Photoproduction of scalar mesons on protons and nuclei

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We study the photoproduction of scalar mesons close to the threshold of $f_0(980)$ and $a_0(980)$ using a unitary chiral model. Peaks for both resonances show up in the invariant mass distributions of pairs of pseudoscalar mesons. A discussion is made on the photoproduction of these resonances in nuclei, which can shed light on their nature, a subject of continuous debate. [S0556-2813(99)01507-1]

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I. INTRODUCTION

The understanding of the scalar sector of mesons has been traditionally very problematic. The low energy scalar states, like the $f_0(980) I^G(J^{PC})=0^+(0^{++})$ and $a_0(980) I^G(J^{PC})=1^-(0^{++})$, have been ascribed to conventional $q\bar{q}$ mesons [1,2], $q^2\bar{q}^2$ states [3,4], $K\bar{K}$ molecules [5,6], glueballs [7], and hybrids [8]. An important step in the understanding of the nature of these states has been made possible in terms of chiral Lagrangians [9] by using a nonperturbative unitary model in coupled channels based on a $O(p^2)$ expansion of the inverse of the K matrix [10], similar to the effective range expansion in quantum mechanics. Within this method, a good reproduction of all data on meson-meson interactions up to 1.2 GeV is obtained, including the scalar and vector resonances, with their position, width and partial decay rates well described.

A further insight into the problem is offered in [11] where an investigation of the meson-meson data up to 1.4 GeV using arguments based on the large N_c limit of QCD is done. This allows one to distinguish between meson resonances which survive in the large N_c limit, which are genuine QCD meson states (essentially built from $q\bar{q}$), and other states which appear from multiple scattering of the mesons and which qualify as quasibound meson-meson states or scattering resonances. The genuine $q\bar{q}$ states in the scalar sector $(L=0)$ would be one octet around 1.4 GeV and a singlet around 1 GeV. The $\sigma(500)$ and $a_0(980)$ appear then as a $\pi\pi$ resonance and a quasibound meson-meson state, respectively. The $f_0(980)$ becomes a mixture of the genuine 1 GeV singlet with large components of a meson-meson quasibound state. This 1 GeV singlet could be associated with the $I=0$ state predicted around this energy in QCD inspired models [12,13] and also has been advocated in phenomenological analyses [14,15]. The effects of the 1 GeV singlet are essentially seen in the $\eta\eta$ decay channels at energies above 1.1 GeV, but it has no practical effect on other channels. This is indicative of the large weight of the meson-meson molecular component in the $f_0(980)$ resonance. For this reason it was possible to obtain a good reproduction of the data of the scalar sector below 1.2 GeV in terms of multiple scattering of the mesons alone with a ''pseudopotential'' provided by the lowest order chiral Lagrangian [16]. In addition, one needs there a cut off to regularize the loop integrals and account for the effect of higher order contributions from the second order Lagrangian. This latter picture is technically very simple and it is the one we shall employ here.

Consistency of these ideas, and in any case information on the nature of these states, can be obtained by producing them on proton targets and also in a nuclear environment which modifies the properties of the building blocks, in this case pseudoscalar mesons, and should have repercussions on the scalar mesons. Hence, we make some suggestions on how the properties of the $f_0(980)$ and $a_0(980)$ scalar resonances in the nuclear medium could be investigated.

II. PHOTOPRODUCTION OF THE SCALAR RESONANCES

The reaction proposed is

$$\gamma p \rightarrow pM, \tag{1}$$

where M is either of the resonances f_0 or a_0. In practice, the meson M will decay into two mesons, $\pi\pi$ or $K\bar{K}$ in the case of the $f_0(980)$ or $K\bar{K}$, $\pi\eta$ in the case of the $a_0(980)$.

The f_0 and a_0 are $L=0$ resonances. Hence we introduce the basic mechanisms that will lead to the s-wave production of the pair of mesons. Photoproduction of pairs of mesons has been the subject of theoretical studies [17–19], particularly for $\pi\pi$ production. The approaches of [17,18] for $\pi\pi$ production are meant for energies of the photon below the K^+K^- production threshold. Without going into detail in these latter models, we can refer to the dominant term, depicted in Fig. 1(a), in order to see that it does not involve the production of the two pions in s-wave. Indeed, the upper vertex corresponds to $\Delta \rightarrow \pi N$ decay, where the pion is produced in a p-wave and the lower vertex is of the type $\vec{S}^+ \cdot \vec{\epsilon}$ leading to an s-wave pion. Another term of relevance in [17,18] is the photoproduction of the $N^*(1520)$ resonance which later decays into a $\Delta\pi$ (with s and d-waves) and the Δ again decays into a πN with the pion in p-wave. Other terms containing explicitly the production of a ρ meson involve directly the pions in p-wave.

One can think of other resonance excitation like the $N^*(1535)$ and its decay into $N\pi\pi$, but this we will show later on that does not lead to 2π in s-wave.

A different point of view can be taken by means of which the production of the s-wave pair of mesons is isolated. This can be accomplished easily in the context of chiral effective

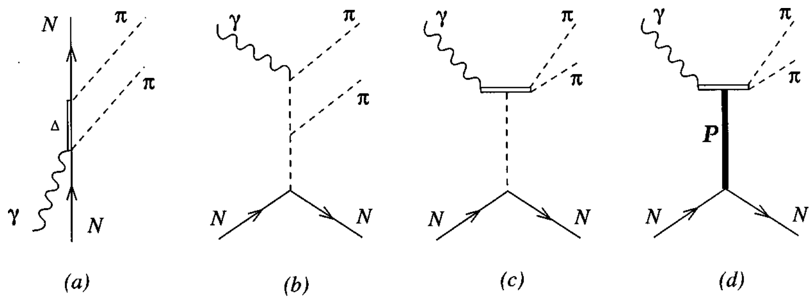

FIG. 1. Mechanisms for two pion production considered in Refs. [17–20].

theories which are meant to work at relatively low energies. This is the approach which we shall follow and for this reason we shall concentrate at energies of the photon close to the threshold production of the $f_0(980)$ and $a_0(980)$ resonances.

In [19] a combined analysis of $\pi\pi$ and $K\bar{K}$ photoproduction in s-wave is conducted. The study is done at higher photon energies than in the present paper, $E_\gamma = 4$ GeV in [19] while here we shall evaluate cross sections for $E_\gamma = 1.7$ GeV. In [19] a particular mechanism for pair production is used, which is depicted in Fig. 1(b). The intermediate meson lines stand for π, ρ, ω. Alternatively a Regge exchange model is used with a strength to be adjusted to data. At the high energies explored and with the model used, there is a strong t dependence of the cross section and the pair of pions are produced in many partial waves, out of which the s-wave is projected out.

In [20] which is concerned about K^+K^- photoproduction, other mechanisms are suggested. They are depicted in Fig. 1(c), where an s-wave resonance is produced from the photon and a virtual meson or a Pomeron [as in diagram 1(d)] is exchanged from a vector meson produced by the photon. In addition, the bremsstrahlung diagrams depicted in Figs. 2(a) and 2(b) are also suggested. It looks clear to us that at high energies the mechanisms of production can be rather complicated, as the complexity of the ϕ photoproduction model of [21], followed by $\phi \to K^+K^-$ decay, shows. Also, as shown in [21], there are many unknown parameters in the theory. The same can be said about the diagrams (c) and (d) of Fig. 1, the strength of which is unknown.

Our approach is different to all of these and relies upon the use of effective chiral Lagrangians. They provide us with Lagrangians for s-wave coupling of pairs of mesons to the baryons, from where the coupling of the external photon becomes straightforward. However, we have the limitation of relatively low energies for the use of these Lagrangians and this is the reason why we concentrate around threshold of the scalar mesons production.

The lowest order Lagrangians for meson-meson and meson-baryon interactions are given by [9]

$$\mathcal{L}_{\rm I} = \mathcal{L}_{\rm I}^{\rm (M)} + \mathcal{L}_{\rm I}^{\rm (MBE)} + \mathcal{L}_{\rm I}^{\rm (MBO)}, \tag{2}$$

with $\mathcal{L}_{\rm I}^{\rm (M)}$ for pure meson-meson interaction, $\mathcal{L}_{\rm I}^{\rm (MBE)}$ for meson-baryon vertices containing an even number of mesons and $\mathcal{L}_{\rm I}^{\rm (MBO)}$ for meson-baryon vertices containing an odd number of mesons. These interaction Lagrangians are given by

$$\mathcal{L}_{\rm I}^{\rm (M)} = \frac{1}{12f^2}\langle (\partial_\mu \Phi \Phi - \Phi \partial_\mu \Phi)^2 + M\Phi^4 \rangle, \tag{3}$$

$$\mathcal{L}_{\rm I}^{\rm (MBE)} = \frac{1}{4f^2}\langle \bar{B} i \gamma^\mu [(\Phi \partial_\mu \Phi - \partial_\mu \Phi \Phi) B - B(\Phi \partial_\mu \Phi - \partial_\mu \Phi \Phi)] \rangle, \tag{4}$$

$$\mathcal{L}_{\rm I}^{\rm (MBO)} = \frac{D+F}{2}\langle \bar{B} \gamma^\mu \gamma^5 u_\mu B \rangle + \frac{\mathcal{D}-F}{2}\langle \bar{B} \gamma^\mu \gamma^5 B u_\mu \rangle, \tag{5}$$

with u_μ up to three meson fields given by

$$u_\mu = -\frac{\sqrt{2}}{f}\partial_\mu \Phi + \frac{\sqrt{2}}{12f^3}(\partial_\mu \Phi \Phi^2 - 2\Phi \partial_\mu \Phi \Phi + \Phi^2 \partial_\mu \Phi), \tag{6}$$

with f the pion decay constant, the symbol $\langle\,\rangle$ standing for the trace of the SU(3) matrices and Φ, B the meson and baryon SU(3) matrices given by

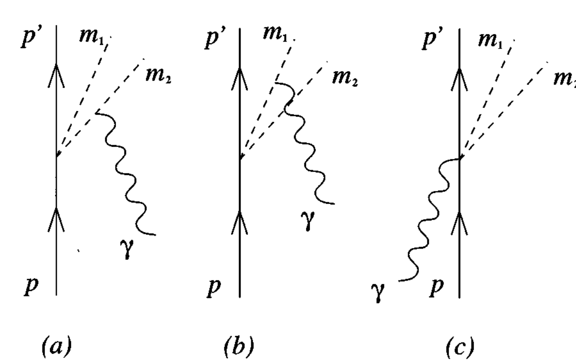

FIG. 2. Feynman diagrams for the process $\gamma p \to p m_1 m_2$: (a) and (b), meson pole processes; (c) contact term required by gauge invariance.

$$\Phi=\begin{pmatrix} \frac{1}{\sqrt{2}}\pi^0+\frac{1}{\sqrt{6}}\eta & \pi^+ & K^+ \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^0+\frac{1}{\sqrt{6}}\eta & K^0 \\ K^- & \bar{K}^0 & -\frac{2}{\sqrt{6}}\eta \end{pmatrix}, \tag{7}$$

$$B=\begin{pmatrix} \frac{1}{\sqrt{2}}\Sigma^0+\frac{1}{\sqrt{6}}\Lambda & \Sigma^+ & p \\ \Sigma^- & -\frac{1}{\sqrt{2}}\Sigma^0+\frac{1}{\sqrt{6}}\Lambda & n \\ \Xi^- & \Xi^0 & -\frac{2}{\sqrt{6}}\Lambda \end{pmatrix}. \tag{8}$$

Our starting point will be the meson-baryon→meson-baryon vertex originated from the Lagrangian of Eq. (4). We need actually only the $K^-p\to K^-p$ and $\pi^-p\to\pi^-p$ couplings which are given by

$$V_{\pi(K)}=-C_{\pi(K)}\frac{1}{4f^2}\overline{u}(p')\gamma^\mu u(p)(k_\mu+k'_\mu), \tag{9}$$

where k,k' are the momenta of the incoming and outgoing mesons and $C_\pi=1$, $C_K=2$. As mentioned above, the constant f is the pion decay constant ($f_\pi=93\,\mathrm{MeV}$). In the present work the $K\bar{K}$ states are the most relevant building blocks of the resonances, and thus the $K^-p\to K^-p$ is the relevant ingredient. In Refs. [22, 24] a study of the latter reaction and coupled channels was done using a unitary chiral method. We follow here the approach of [24] where an intermediate value of f between the one of kaons and pions $f=1.15f_\pi$ was chosen and this will be used here too for this latter amplitude. The choice of an average value for f in [24], together with the choice of a cut off, is an approximate way to incorporate the effects of higher order Lagrangians, which is possible in the K^-p sector but not in other meson baryon channels [22,23].

For the low energies which we will consider here and only s-wave of the pair of mesons, the vertex of Eq. (9) simplifies to

$$V_{\pi(K)}=-C_{\pi(K)}\frac{1}{4f^2}(k^0+k'^0). \tag{10}$$

This vertex, together with the standard electromagnetic coupling of the photon to the mesons, allows one to evaluate diagrams (a) and (b) of Fig. 2. However, gauge invariance requires the presence of the contact term of Fig. 2(c), which we also need, for $\gamma\pi^-p\to\pi^-p$ and $\gamma K^-p\to K^-p$, or analogously $\gamma p\to\pi^+\pi^-p$, $\gamma p\to K^+K^-p$. The vertex is given by

$$V^\gamma_{\pi(K)}=-C_{\pi(K)}\frac{e}{2f^2}\overline{u}(p')\gamma^\mu u(p)\epsilon_\mu. \tag{11}$$

As argued above we choose an energy of the photon around $E_{\gamma\mathrm{lab}}=1.7\,\mathrm{GeV}$. This allows one to produce the scalar resonances close to threshold. This kinematics allows us to simplify Eq. (11) which becomes now in the c.m. of the γp system and using the Coulomb gauge, $\epsilon^0=0$, $\vec{\epsilon}\cdot\vec{k}=0$,

$$V^\gamma_{\pi(K)}=C_{\pi(K)}\frac{e}{2f^2}\frac{i(\vec{\sigma}\times\vec{q})\vec{\epsilon}}{2M}, \tag{12}$$

with M the mass of the proton and $\vec{q}$ the photon momentum.

For the case of K^+K^- production, with the energy of the photon chosen, the kaon momenta are very small. In this case, the kaon bremsstrahlung diagrams, 2(a) and 2(b), give a negligible contribution (less than 5%) and we shall neglect them. This is not the case for the pions, which carry a larger momentum and these mechanisms become important. On the other hand, there are many mechanisms for $\pi^+\pi^-$ production around this region, as can be seen by the relative complexity of the models used to study the process $\gamma p\to\pi^+\pi^-p$ up to $E_\gamma\simeq 1$ GeV in [17,18].

In the case of $\pi\pi$ production we shall evaluate the contribution from the f_0 resonance and we will estimate the background from the experimental cross section. This should give us an idea of the ratio of the signal for f_0 excitation to the background to be found in actual experiments. On the other hand, the near threshold cross sections for K^+K^- production evaluated here should be rather realistic, since other terms which can be constructed in analogy to the model of [17] would vanish at threshold.

The next step necessary to build up the scalar meson resonances is the final state interaction of the mesons. For this purpose we follow the approach of [16], where the resonances are obtained through an iteration of the lowest order chiral Lagrangian vertex considered as a potential in the Bethe-Salpeter equation. This is depicted in Figs. 3(a), 3(b), and 3(c).

However, unlike the tree level bremsstrahlung diagrams of Figs. 2(a) and 2(b), which are either negligible at threshold of the meson pair production, or have a strong angular dependence when the meson momenta are not small, the loops considered in Figs. 2(d)–2(g) directly contribute to s-wave pair production and are also required by gauge invariance. The set of diagrams in Fig. 3 build up the s-wave resonance production and are evaluated below.

We have mentioned above how the main terms in $\pi\pi$ production in [17,18] do not produce the two pions in s-wave. One can envisage other mechanisms for the s-wave resonance production like the one corresponding to diagrams 3(a), 3(b), 3(c), where the photon couples to the nucleon before of after the *NNMM* vertex. In this case the dominant component would vanish at threshold of resonance production since it involves the amplitude of Eq. (10) but with $(k^0-k'^0)$ rather than $(k^0+k'^0)$. Smaller components from $\vec{\gamma}\cdot(\vec{k}-\vec{k}')$ from Eq. (9) would be even further suppressed, since at threshold of resonance production there is a cancel-

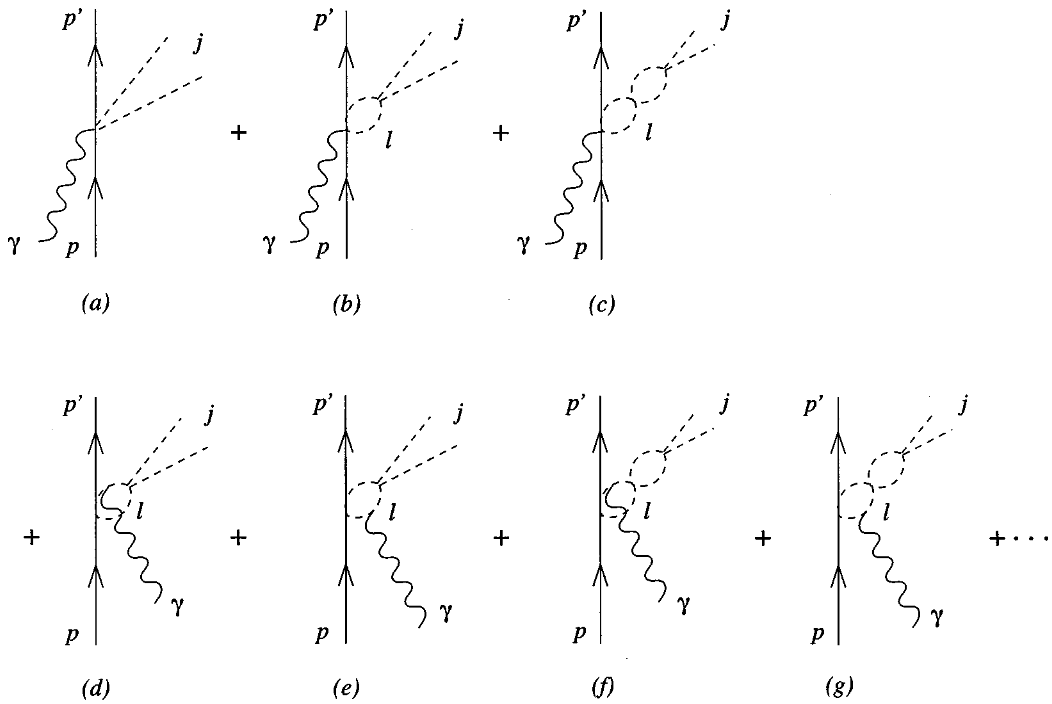

FIG. 3. Iterated terms from the contact term and from mesonic bremsstrahlung.

lation between the diagrams where the photon couples before and after the *NNMM* vertex. It is also easy to see that for parity reasons, terms like those in Figs. 3(f) and 3(g) with the photon coupling to the second loop do vanish.

Unitarity in coupled channels for the two strongly interacting mesons is one of the important ingredients here in order to produce the $f_0(980)$ and $a_0(980)$ resonances. One could also think of the coupling of the final *BMM* system to intermediate *BM* states. In this case one should select the case with *BM* in *s*-wave. The coupling of γN to *BM* in *s*-wave has been worked out in [22] and diagrammatically it is depicted in Fig. 4(a), incorporating the *BM* final state interaction.

Within the set of chiral Lagrangians written in Eqs. (3)–(5) the way to couple an extra meson on top of the intermediate *MB* state is through diagrams like those depicted in Figs. 4(b)–4(d). All these terms require the use of the $\mathcal{L}_{\rm I}^{\rm (MBO)}$ Lagrangian of Eq. (5), where the $\gamma^\mu\gamma_5\partial_\mu\Phi$ combination leads to a $\vec{\sigma}\vec{p}$ vertex in the nonreactivistic reduction, and hence a *p*-wave meson, leading to terms which do not contribute to the scalar meson production.

In the particular case of the $N^*(1535)$ excitation, the arguments above can be expressed by stating that the N^* with negative parity cannot decay into a nucleon and two pions in *s*-wave.

If one goes beyond the chiral approach and considers mechanisms for two pion production involving the excitation of resonances, we find two types of diagrams which would provide contribution. One of them is the $\gamma N \to N^*$ process followed by $N^* \to N\pi\pi$ ($I=0$, *s*-wave). The N^* should be a $1/2^+$ state in this case, and restricting ourselves to resonances below $\sqrt{s}=2000$ MeV, we find the $N^*(1440)$ and the $N^*(1710)$. The mechanism mentioned was considered for the case of the $N^*(1440)$ in [17] and found to be relevant only at threshold, but negligible at higher energies. Here we consider photons at higher energies than in [17] and the $N^*(1710)$ has more chances to be relevant. Unfortunately both the uncertainty in the width, $\Gamma = 50-250$ MeV, and the branching ratio for decay into $N\pi\pi$ ($I=0$, *s*-wave), $B=10-40\,\%$, introduce large uncertainties in this contribution. On the other hand, the helicity amplitude for this resonance has also large uncertainties but seems to be reasonably smaller than in the $N^*(1440)$ case [25].

The second mechanism would correspond to diagrams like in Fig. 1(a) with an N^* instead of a Δ in the intermediate baryon. In this case we should have a $1/2^-$ state to allow for an *s*-wave pion in the $N^* \to N\pi$ decay. Here we would have the $N^*(1535)$ and $N^*(1650)$ resonances. The uncertainties here stem from the contact vertex gamma $NN^*\pi$, which from minimal coupling from the leading constant $NN^*\pi$ vertex would be zero, and in practice should be relatively small.

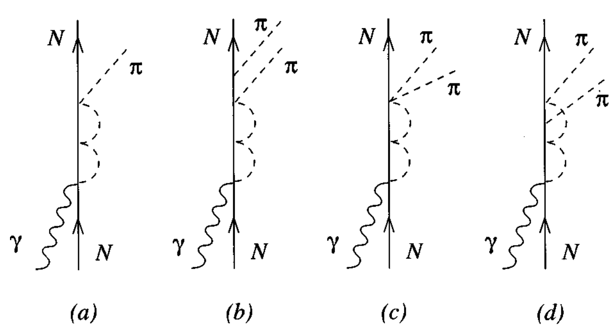

FIG. 4. Diagrams for one and two pion production including coupled channels.

The presence of such mechanisms would introduce some elements of uncertainty in the cross sections evaluated here, although from the arguments used above these terms do not seem to be large. There is still another argument that would favor the mechanism chosen here. Indeed, all the resonant terms discussed above would provide the maximum contribution when their propagators are placed on shell in the diagrams, providing an imaginary part of the amplitude. As we shall see later on, the mechanisms considered here lead to a peak in the real part of the amplitude when the $f_0(980)$ resonance is excited which can interfere with the largely real amplitude of the whole $\gamma N \rightarrow \pi\pi N$ process, hence magnifying the effect of the resonance. This would not be the case of the N^* resonance excitation mechanisms which would contribute mostly an imaginary part to the amplitude in the case when the resonance is placed on shell and the mechanism is most relevant.

Accepting some uncertainties, the arguments given above indicate that the mechanism considered here can provide a fair estimate of the strength of the scalar resonance excitation, which is sufficient for the purpose of the present paper, where an exploration of the possibilities of observation of these resonances in gamma induced reactions is made.

Coming back to our model, the final states with a pair of mesons which can be produced in the reactions are $\pi^+\pi^-, \pi^0\pi^0, K^+K^-, K^0\overline{K^0}, \pi^0\eta$. Note that even if the $\pi^0\pi^0$, $K^0\overline{K^0}$ and $\pi^0\eta$ do not couple to the photon vertex in diagram (c) of Fig. 3, they can appear in the final states through the iterated terms of diagrams (b), (c)–(g) when we sum over the intermediate state, l, which can be either $\pi^+\pi^-$ or K^+K^-.

The sum of diagrams in Fig. 3 bears a close resemblance to the ϕ decay into $K\overline{K}\gamma$ which has been studied in [26–28]. Indeed, the vertex of Eq. (9) changes $k_\mu + k'_\mu$ by $k_\mu - k'_\mu$ when the two mesons are outgoing and has the same structure as the $\phi - K\overline{K}$ vertex which goes as $\epsilon^\mu(\phi)(k_\mu - k'_\mu)$. The results of [26–28] are very useful. Using arguments of gauge invariance it is found there that the sum of the loops in Figs. 3(b), 3(d), and 3(e) is convergent and replaces the two meson loop of Fig. 3(b)

$$G(P) = i\int \frac{d^4k}{(2\pi)^4}\frac{1}{k^2 - m^2 + i\epsilon}\frac{1}{(P-k)^2 - m^2 + i\epsilon} \quad (13)$$

by

$$\tilde{G}(Q,P) = -\frac{Q\cdot q}{(2\pi)^2}\int_0^1 dx \int_0^x \times dy \frac{(1-x)y}{Q^2 x(1-x) - 2Q\cdot q(1-x)y - m^2 + i\epsilon}, \quad (14)$$

where P is the four momentum of the two mesons, m the mass of the meson in the loop (π^+ or K^+ in our case) and $Q = p - p'$. In addition, M_I is the invariant mass of the pair of mesons and the invariant product $Q\cdot q$ is given here by

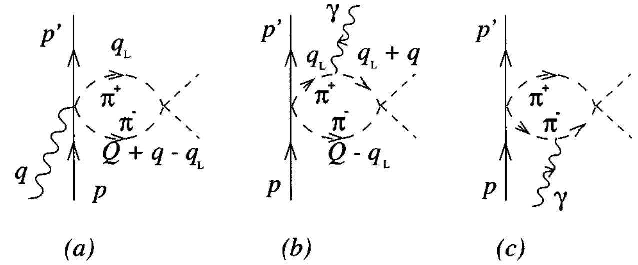

FIG. 5. One loop diagrams for the process $\gamma p \rightarrow pMM$.

$$Q\cdot q = \frac{1}{2}(Q^2 - M_I^2),$$

$$Q^2 = 2M^2 - 2E(\vec{p})E(\vec{p}\,') + 2|\vec{p}||\vec{p}\,'|\cos\theta. \quad (15)$$

This introduces a dependence of the t matrix for the process in the angle of $\vec{p},\vec{p}\,'$, but not the angle of the mesons with the photon.

The evaluation proceeds as follows. We must evaluate the contribution of the three loops which we show now in Fig. 5 with the appropriate labels for the momenta of the particles.

Diagrams (b) and (c) contribute on equal amount. We shall call contact term (C) the one coming from diagram (a) and bremsstrahlung (B) the one coming from diagrams (b) and (c). The contribution goes as (for the case of intermediate $\pi^+\pi^-$)

$$-it^{(\gamma C)} = -C_\pi \frac{e}{2f^2}\overline{u}(\vec{p}\,')\gamma^\mu u(\vec{p})\epsilon_\mu \times \int \frac{d^4q_L}{(2\pi)^4}\frac{1}{q_L^2 - \mu^2 + i\epsilon} \times \frac{1}{(Q+q-q_L)^2 - \mu^2 + i\epsilon}t^{(S)}, \quad (16)$$

$$-it^{(\gamma B)} = C_\pi \frac{e}{2f^2}\overline{u}(\vec{p}\,')\gamma^\mu u(\vec{p}) \times \int \frac{d^4q_L}{(2\pi)^4}2eq_L^\nu\epsilon_\nu(2q_L - Q)_\mu \frac{1}{q_L^2 - \mu^2 + i\epsilon} \times \frac{1}{(Q-q_L)^2 - \mu^2 + i\epsilon}\frac{1}{(q_L+q)^2 - \mu^2 + i\epsilon}t^{(S)}, \quad (17)$$

where $t^{(S)}$ is the strong meson-meson amplitude. Since Q,q are the only vectors not integrated in Eqs. (16) and (17), the sum of the two terms has a structure of the type

$$\gamma^\mu\epsilon^\nu\{ag_{\mu\nu} + bQ_\mu Q_\nu + cQ_\mu q_\nu + dQ_\nu q_\mu + eq_\mu q_\nu\}, \quad (18)$$

where the contact term only contributes to $ag_{\mu\nu}$ while the B term contributes to all.

Gauge invariance of the sum of all terms requires that the expression of Eq. (18) vanishes with the substitution $\epsilon^\nu \rightarrow q^\nu$. This implies

TABLE I. Elements of the transition t matrix from the state l to j ($T_{lj}=T_{jl}$). The isospin $I=0$, $I=1$ matrix elements can be found in [16].

	K^+K^-	$K^0\bar{K}^0$	$\pi^+\pi^-$	$\pi^0\pi^0$	$\pi^0\eta$
K^+K^-	$\frac{1}{2}\{t^{I=0}_{K\bar{K},K\bar{K}}+t^{I=1}_{K\bar{K},K\bar{K}}\}$	$\frac{1}{2}\{t^{I=0}_{K\bar{K},K\bar{K}}-t^{I=1}_{K\bar{K},K\bar{K}}\}$	$\sqrt{\frac{1}{3}}\,t^{I=0}_{K\bar{K},\pi\pi}$	$\sqrt{\frac{1}{3}}\,t^{I=0}_{K\bar{K},\pi\pi}$	$-\sqrt{\frac{1}{2}}\,t^{I=1}_{K\bar{K},\pi^0\eta}$
$K^0\bar{K}^0$		$\frac{1}{2}\{t^{I=0}_{K\bar{K},K\bar{K}}+t^{I=1}_{K\bar{K},K\bar{K}}\}$	$\sqrt{\frac{1}{3}}\,t^{I=0}_{K\bar{K},\pi\pi}$	$\sqrt{\frac{1}{3}}\,t^{I=0}_{K\bar{K},\pi\pi}$	$\sqrt{\frac{1}{2}}\,t^{I=1}_{K\bar{K},\pi^0\eta}$
$\pi^+\pi^-$			$\frac{2}{3}\,t^{I=0}_{\pi\pi,\pi\pi}$	$\frac{2}{3}\,t^{I=0}_{\pi\pi,\pi\pi}$	0
$\pi^0\pi^0$				$\frac{2}{3}\,t^{I=0}_{\pi\pi,\pi\pi}$	0
$\pi^0\eta$					$t^{I=1}_{\pi^0\eta,\pi^0\eta}$

$$aq^\mu+bQ^\mu(Q\cdot q)+dq^\mu(Q\cdot q)=0 \tag{19}$$

or equivalently

$$b=0;\quad a=-d(Q\cdot q). \tag{20}$$

On the other hand, in the Coulomb gauge chosen where $\epsilon^0=0$, $\vec{\epsilon}\vec{q}=0$ in the γp c.m. frame, the expression of Eq. (18) is greatly simplified since we have that

$$\epsilon^\mu Q_\mu=-\vec{\epsilon}\vec{Q}=-\vec{\epsilon}(\vec{p}-\vec{p}\,')\simeq-\vec{\epsilon}\vec{p}=\vec{\epsilon}\vec{q}=0, \tag{21}$$

$$\epsilon^\mu q_\mu=-\vec{\epsilon}\vec{q}=0, \tag{22}$$

where we have assumed $\vec{p}\,'\simeq 0$ because we work close to the scalar meson photoproduction threshold. Hence only the $g_{\mu\nu}$ term of Eq. (18) contributes to the amplitude. The trick then is to evaluate the term which goes like $dQ_\nu q_\mu$ to which only the B diagrams contribute and from it via Eq. (20) obtain the coefficient a which is the only one needed to evaluate the amplitudes. Equation (17) is then evaluated using the Feynman technique and the terms proportional to $Q_\nu q_\mu$ are kept. For dimensional reasons the rest of the integral has two powers less in the loop variable q_L, and hence the contribution to this term is convergent and via Eq. (20) leads to the result of Eq. (14). Further details on the integration technique can be seen in [26–28].

The work of [28] adds some new relevant ingredients to the work of [26,27], since it proves that by using chiral Lagrangians to deal with the final state interaction of the mesons after the first loop which involves the photon, the strong t matrix for meson-meson interaction factorizes on shell. This occurs because the $MM\to MM$ amplitudes in lowest order have the structure $\alpha s+\beta\Sigma_i p_i^2$, which can be recast into $\alpha s+\beta\Sigma_i m_i^2+\beta\Sigma_i(p_i^2-m_i^2)$, where the first two terms account for the on-shell part. The last term in this former expression kills one of the meson propagators in Eqs. (16) and (17) and does not provide contribution to the $Q_\nu q_\mu$ term.

With all these ingredients we can write the sum of the diagrams in Fig. 3 which leads to the amplitude

$$t_j^\gamma=\frac{e}{4f^2}\,\frac{i(\vec{\sigma}\times\vec{q})\vec{\epsilon}}{2M}\left(D_j+\sum_l D_l\tilde{G}_l T_{lj}\right), \tag{23}$$

where D_j is the vector (4,0,2,0,0) counting the channels in the following order, $K^+K^-, K^0\bar{K}^0, \pi^+\pi^-, \pi^0\pi^0, \pi^0\eta$. The matrix T_{lj} in Eq. (23) is the transition t matrix from the meson state l to j. These matrix elements are easily obtained from [16] using the isospin decomposition of the states and we find the matrix of Table I, in terms of the isospin $I=0$, $I=1$ matrix elements derived in [16].

One should note that the matrix elements involving pions use a unitary normalization in [16] including an extra factor $1/\sqrt{2}$ per each pair of pion states. This normalization is convenient to account for factors due to the identity of the particles when summing over intermediate states. The amplitudes of Table I are the physical ones, where the proper normalization of the states is used.

The resonance structure of the pair of mesons comes from the term $\Sigma_l D_l\tilde{G}_l T_{lj}$ in Eq. (23). Hence, for the case of pion pair production, we remove the isolated term D_j in Eq. (23) which, together with other terms will build up the background for this process. In the case of $K\bar{K}$ production the threshold is above the f_0 and a_0 mass and the cross section does not exhibit the resonance structure, although the amplitudes are affected by it. In this case we keep all terms since with the amplitude of Eq. (23) we are producing absolute cross sections.

The function $\tilde{G}$ of Eq. (14) can be written in an analytical form following [26,27]. However, there are novel ingredients here since Q^2 can be negative, unlike the case of the ϕ decay, where m_ϕ^2 is positive. For this reason we give below the analytical expressions valid in all the range of values of Q^2, M_I^2:

$$\tilde{G}(Q^2,M_I^2)=\frac{1}{8\pi^2}\left\{\frac{1}{2}-\frac{2}{a-b}\left[f\left(\frac{1}{b}\right)-f\left(\frac{1}{a}\right)\right]+\frac{a}{a-b}\left[g\left(\frac{1}{b}\right)-g\left(\frac{1}{a}\right)\right]\right\}, \tag{24}$$

where $a=Q^2/m^2$, $b=M_I^2/m^2$, m is the mass of the meson in the loop and $f(x)$ and $g(x)$ are given by

$$f(x)=\begin{cases} -\left[\arcsin\left(\frac{1}{2\sqrt{x}}\right)\right]^2 & \text{for } x>\frac{1}{4} \\ \frac{1}{4}[\ln(\eta+/\eta_-)-i\pi]^2 & \text{for } 0<x<\frac{1}{4} \\ \left[\text{argsinh}\left(\frac{1}{2\sqrt{-x}}\right)\right]^2 & \text{for } x<0, \end{cases}$$

$$g(x)=\begin{cases} (4x-1)^{1/2}\arcsin\left(\frac{1}{2\sqrt{x}}\right) & \text{for } x>\frac{1}{4} \\ \frac{1}{2}(1-4x)^{1/2}[\ln(\eta_+/\eta_-)-i\pi] & \text{for } 0<x<\frac{1}{4} \\ (1-4x)^{1/2}\text{argsinh}\left(\frac{1}{2\sqrt{-x}}\right) & \text{for } x<0, \end{cases}$$

$$\eta_\pm=\frac{1}{2x}[1\pm(1-4x)^{1/2}]. \tag{25}$$

The particular structure of Eq. (23) allows one to obtain an easy formula for the invariant mass distribution of the two mesons

$$\left.\frac{d\sigma}{dM_I}\right|_j=\frac{1}{(2\pi)^3}\frac{1}{4s}\frac{M^2}{s-M^2}\frac{1}{M_I} \times S\lambda^{1/2}(s,M_I^2,M^2)\lambda^{1/2}(M_I^2,m_1^2,m_2^2) \times\frac{1}{2}\int_{-1}^{1}d\cos\theta\bar{\sum}\sum|t_j^\gamma|^2, \tag{26}$$

where m_1,m_2 are the masses of the two mesons in the final meson-meson state, λ is the ordinary Källen function and S is a symmetry factor, 1/2 for $\pi^0\pi^0$ in the final state and 1 for the other channels.

The technique used here has been recently applied to the study of the radiative ϕ decay, $\phi\rightarrow\pi^+\pi^-\gamma$, which proceeds via K^+K^- loops [30]. An invariant mass distribution is predicted in [30] with a clear peak for the $f_0(980)$ excitation. Recently the measurements have been concluded at Novosibirsk [31] and the experimental distribution is in perfect agreement with the predictions of [30]. This finding gives us extra confidence in the techniques used here for the photoproduction processes.

III. RESULTS

In Fig. 6 we show the results for the 5 channels considered. We observe clear peaks for $\pi^+\pi^-$, $\pi^0\pi^0$ and $\pi^0\eta$ production around 980 MeV. The peaks in $\pi^+\pi^-$ and $\pi^0\pi^0$ clearly correspond to the formation of the $f_0(980)$ resonance, while the one in $\pi^0\eta$ corresponds to the formation of the $a_0(980)$. The $\pi^0\pi^0$ cross section is 1/2 of the $\pi^+\pi^-$ one due to the symmetry factor S in Eq. (26). As commented

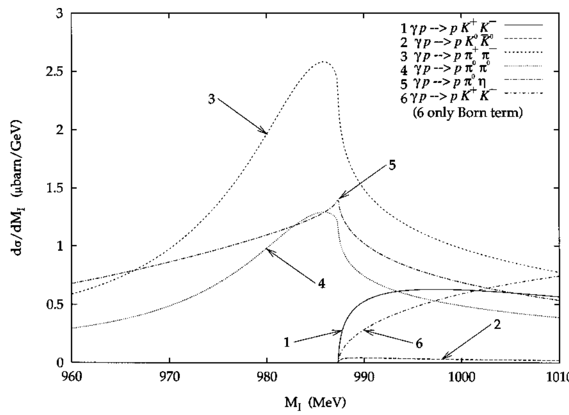

FIG. 6. Results for the cross section on protons as a function of the invariant mass of the meson-meson system.

above, the K^+K^- and $K^0\overline{K^0}$ production cross section appears at energies higher than that of the resonances and hence do not show the resonance structure. Yet, final state interaction is very important and increases appreciably the K^+K^- production cross section for values close to threshold with respect to the Born approximation [only the D_j term in Eq. (23), or diagram (a) of Fig. 3].

It is also interesting to see the shapes of the resonances which differ appreciably from a Breit-Wigner, due to the opening of the $K\overline{K}$ channel just above the resonance [29].

We would like to stress here that the invariant mass distributions for resonance excitation into the various pseudoscalar channels depicted in Fig. 6 are theoretical predictions of a chiral unitary model, in this case the one of [16], where only one parameter was fitted to reproduce all the data of the meson-meson interaction in the scalar sector.

A small variant of this reaction would be the $\gamma p\rightarrow nM\overline{M}$. In this case the $M\overline{M}$ system has charge +1 and hence $I=0$ is excluded, hence, one isolates the a_0 production.

It is interesting to notice the origin of the peak structure for $\pi\pi$ production. Indeed, the cross section for $\pi\pi\rightarrow\pi\pi$ in $I=0$ exhibits a minimum at the f_0 energy because of the interference between the f_0 contribution and the $\sigma(500)$

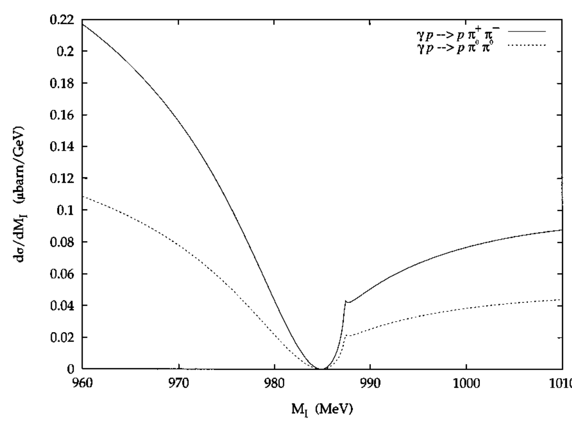

FIG. 7. Results for the cross section on protons as a function of the invariant mass of the meson-meson system including only the $\pi^+\pi^-$ contribution on the first loop.

broad resonance. We can see this here also by killing the K^+K^- first loop in the diagrams of Fig. 3. The results are shown in Fig. 7 for $\pi^+\pi^-$ and $\pi^0\pi^0$ production. Very small cross sections and a clear minimum around the f_0 position can be seen in the figure. This means that the resonant structure for $\pi\pi$ production of Fig. 6 is due to the K^+K^- first loop which factorizes the $K^+K^-\to\pi\pi$ amplitude in the final state interaction. This amplitude has a peak at the f_0 position but cannot be seen in the $K^+K^-\to\pi\pi$ cross section because the resonance is below threshold. A reaction like the present one which factorizes this amplitude at energies below the physical threshold can then evidence the peak, as is indeed visible in the figure.

However, we should bear in mind that we have plotted there the contribution of the f_0 resonance alone. The three level contact term and bremsstrahlung diagrams, plus other contributions which would produce a background, are not considered there. We estimate the background from the experimental cross section for $\gamma p\to p\pi^+\pi^-$ of [32], which is around 45 μb at $E_\gamma=1.7\,\mathrm{GeV}$. Using Eq. (26), assuming $\overline{\Sigma}\Sigma|t_j^\gamma|^2$ constant and integrating over the range of M_I allowed, we determine that constant from the experimental cross section and then the same Eq. (26) gives us the background for $d\sigma/dM_I$. This provides a background of around 55 μb/GeV while the resonant peak has about 2.5 μb/GeV strength. This gives a ratio of 5% signal to background assuming that the background is mostly real versus an imaginary contribution from the resonance and hence there would be no interference. The situation with the $\pi^0\pi^0$ channel should be better because the $\gamma p\to\pi^0\pi^0 p$ cross section is about eight times smaller than the one for $\gamma p\to\pi^+\pi^- p$ [33,34]. Considering that the resonant signal now is a factor two smaller than the $\gamma p\to\pi^+\pi^- p$ cross section, this would give a ratio of signal to background of 20%, which should be more clearly visible in the experiment. The same or even better ratios than in the $\pi^0\pi^0$ case are expected for $\pi^0\eta$ production in the a_0 channel, since estimates of the background along the lines of present models for $\pi^0\pi^0$ production [17,18] would provide a cross section smaller than for $\pi^0\pi^0$ production.

However, there is a distinct feature about the f_0 resonance which makes its contribution, in principle, bigger than the estimates given above. Indeed, the f_0 is approximately a Breit-Wigner resonance with an extra phase of $e^{i\pi/2}$. This means that the real part has a peak while the imaginary part changes sign around the resonance energy. This was the case in the $K\bar{K}\to\pi\pi$ amplitude [35] and we have the same situation here as can be seen in Fig. 8. This means that assuming the background basically real, there would be an interference with the f_0 resonance which would lead to an increase of about 50% over the background, or a decrease by about 40% (depending on the relative sign) for the $\pi^+\pi^-$ case and larger effects for the $\pi^0\pi^0$ case. This is of course assuming weak dependence on momenta and spin of the background amplitudes. In any case, due to the particular feature of the f_0 resonance discussed above, it is quite reasonable to expect bigger signals than the estimates based on a pure incoherent sum of cross sections.

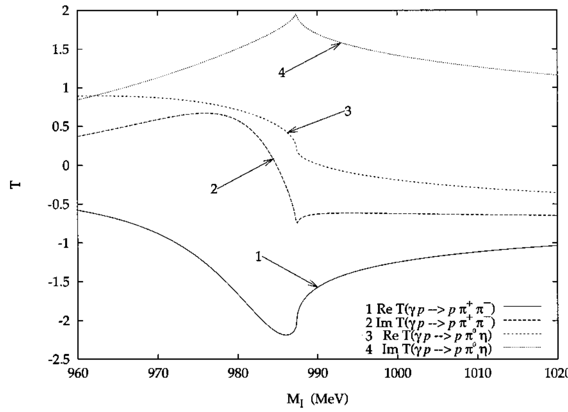

FIG. 8. Real and imaginary parts of the resonant piece of the amplitude t_j^γ $[\Sigma_l D_l \tilde{G}_l T_{lj}$ of Eq. (23)], for the channels $\pi^+\pi^-$ and $\pi^0\eta$.

Certainly it is possible to obtain better ratios if one looks at angular correlations. If one looks in a frame where the two mesons are in their c.m., the bremsstrahlung pieces (both from the squared of the bremsstrahlung term as well as from interference with s-wave terms) have a $\sin^2\theta$ dependence, with θ the angle between the meson and the photon. Other terms from [17,18] exhibit equally strong angular dependence, for what extraction of the angle independent part of the cross section would be an interesting exercise which would select the part of the cross section to which the resonant contribution obtained here belongs to.

IV. MESON RESONANCE PRODUCTION IN NUCLEI

Now we turn our attention to nuclei. As mentioned in the introduction there is much debate about the nature of the scalar meson resonances. The modification of the properties of these resonances in nuclei should depend on their nature. For instance, it would not be modified in the same way if it is a $q\bar{q}$ state than if it is a $K\bar{K}$ molecule. Also, our scheme does not rely upon any of these pictures, although it gives some support to the quasimolecular nature of the states. In any case, we saw that the loop structure, with mostly $K\bar{K}$ in the loop, is what leads to the f_0 production. Thus, the production in nuclei would be modified due to the $K,\bar{K}$ modification in a nuclear medium, but in a particular way, due to the modification of the $\tilde{G}$ function in a medium when the $K,\bar{K}$ propagators are substituted by their renormalized ones in the medium. The changes expected would certainly differ form those expected on the base of the assumption of a $K\bar{K}$ molecule and particularly a $q\bar{q}$ state for this resonance. In this sense modifications of the resonance properties in nuclei are bound to offer us some information on the nature of the states, eventually reinforcing the chiral unitary approach interpretation of those states. The evaluation of the nuclear modifications would require the use of $K,\bar{K}$ self-energies in the nuclear medium, or equivalently their optical potentials.

The interaction of $K,\bar{K}$ with nuclei is a subject that has attracted much attention [36]. Interesting developments have been done recently looking at K^-N scattering from a chiral

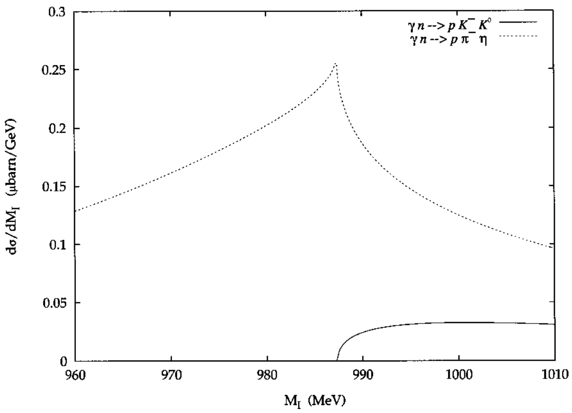

FIG. 9. Results for the cross section on neutrons as a function of the invariant mass of the meson-meson system.

perspective, which have allowed to tackle the problem of the $K,\bar{K}$ nucleus interaction with some novel results [37–39]. The issue is not yet settled since there are still important discrepancies between the different results. It looks wise to allow some time for the issue to get clarified before one tackles the problem suggested here, which looks certainly quite interesting. Meanwhile we can make some exploration following the lines of the former sections. The first thing which we observe is that if one looks for a proton in the final state, one can have the $\gamma n \to p \pi^- \eta (K^- K^0)$ and approximately one would expect a cross section

$$\left.\frac{d\sigma}{dM_I}\right|_A \simeq Z \frac{d\sigma}{dM_I}(p) + N \frac{d\sigma}{dM_I}(n). \tag{27}$$

The latter cross section can proceed through the meson channels $K^- K^0$ and $\pi^- \eta$, both in $I=1$. The cross sections in this case, for these two channels and in this order, are given again by means of Eqs. (23) and (26), by taking in Eq. (23) the vector $D=(1,0)$ and the matrix T_{lj} as

$$T_{lj} = \begin{pmatrix} t^{I=1}_{K\bar{K},K\bar{K}} & t^{I=1}_{K\bar{K},\pi\eta} \\ t^{I=1}_{K\bar{K},\pi\eta} & t^{I=1}_{\pi\eta,\pi\eta} \end{pmatrix}. \tag{28}$$

These cross sections are about one order of magnitude smaller than those on the proton target, as seen in Fig. 9.

Hence, in nuclei we should expect a cross section roughly Z times the one of the proton, unless the properties of the resonances a_0 and f_0 are drastically modified in the medium, which is, however, what one expects. One should recall that the relatively small widths of the f_0 and a_0 resonances are due to the small coupling to the $\pi\pi$ and $\pi\eta$ channels respectively. The resonances, however, couple very strongly to the $K\bar{K}$ system but the decay is largely inhibited because the $K\bar{K}$ threshold is above the resonance mass. Only the fact that the resonances have already a width for $\pi\pi$ and $\pi\eta$ decay, respectively, allows the $K\bar{K}$ decay through the tail of the resonance distribution. If the K^- develops a large width on its own this enlarges considerably the phase space for $K\bar{K}$ decay and the a_0, f_0 width should become considerably larger.

Given the interest that the modifications of meson resonances in nuclei, like the σ [40,41], ρ [42,43], etc., is raising, the study of the modifications of the f_0 and a_0 is bound to offer us some insight into the nature of these resonances, that has been so much debated, and eventually into the chiral approach to these resonances which we have discussed in this paper.

V. CONCLUSIONS

In summary, we have studied the photoproduction of the $f_0(980)$ and $a_0(980)$ resonances for photon energies close to the $K\bar{K}$ production threshold using tools of chiral unitary theory. The $K\bar{K}$ production cross sections were evaluated and the effect of the resonances was shown to modify drastically the cross sections with respect to the Born approximation. The f_0 and a_0 resonances led to peaks in the invariant mass distributions of $\pi\pi$ and $\pi\eta$ production. Although large backgrounds are expected, the signals could be visible particularly if angular correlations are also studied. The (γ,p) experiment in nuclei would also lead to the f_0 and a_0 excitation mostly from the collision of photons with protons, since the neutrons provided only a contribution of about an order of magnitude smaller than protons for a_0 production and do not contribute to the f_0 production. The studies in nuclei would provide information on the f_0, a_0 properties in a nuclear medium, where large modifications are expected in view of present results for the modification of the $\bar{K}$ properties in a nuclear medium based on chiral unitary approaches. Such experimental studies are possible with present facilities like TJNAF and Spring8/RCNP and they would provide novel tests for our understanding of the nature of the scalar resonances and about current ideas on chiral unitary theory, which is emerging as a powerful tool for the study of meson-meson and meson-baryon interactions.

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