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COHERENT PHOTOPRODUCTION OF NEUTRAL PIONS IN NUCLEI AS A DETECTOR OF NUCLEAR PIONIC MODES

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This paper is dedicated to the memory of Professor A.B. Migdal whose pioneering work on the inclusion of pion degrees of freedom in the Fermi liquid theory of nuclei has been at the basis of most vivid discussions and continuous developments in nuclear physics.

Abstract

Common belief is that photon induced reactions are not sensitive to pion propagation effects because of the transversal-longitudinal spin mismatch. We show in this work that the coherent (γ, π^0) reaction on light nuclei can be used as a detector of the collective nuclear state at high excitation energy, called the pionic branch. Its identification had been possible up to now only through charge exchange reactions of the type ($^3\text{He}, t$) which present the usual complications of hadronic interactions and mechanisms. The better experimental knowledge of photon reactions would be of great help for the study of this new mode.

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1. Introduction

Collective states of the nucleus at very high excitation energies represent a novel aspect of nuclear dynamics. A broad state has been predicted in an energy band from the vicinity of the pion mass to $200 - 300 \text{ MeV}$, the so-called pionic branch which is constructed from many multipolarities consisting of coherent superposition of pion and Δ - h states. The prediction is based on our understanding of the propagation of physical pions and therefore rests on solid grounds. Its detection would bring information on the hitherto unexplored domain of the nucleus at very high excitation. This needs a probe which couples, as the pion, to the nucleonic spins in the longitudinal way i.e parallel to the momentum transfer. Few probes satisfy this requirement. The previously mentioned one i.e the physical pion is limited in exploration domain by the energy momentum relation $\omega^2 = q^2 + m_\pi^2$. In this situation the pionic branch does not show up directly but through the decrease of the pionic strength in the total cross-section of pions on nuclei (σ_T/A has been found¹) to vary as $A^{-0.4}$). The situation in that respect is parallel to that which existed for the Gamow-Teller strength, before the discovery of the giant Gamow-Teller state, when only a strong quenching of the G-T strength for the low lying states could be observed. For the pionic branch this evidence through the depopulation effect is so far the most convincing one.

However it is an indirect proof and one would like to display the collective state itself. This implies off-shell pions i.e the use of hadronic probes which can exchange a pion with the target. Indeed the ($^3\text{He}, t$) reaction, with the observation of the downward shift of the delta peak²), has supplied a good indication in favour of the collective state, in spite of the highly peripheral character of this probe^{3,4}). The photon would be ideal since it penetrates the nuclear interior and it probes the spin-isospin excitations, but with a transversal spin coupling instead of a longitudinal one. Thus the inclusive reactions with photons are generally not suited. The aim of this article is to show that one exclusive reaction, namely the process of coherent pion photoproduction $\gamma + A \rightarrow \pi^0 + A$, can provide in the region of photon energy around 200 MeV an evidence for the pionic branch which lies at these energies.

2. Formalism

The coherent production represents only a small fraction of the total photoabsorption cross-section since it implies two mismatches: one between the transversal coupling of the photon and the longitudinal coupling of the pion, the other between the values of the photon and pion momenta, k and $q = \sqrt{k^2 - m_\pi^2}$. These mismatches are only tolerated in light systems and will therefore limit our study to a few light nuclei, the cross-section decreasing rapidly with the mass number. Previous theoretical investigations of this process have been performed, for instance in the isobar doorway models (see e.g. ref.⁵). Our aim is not to improve over these approaches but to show that this reaction can offer an interesting way of detection of the pionic branch.

The essence of our statement is that in this process a transversal-longitudinal conversion takes place and a pionic mode is created inside the nucleus, propagates and then transforms into a physical pion. An advantage of this process is that non-resonant contributions such as the pion pole and Kroll-Ruderman terms which dominate in charged pion photoproduction reactions, are either absent or small, leaving the resonant amplitude as the only important contribution to the cross-section. This makes it well suited to explore the modifications of the Δ - h excitations. Keeping only the resonant contribution and assuming for simplicity that the process takes place on a nuclear state of spin zero $|0\rangle$, the cross-section for pion production by a photon of polarization λ writes :

$$\frac{d\sigma}{d\Omega} = \frac{q}{2k} \frac{1}{(4\pi)^2} \frac{M_A^2}{s_A} \sum_{\lambda} | \langle 0, \varphi(\mathbf{q}) | T_{\Delta} | \gamma(\mathbf{k}, \lambda), 0 \rangle |^2, \quad (1)$$

with s_A the squared total energy in the pion-nucleus c.o.m. system and the resonant amplitude T_{Δ} :

$$\langle 0, \varphi(\mathbf{q}) | T_{\Delta} | \gamma(\mathbf{k}, \lambda), 0 \rangle = \frac{f^* f_{\gamma}^*}{m_{\pi}^2} \sum_{\Delta-h} \langle 0 | O_{\alpha}^+(q) e^{i\mathbf{q}\cdot\mathbf{x}} | \Delta h \rangle G_{\Delta h} \langle \Delta h | O_{\beta}(k) e^{-i\mathbf{k}\cdot\mathbf{x}} | 0 \rangle, \quad (2)$$

$$O_{\alpha}(\mathbf{q}) = \mathbf{S}^+ \cdot \mathbf{q}^* T^0, \quad O_{\beta}(\mathbf{k}) = (\mathbf{S}^+ \times \mathbf{k}^*) \cdot \epsilon^{(\lambda)} T^0,$$

where $f^*(f_{\gamma}^*)$ is the $\pi N \Delta(\gamma N \Delta)$ coupling constant, and $S(T)$ the transition operator from $1/2$ to $3/2$ spin (isospin) states. The vector $\mathbf{q}^*(\mathbf{k}^*)$ is the pion (photon) momentum in the pion (photon)-nucleon CM frame related to that in the pion (photon)-nucleus CM by the so-called angle transformation relations (valid up to relativistic corrections, see e.g. ref.⁶):

$$\mathbf{q}^* = \frac{M\mathbf{q} - \omega\mathbf{p}_f}{M + \omega}, \quad \mathbf{k}^* = \frac{M\mathbf{k} - \omega\mathbf{p}_i}{M + \omega}$$

where $\mathbf{p}_i, (\mathbf{p}_f)$ is the moment of incoming (outgoing) nucleon. As in ref.⁷, we have assumed that on the average the nucleon momenta can be taken as those of the Breit system for an individual nucleon :

$$\mathbf{p}_i = \frac{\mathbf{q} - \mathbf{k}}{2}, \quad \mathbf{p}_f = \frac{\mathbf{k} - \mathbf{q}}{2}.$$

A multipole expansion of the amplitude can be easily performed

$$\begin{aligned} \langle 0, \varphi(\mathbf{q}) | T_{\Delta} | \gamma(\mathbf{k}, \lambda), 0 \rangle &= -i \frac{f^* f_{\gamma}^*}{m_{\pi}^2} \sum_{JM} \langle JM - \mu; 1\mu | JM \rangle Y_{JM-\mu}^*(\hat{\mathbf{k}}) Y_{JM}(\hat{\mathbf{q}}) \\ &\quad \sum_{L=J\pm 1} a_{JL} \Pi^L(k, q, \omega) b_{JL} \\ &\equiv -i \frac{f^* f_{\gamma}^*}{m_{\pi}^2} \sum_{JM} \langle JM - \mu; 1\mu | JM \rangle Y_{JM-\mu}^*(\hat{\mathbf{k}}) Y_{JM}(\hat{\mathbf{q}}) \sum_{L=J\pm 1} \Pi_{tL}^L(k, q, \omega) \end{aligned} \quad (3)$$

with

$$a_{JL} = \begin{cases} \sqrt{\frac{J}{2J+1}} & \text{for } J=L-1 \\ -\sqrt{\frac{J+1}{2J+1}} & \text{for } J=L+1 \end{cases}$$

$$b_{JL} = \begin{cases} \sqrt{\frac{J+1}{2J+1}} & \text{for } J=L-1 \\ \sqrt{\frac{J}{2J+1}} & \text{for } J=L+1 \end{cases}$$

$$a_{JL} = 0, \quad b_{JL} = 1 \quad \text{otherwise.}$$

The quantity $\Pi_{tl}(k, q, \omega)$ is the transverse-longitudinal component of spin-isospin polarization propagator $\Pi(k, q, \omega)$. This Π_{tl} is a crucial ingredient in our work and it is affected as well as the purely longitudinal component by the collective nature of the response, in a way that will be discussed later. Our first task is to evaluate the free quantity $\Pi^\circ(k, q, \omega)$. For this we have used a Wigner-Kirkwood expansion limited to first order. This avoids the cumbersome diagonalizations in the Δ - h states and allows an easy evaluation, even in heavy nuclei. In this semi-classical approximation $\Pi^\circ(k, q, \omega)$ can be cast in the form

$$\Pi^\circ(k, q, \omega) = \int dR e^{i(k-q) \cdot R/2} \Pi^\circ((k+q)/2, R, \omega) \quad (4)$$

where $\Pi^\circ((k+q)/2, R, \omega)$ is the Δ - h polarization propagator in a free Fermi gas of nucleons evaluated at a density which depends on the R position as follows. The R dependent Fermi momentum is related to the mean field potential

$$K_F(R) = \left(2M(E_F - V(R)) \right)^{1/2} \quad (5)$$

In our work a Wood-Saxon form was adopted

$$V(R) = V_0 / (1 + e^{(R-R_0)/a}) \quad (6)$$

The parameters V , a and R_0 have been taken from ref.⁸). The Δ - h polarization propagator has been written under a relativistic form

$$\Pi_{\Delta-h}^\circ = \frac{8M_\Delta}{3} \sum_j \int \frac{d\mathbf{p}}{(2\pi)^3} \left\{ \frac{1}{s - M_\Delta^2 + iM_\Delta \Gamma_\Delta} + \frac{1}{u - M_\Delta^2} \right\} \left| < \frac{3}{2}, j | T^\circ | \frac{1}{2}, \pm \frac{1}{2} > \right|^2 \quad (7)$$

u are the usual Mandelstam variables and Γ_Δ is the delta width. We have introduced the Pauli blocking according to the recipe of ref.⁹). We have also introduced the coupling of the delta to 2p-1h and 3p-2h states⁹). This incorporates the 2p-2h contributions to the polarization propagator which can be reduced to Δ self-energy graphs. We have also included the other ones through an extrapolation of calculations of physical pion absorption at threshold¹⁰). We

have limited the medium modifications of $\Pi_{\Delta-h}^{\circ}$ to its imaginary part. Indeed comparison of our approach to π -nucleus scattering has not shown the necessity of introducing a real spreading potential.

The next step is to include the effects of the Δ - h interaction. The semi-classical method is not suited to explore the collective aspect of the response. It is a nuclear matter approach and it tends to overestimate the collective effects in that energy domain. For instance it produces a complete depopulation of strength along the physical pion line, which is not realistic. We have therefore solved the RPA equations for finite systems as in previous works on the ($^3\text{He}, t$) reaction³), the semi-classical method being used only to evaluate the zeroth order polarization propagators. We have to keep in mind that this method is certainly not very appropriate to describe nuclei as light as ^4He but is used in this case only for illustration purposes. We have used the ring approximation i.e we have ignored the antisymmetrization diagrams. These can be considered as effectively included in a Landau-Migdal parameter for the Δ - h interaction sector as commonly done for the N - h case. We have solved the following system :

$$\begin{aligned}\Pi_{ll}^{(J)}(q_1, q_2, \omega) &= \Pi_{ll}^{\circ(J)}(q_1, q_2, \omega) + \int \frac{q_i^2 dq_i}{(2\pi)^3} \Pi_{ll}^{\circ(J)}(q_1, q_i, \omega) V_l(q_i) \Pi_{ll}^{(J)}(q_i, q_2, \omega) \\ &\quad + \int \frac{q_i^2 dq_i}{(2\pi)^3} \Pi_{lt}^{\circ(J)}(q_1, q_i, \omega) V_t(q_i) \Pi_{tl}^{(J)}(q_i, q_2, \omega)\end{aligned}\quad (8)$$

$$\begin{aligned}\Pi_{tl}^{(J)}(q_1, q_2, \omega) &= \Pi_{tl}^{\circ(J)}(q_1, q_2, \omega) + \int \frac{q_i^2 dq_i}{(2\pi)^3} \Pi_{tl}^{\circ(J)}(q_1, q_i, \omega) V_l(q_i) \Pi_{ll}^{(J)}(q_i, q_2, \omega) \\ &\quad + \int \frac{q_i^2 dq_i}{(2\pi)^3} \Pi_{tt}^{\circ(J)}(q_1, q_i, \omega) V_t(q_i) \Pi_{tl}^{(J)}(q_i, q_2, \omega)\end{aligned}\quad (9)$$

where V_l and V_t are the components of the Δ - h interaction in the longitudinal and transversal spin-isospin channels, respectively:

$$V_l(q) = F_{\pi}^2 \left(g'_{\Delta} + \frac{q^2}{\omega^2 - q^2 - m_{\pi}^2 + i\epsilon} \right) \quad V_t(q) = F_{\pi}^2 g'_{\Delta} + F_{\rho}^2 \left(\frac{q^2}{\omega^2 - q^2 - m_{\rho}^2 + i\epsilon} \right) \quad (10)$$

Here F_{π} and F_{ρ} are the pion and rho N - Δ coupling constants including monopole form factors:

$$F_{\pi}(q) = f^* \left(\frac{\Lambda_{\pi}^2 - m_{\pi}^2}{\Lambda_{\pi}^2 - \omega^2 + q^2} \right) \quad F_{\rho}(q) = f_{\rho}^* \left(\frac{\Lambda_{\rho}^2 - m_{\rho}^2}{\Lambda_{\rho}^2 - \omega^2 + q^2} \right)$$

with $\Lambda_{\pi} = 1. \text{GeV}$, $\Lambda_{\rho} = 1.5 \text{GeV}$ and $\frac{f_{\rho}^{*2}}{f^{*2}} = 2.3$ (strong ρ coupling).

The quantity g'_{Δ} is the Landau-Migdal parameter of the Δ - h interaction. The value that we have adopted ($g'_{\Delta} = 0.5$) is assumed to incorporate besides the genuine repulsion the effect of antisymmetrization (and maybe that of the spreading potential that we have not taken explicitly into account). We have tested the validity of our description on the similar process of π^0 elastic scattering. The results are illustrated in fig. (1) which gives the angular

distribution of this cross-section for ^{16}O at the energy $\omega = 303 \text{ MeV}$ and in fig. (2) which shows the angle integrated cross-section for the ^{16}O nucleus as a function of the energy. The total pion-nucleus cross-section is also well reproduced by our description as seen in fig. (3). We can thus be confident in the reliability of our approach.

We point out at this level the importance of the angle transformations which had been previously mentioned by other authors⁵). It introduces a multiplicative factor of $(\frac{M}{\omega + M})$ in $\Pi_{il}^o(k, q, \omega)$ and almost the same factor in the other components.

3. Results

In order to display the effect of the Δ - h interaction we have performed evaluations both with the bare and the full polarization propagator (i.e. RPA). The results for the total (γ, π^0) coherent cross-section are shown in fig. (4) for the case of ^{16}O . The RPA produces a quenching of the cross-section by a factor of about 3. Moreover the delta peak is shifted by 49 MeV towards lower energies. For other nuclei of interest the shift is found to vary from about 20 MeV in ^4He (fig. (5)) to 37 MeV in ^{12}C (fig. (6)).

We will now comment on the interpretation of these results and their relevance for the pionic branch. It is clear from the RPA equations (8) and (9) that the sensitivity of the photoproduction process to pion propagation effects arises from the second term of the r.h.s of eq. (9) which contains the longitudinal-longitudinal component of the polarization propagator in which the pionic branch is known to manifest its existence. The rôle of this branch on the cross-section has to be disentangled. We have first checked that the transversal channel (third term of eq. (9)) is not responsible of the quenching and the peak downward shift in the cross-section; it produces only a moderate repulsive effect (upward shift of the peak). The main features are produced by the second term of the r.h.s. of equation (9) where the longitudinal interaction given in eq. (10) contains the exchange of a pion that can either be on-shell or off-shell. The on-shell part corresponds to the elastic broadening where the photon transforms directly into a physical pion which subsequently undergoes an elastic scattering by the nucleus. This elastic broadening produces a quenching of the cross-section but practically no shift. This is illustrated in fig. (4) for the case of ^{16}O by the dotted curve where we have retained in the longitudinal interaction only its imaginary part $\text{Im}V_l = -\pi F_\pi^2 k^2 \delta(q^2 - \omega^2 + m_\pi^2)$, which corresponds to on-shell pions.

The shift in the peak energy is instead due to off-shell pions and more precisely to the pionic branch as shown below. In order to prove this we use fig. (7) to choose two typical energies. The first one $\omega = 220 \text{ MeV}$ corresponds to the crossing of the photon line and the pion branch ($\omega = k \approx q_{PB}$ where q_{PB} is the momentum of the pionic branch corresponding to the energy ω). For the second one $\omega = 280 \text{ MeV}$, the momenta of the two lines differ appreciably ($k = 280 \text{ MeV}/c$, $q_{PB} \approx 420 \text{ MeV}/c$). In the first case where the momenta of the photon line and the pionic branch match, the transformation of the photon into a pionic mode is favoured. This last one propagates in the nucleus before converting finally into the physical pion. The predominance of the pion branch is illustrated in the following. The

quantity which governs the cross-section of the coherent photoproduction is $\Pi_{tl}(k, q_\pi, \omega = k)$ which according to eq. (9) writes:

$$\Pi_{tl}^{(J)}(k, q_\pi, \omega) = \Pi_{tl}^{o(J)}(k, q_\pi, \omega) + \int \frac{q_i^2 dq_i}{(2\pi)^3} \Pi_{tl}^{o(J)}(k, q_i, \omega) V_L(q_i) \Pi_{ll}^{(J)}(q_i, q_\pi, \omega) \quad (11)$$

where we have for illustrative purpose omitted the transversal contribution.

In order to show which momenta contribute most in the integrand of eq. (11) we display in fig. (8) for a typical multipole $J=2$ the integrand as a function of q_i . We give its imaginary part, normalized to that of the leading term $\Pi_{tl}^{o(J)}(k, q_\pi, \omega)$. This ratio becomes large and positive in the vicinity of the pionic branch momentum q_{PB} , thus enhancing the cross-section. At $\omega = 280 \text{ MeV}$ instead the pionic branch does not show up in the integrand (see fig. (9)). The ratio alternates in sign and the overall integral is small with a resulting small effect in the cross-section. From the two above instances one understands how the pionic branch explains the shift in the coherent photoproduction cross-section.

One can wonder whether a similar shift occurs in the elastic scattering which looks at first sight quite similar to the photoproduction process. However it presents in fact an important difference concerning the number of contributing multipoles. In the scattering case a multipole of rank J is the sum of two additive contributions with $L = J \pm 1$ whereas photoproduction involves the difference between these contributions which become rapidly equal for increasing L so that they cancel and only a few multipoles actually participate into this reaction. This explains why the elastic cross-section falls less rapidly with energy than the photoproduction one. The shift of the peak is absent in the elastic case (see fig. (2)) because of the large multipolarities contribution.

We now come to the comparison with experimental data. The measured coherent cross-section for ^{12}C is shown by the full circles of fig. (6). They start to deviate from our calculation already at $\omega \approx 210 \text{ MeV}$. It is however accepted that the experiment may be contaminated by the contribution of other final nuclear states within the experimental resolution of 15 MeV . For the angular distributions which are available at 235 and 295 MeV photon energies we have added to our calculated values the estimates of Takaki et al.¹²⁾ for the excited isoscalar ^{12}C states $J^\pi = 0^+, 2^+$ and 3^+ . This restores the agreement at the two energies (see fig. (10) for $\omega = 235 \text{ MeV}$). More significant data with clear separation of the ground state are required to sign the pionic branch. It is obvious that this would be an easier task in ^4He where the absence of excited states below 20 MeV should permit a clean measurement of the coherent photoproduction cross-section. In this case our model should be improved on since it is not suited for such a light system. First one should account of $1/A$ corrections. Moreover our semi-classical approach for the bare polarization propagator should be replaced by an exact calculation. However since our RPA equations have been solved exactly we believe that the shift of the order of 20 MeV that we have found here will persist.

4. Conclusion

We have shown in this work that the collective mode at high energy known as the pionic branch affects the coherent (γ, π^0) cross-section on nuclei, shifting the maximum to a lower energy. We have tested our model on the elastic and total π -nucleus cross-sections where it gives very good results for both differential and total cross-sections (cf. figs. (1) to (3)). In the case of the more easily feasible coherent reaction on a ${}^4\text{He}$ target we predict a similar though somewhat smaller shift. According to our model the predicted shifts are sizeable (from 20 MeV in ${}^4\text{He}$ to 50 MeV in ${}^{16}\text{O}$). Their experimental display would offer a clean signature for the somewhat elusive pionic branch, this nuclear collective state of very high excitation energy. We have thus shown that contrary to the inclusive photon reactions which bring no information on the longitudinal spin-isospin channel, a specific photon induced exclusive reaction can be equivalent to a hadronic one for generating a collective effect associated with pion propagation.

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Figure Captions

- Fig. (1). The differential cross-section of pion- ^{16}O elastic scattering at 303 MeV : theory (full line) vs. average of $\pi^\pm$ experiments¹⁴⁾
- Fig. (2). The angle integrated cross-section of pion- ^{16}O elastic scattering as a function of the pion energy. The experimental points correspond to $\pi^\pm$ average¹³⁾. Full and dashed lines are for calculations with and without RPA.
- Fig. (3). The total cross-section of pion- ^{16}O and pion- ^{12}C as a function of the energy: theory (respectively dotted and dashed lines) vs. average of $\pi^\pm$ experiments¹⁵⁾
- Fig. (4). The calculated $^{16}\text{O}(\gamma, \pi^0)^{16}\text{O}$ total cross-section. The full, dashed and dotted lines represent respectively the RPA result, the bare one and the RPA with only on-shell pion propagation.
- Fig. (5). Same legend as in fig. (4) for the ^4He case.
- Fig. (6). Same legend as in fig. (4) for the ^{12}C case: experimental points are from ref. ¹¹⁾.
- Fig. (7). The position of the pionic branch illustrated by equal contour lines of the quantity $\text{Im}\tilde{\Pi}_n(q, q, \omega)$ in the (q, ω) plane. The crest is indicated by the line labelled q_{PB} . The photon line is also shown.
- Fig. (8). The ratio between the imaginary part of the integrand of eq.(11) and that of the bare quantity $\Pi^0(k, q)$ as a function of the momentum of integration q_i at $\omega = 220\text{ MeV}$ for the multipole $J=2$.
- Fig. (9). Same legend as fig. (8) at $\omega = 280\text{ MeV}$
- Fig. (10). The angular distribution of the reaction $^{12}\text{C}(\gamma, \pi^0)^{12}\text{C}$ at $\omega = 235\text{ MeV}$ uncorrected (dashed line) and corrected (full line) for the excited states calculated by Takaki et al.¹²⁾. Experimental points are from ref.¹¹⁾.

Figure 1

Figure 2

Figure 3

Figure 4

Figure 5

Figure 6

Figure 7

Figure 8

Figure 9

Figure 10