons is the Pauli blocking of the occupied states.

When we renormalize the I=0 amplitude to account for the pion final state interaction, we change $G_{\pi\pi}$ and $t_{\pi\pi}^{I=0}$ by their corresponding results in nuclear matter [17] evaluated at the local density of the point $\vec{r}$ in the integral of Eq. (9). We do not include here the direct coupling of the two pions to the $N^*(1440)h$ which was found extremely small in [30]. We do not include either the direct coupling of the two pions to a $p-h$ excitation. The s-wave would proceed only through the tiny isoscalar $\pi N \rightarrow \pi N$ interaction, and the p-wave

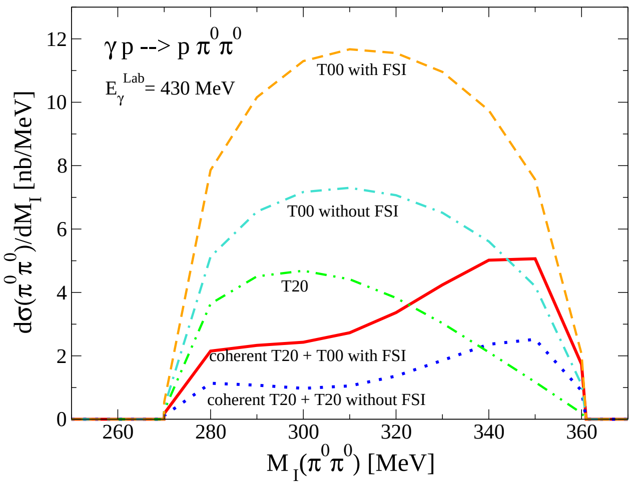

Figure 4: Contributions of the different isospin channels to the 2π mass distribution.

part, weak in principle at the relatively small momenta of the pions, would be only relevant at excitation energies of about 30-50 MeV, far below the region of interest to us in the present problem. On the other hand, one should note that other alternatives have been proposed to study the modification of the isoscalar $\pi\pi$ interaction in the nuclear medium based on the reduction of the pion decay coupling constant in the medium, which is tied to the dropping of the quark condensate via the Gell-Mann–Oakes–Renner relation [31, 32], or a combination of this effect and the dressing of the pion in the medium [33]. However, in the approach which we follow, one must be cautious not to include the dropping of f on top of the many body corrections done. Indeed, the change of f as a function of density in [31, 34] and related works, obtained through the GOR relation, is linked to the renormalization of the time component of the axial current in nuclei, but in standard many body theory these currents are renormalized using the same lagrangians as in the evaluation of the pion selfenergy [35, 36, 37, 38]. Therefore, one should not modify f in these latter approaches to avoid double counting if using an explicit perturbative calculation with effective chiral lagrangians. Actually, a large change of f at normal nuclear density would be difficult to accommodate to the quite well known pion nucleus phenomenology (pionic atoms, pion absorption, etc). Concerning this point it is worth noting the discussion in [39, 40, 41] about the difference between the $f(\rho)$ constant related to the quark condensate and the f used in perturbative calculations with chiral lagrangians. These differences were also stressed in [42], where it was also shown that taking the $f(\rho)$ coupling related to the quark condensate in the many body evaluation of the pion nucleus optical potential led to unacceptably large widths of pionic atoms.

We have also used the complex Δ selfenergy from [43] to dress the Δ propagator. In addition to the proper real part of the selfenergy in [43] we add the effective contribution to the selfenergy $4/9(f^*/\mu)^2 g'\rho$ coming from the iterated Δh excitation driven by de Landau Migdal interaction [44, 43].

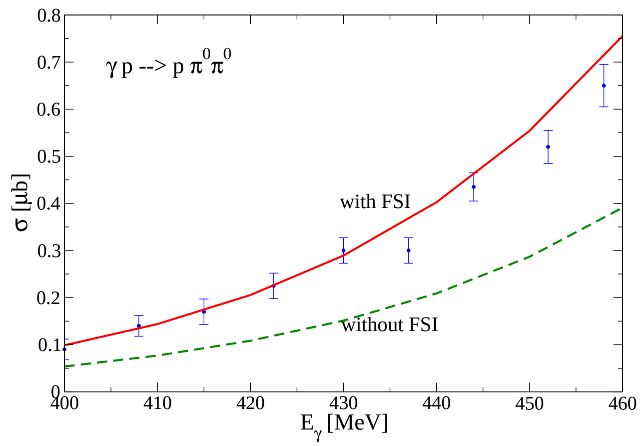

Figure 5: Total cross section for $\gamma p \to \pi^0\pi^0 p$ with and without pion final state interaction. Experimental data from ref. [46].

In Fig. 4 we show the results of the invariant mass distribution of the two pions for the $\gamma p \to \pi^0\pi^0 p$ reaction. We can see the contribution of the I=0 part alone (T00), the I=2 part alone (T20) and the coherent sum of the two, both with and without renormalization (FSI) of the I=0 amplitude. The renormalization of the I=0 amplitude has important effects, nearly doubling the cross section. This is reminiscent of the similar enhancement found from chiral loops at threshold in [45]. When we sum coherently the I=0 and I=2 amplitudes the shape of the distribution exhibits a double hump, one at low invariant masses and the other one at the high mass part of the spectrum. This shape is corroborated by the preliminary experimental results of [18].

The integrated cross section compared with experiment can be seen in Fig. 5 where we can appreciate that the inclusion of final state interaction improves the agreement with the data. The amplitude for the $(\gamma, \pi^0\pi^0)$ reaction on neutrons is calculated along the same lines and leads to a similar mass distribution as that of Fig. 4 albeit with a little smaller cross section.

In Fig. 6 a) and b) we show the cross section for ^{12}C and ^{208}Pb assuming the FSI of the pions in the free case ($\rho = 0$), and we show the results without pion absorption or quasielastic steps, with pion absorption alone, and a third curve which corresponds to the case where the pions which undergo quasielastic collisions together with those absorbed are eliminated from the outgoing pion flux. This is actually not the case but the comparison of the two lower curves gives us a measure of the amount of quasielastic collisions undergone by the pions in their way out from the production point. The figures show that more pions undergo absorption than quasielastic collisions. The larger part of the quasielastic collisions would not change the charge of the pions, only their energy and momentum would be changed. In this case the two π^0 would still be there and their invariant mass would be somewhat changed, sometimes leading to larger two pion invariant masses and other to smaller ones. Hence, as an average the invariant mass distribution of the two pions should not be much

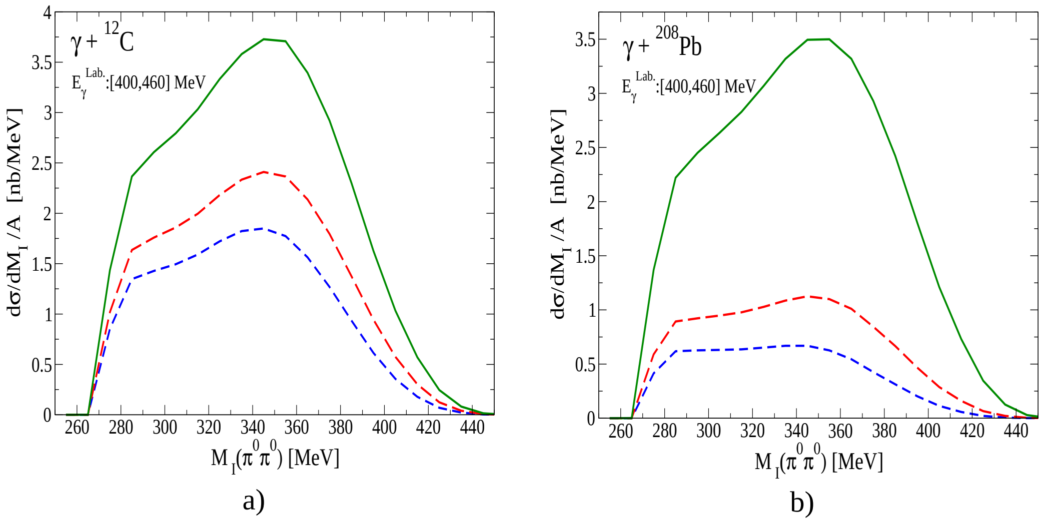

Figure 6: Two pion invariant mass distribution for $2\pi^0$ photoproduction in ^{12}C and ^{208}Pb. All three curves are calculated using the 2π final state interaction at 0 density but they differ in the final pion distortion: continuous line: without absorption nor quasielastic scattering. Long dashed line: only final pions absorption. Short dashed line: final pions absorption and quasielastic scattering.

modified by these quasielastic collisions. In other cases there could be change of charge and then we would not have two π^0 in the final state. However, this could also be compensated by having originally a $\pi^+\pi^0$ production followed by a collision of the π^+ with charge exchange. The distortion of the final pions has been done simply by using the imaginary part of the optical potential and leads to a reduction of the cross section. The results for the case of removal of only the pions which are absorbed are obtained by putting in the imaginary part of the pion selfenergy, Π, in Eq. (10), the part which comes from the absorption and omitting the one that comes from quasielastic, which have been separated in [29]. Additional effects from the real part of the potential can be taken into account following the lines of [47] by introducing the real part of the pion selfenergy in the pion propagators cut by the horizontal line in Fig. 3 when one evaluates the photon selfenergy with a ph and two pions in the intermediate state. We find a moderate increase of the cross section by an amount of 20 percent in ^{12}C from this effect when the $\pi\pi$ interaction in the medium is considered in addition, but again it does not modify at all the shape of the invariant mass distribution. All these things discussed, we can conclude that while there are uncertainties of about 20 percent in ^{12}C, and a bit more in ^{208}Pb, in the size of the total cross section, the shape of the mass distribution is still determined basically by the implementation of the $\pi\pi$ interaction in the medium.

In Fig. 7 we can see the results for the invariant mass distribution of the two pions for ^{12}C and ^{208}Pb including the absorption of the final pions. The

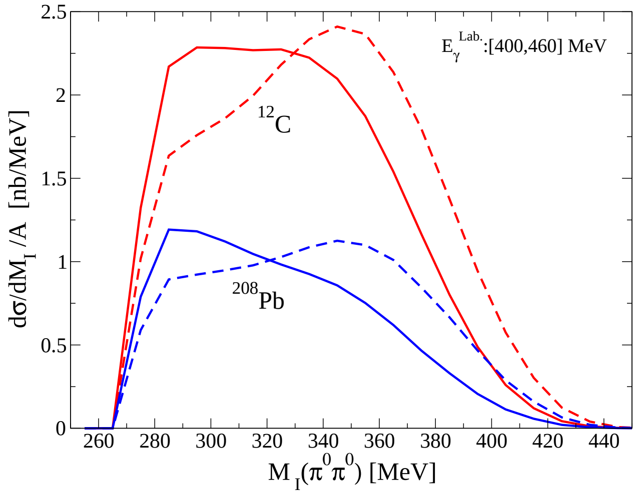

Figure 7: Two pion invariant mass distribution for $2\pi^0$ photoproduction in ^{12}C and ^{208}Pb. Continuous lines: Using the in medium final $\pi\pi$ interaction. Dashed lines: using the final $\pi\pi$ interaction in free space.

difference between the solid and dashed curves is the use of the in medium $\pi\pi$ scattering and G function instead of the free ones, which we take from [17]. As one can see in the figure, there is an appreciable shift of strength to the low invariant mass region due to the in medium $\pi\pi$ interaction. This shift is remarkably similar to the one found in the preliminary measurements of [18].

These results show a clear signature of the modified $\pi\pi$ interaction in the medium. The fact that the photons are not distorted has certainly an advantage over the pion induced reactions and allows one to see inner parts of the nucleus. In this sense it is worth noting that from our calculations we can determine the average nuclear density felt by the reaction which turns out to be 35 per cent and 65 per cent of the normal nuclear density for ^{12}C and ^{208}Pb respectively.

Although we have been discussing the $\pi\pi$ interaction in the nuclear medium it is clear that we can relate it to the modification of the σ in the medium. We have mentioned that the reason for the shift of strength to lower invariant masses in the mass distribution is due to the accumulated strength in the scalar isoscalar $\pi\pi$ amplitude in the medium. Yet, this strength is mostly governed by the presence of the σ pole and there have been works suggesting that the σ should move to smaller masses and widths when embedded in the nucleus [31, 33, 48, 49, 32]. We should stress that in refs. [31, 33, 49, 32] effects of chiral symmetry restoration from the dropping of the condensate and/or the change in the bare σ mass are considered. In our approach which, as we mentioned, relies upon a conventional many body expansion based on standard chiral lagrangians, where the σ in the free space or in the nucleus is generated dynamically, we can look at the σ poles and see their evolution with the nuclear density. A detailed study of the poles in the complex plane for the $\pi\pi$ interaction in the nuclear medium within this approach is now available in [50]. In this work it is indeed found that the pole position of the σ moves to smaller energies as the density increases and the width is also reduced. The present results, when contrasted by the definitive data, if they confirm the preliminary ones of [18], should represent an evidence of this interesting phenomenon which would come to strengthen once more the nature of the σ meson as dynamically generated by the multiple scattering of the pions through the underlying chiral dynamics.

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