Coherent π^0 photoproduction in the isobar-hole formalism

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(Received 16 March 1979)

The isobar-hole formalism for pion-nucleus interactions is applied to coherent π^0 photoproduction on nuclei. The unifying features of this approach for elastic scattering and photoproduction are discussed, and the relationship to distorted-wave impulse appproximation calculations and the phenomenological doorway model are examined. The results are presented for $^{16}O(\gamma,\pi^0)^{16}O$ and show a large suppression from the impulse approximation prediction.

[NUCLEAR REACTIONS ^{16}O (γ, π^0) ^{16}O calculated in isobar-hole formalism; comparison to DWIA.]

I. INTRODUCTION

Recent analyses of pion-nucleus elastic scattering have led to the consensus that the first order pion optical potential is inadequate. The precise form of the higher order terms remains somewhat controversial. One way to further investigate the optical potential is to analyze other pionic reactions, such as inelastic scattering, pion absorption, or pion production from nuclei. Only a few experiments of this type have been performed, but several of these, if correct, indicate that we do not understand fully the pion-nucleus reaction mechanism.

A reaction which appears to be particularly attractive for learning about the optical potential is the coherent photoproduction of π^0 mesons from nuclei. In contrast to elastic scattering, a production reaction depends on the half-off-shell pion-nucleus transition matrix. Furthermore, since the photon interacts weakly with the target, the π^0 mesons are produced throughout the nuclear volume. Therefore, in the framework of the distorted-wave impulse approximation (DWIA), one might expect to extract the pion wave function in the nuclear interior. Saunders[1] analyzed the (γ,π^0) data of Davidson[2] using the DWIA with a standard first order pion optical potential. The results for ^{12}C are shown in Fig. 1, with the experimental cross section several times larger than that calculated. One possible resolution of this discrepancy is that the optical potential is far too absorptive. Other possible explanations lie in more complicated reaction mechanisms or in experimental difficulties.

In this note, we apply the isobar-hole formalism[3-6] for pion-nucleus interactions to the coherent (γ,π^0) process. Detailed studies of π-^{4}He and π-^{16}O elastic scattering have been performed within this framework. In these isobar-hole calculations, several dynamical aspects of pion-nucleus interactions are handled in an essentially exact manner within the framework of the nuclear shell model. These include, in addition to the usual pion multiple scattering, nucleon binding effects, isobar propagation, and Pauli blocking. More complicated mechanisms, such as pion absorption, are treated through a physically motivated phenomenology, namely, a "spreading" interaction

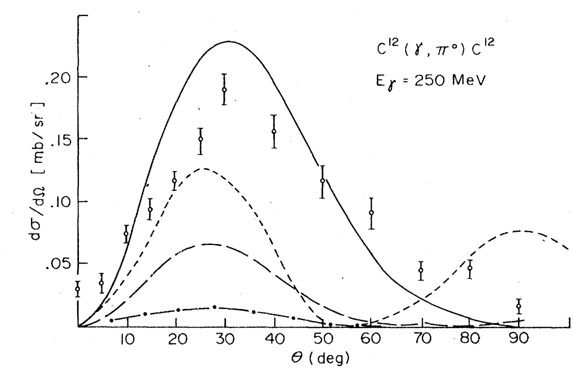

FIG. 1. Coherent π^0 photoproduction from ^{12}C at incident photon energy $E_\gamma = 250$ MeV. Data from Ref. 2. Solid curve calculated in plane wave impulse approximation. Long-dashed and dot-dash curves are DWIA calculations with local and Kisslinger pion optical potentials, respectively. These three curves are taken from Saunders (Ref. 1). The short-dashed curve is the phenomenological isobar doorway result of Woloshyn (Ref. 9).

for the isobar. The approach is reviewed briefly in Sec. II. One aim in applying the formalism to the coherent (γ, π^0) reaction is to obtain a quantitative evaluation of the process using a more refined description of the pion-nucleus interaction than has been used previously. Another aim is to demonstrate the importance of modifying consistently the production operator employed in distorted-wave calculations using higher order optical potentials. Numerical results for coherent π^0 photoproduction from ^{16}O are given in Sec. III. Section IV provides a summary of our findings.

It is important to stress the differences between the phenomenological isobar doorway approach of Kisslinger and Wang[7] and the isobar-hole formalism. The former, which has been applied to elastic scattering,[7] to analog state single charge exchange,[8] and to coherent π^0 photoproduction,[9] involves parametrization of the pion-nucleus transition matrix with a resonant form; the latter provides a vehicle for a *microscopic* calculation, with only certain "higher order" processes, such as absorption, treated approximately. The extent to which the isobar doorway phenomenology[7] is justified by the results of the microscopic calculation will be discussed in Secs. II and III.

II. BRIEF REVIEW OF THE ISOBAR-HOLE FORMALISM

At intermediate energies, πN scattering is dominated by the $J=\frac{3}{2}$, $T=\frac{3}{2}$ resonance, leading to a picture of intermediate excitation of the Δ isobar as indicated in Fig. 2(a). Pion-nucleus elastic scattering then proceeds as indicated in Fig. 2(b): The pion generates Δ-nucleon hole states and the scattering process is given by appropriate matrix elements of the Δ-hole propagator $G_{\Delta h}$. (References 3 and 5 contain detailed accounts of the material in this section.) This can be made explicit quite easily in the projection operator formalism. We define the P space as containing all states consisting of the pion plus ground state nucleus. The D space contains all Δ-hole states, and the Q space consists of all more complicated channels. Given the doorway condition $H_{PQ}=0$, the pion-nucleus elastic transition matrix is given by

$$T_{PP} = H_{PD} G_{\Delta h} H_{DP}$$
$$= H_{PD}[E - E_R + \tfrac{1}{2} i\Gamma(E) - \mathcal{H}_{\Delta h}]^{-1} H_{DP}\,, \tag{1}$$
$$\mathcal{H}_{\Delta h} = H_{DD} + H_{DD}^{\uparrow} + H_{DD}^{\downarrow}\,, \tag{2}$$
$$H_{DD}^{\uparrow} = H_{DP}(E^+ - H_{PP})^{-1} H_{PD}\,, \tag{3}$$
$$H_{DD}^{\downarrow} = H_{DQ}(E^+ - H_{QQ})^{-1} H_{QD}\,. \tag{4}$$

Here H_{DP} represents the π-Δ-h vertex function and the Breit-Wigner denominator $(E - E_R + i\Gamma/2)$ represents the free Δ propagator. The interaction $H_{DD}^{\uparrow}$ generates the elastic width and is represented by Fig. 3(a). The spreading interaction $H_{DD}^{\downarrow}$ sums intermediate coupling to the more complicated Q space. One example of a spreading interaction is indicated in Fig. 3(b); this corresponds to pion absorption, since the intermediate state is a nuclear 2p-2h state. The isobar-hole Hamiltonian $\mathcal{H}_{\Delta h}$ also contains a diagonal interaction H_{DD} which includes isobar propagation and binding effects, Pauli blocking effects, and isobar-hole residual interactions. We can define the normalized doorway state for the Lth pion-nucleus partial wave as

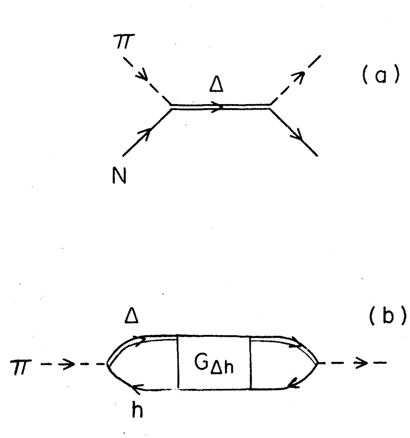

FIG. 2. (a) Pion-nucleon interaction via Δ excitation. (b) Pion-nucleus elastic scattering in the isobar-hole picture.

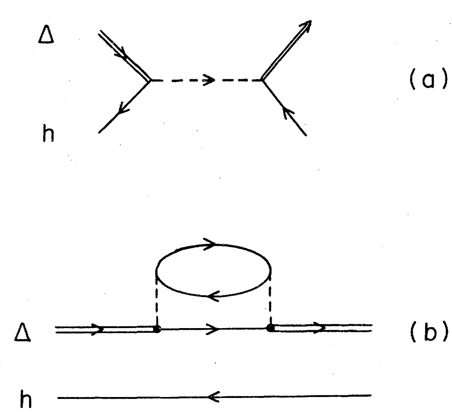

FIG. 3. (a) Pion exchange interaction $H_{DD}^{\uparrow}$. (b) Example for spreading interaction $H_{DD}^{\downarrow}$: coupling to 2p-2h states (pion absorption).

$$|D_0^L\rangle \equiv N_L H_{DP}|(\vec{q})_L;0\rangle\,,\qquad N_L=\langle(\vec{q})_L;0|H_{PD}H_{DP}|(\vec{q})_L;0\rangle^{-1/2}\,, \tag{5}$$

where $|(\vec{q})_L;0\rangle$ represents the partial wave projection of the pion plane wave (wave number $\vec{q}$) and the nuclear ground state (we deal only with $J^\pi T=0^+0$ ground states). The doorway state is a linear superposition of Δ-hole states. The pion-nucleus partial wave transition matrix is simply the expectation value of $G_{\Delta h}$ in the doorway state:

$$T_L=N_L^{-2}\langle D_0^L|G_{\Delta h}|D_0^L\rangle\,. \tag{6}$$

We stress that the doorway state generally is *not* an eigenstate of the Hamiltonian $\mathcal{H}_{\Delta h}$. If it were an eigenstate, the partial wave transition matrix would simply be

$$T_L^{\text{Doorway}}=\frac{N_L^{-2}}{E-E_R+\frac{1}{2}i\Gamma-\langle D_0^L|\mathcal{H}_{\Delta h}|D_0^L\rangle}\,. \tag{7}$$

The phenomenological doorway model[7] *assumes* this form and fits the expectation value to experiment. In practice, an L-independent parametrization is generally used.[7-9]

To make detailed comparisons with data, pion nonresonant interactions must be included also. These can be handled[3,5] within the two-potential framework. However, in the absence of high precision (γ,π^0) data, we shall keep only the resonant interaction throughout this work, thereby gaining some simplicity in the formal expressions and in the calculations. In the energy region of interest here, inclusion of the nonresonant terms would have little effect.

A convenient orthonormal basis[4,5] for evaluating Eq. (6) can be constructed as follows ($i\geq 1$):

$$|D_i^L\rangle \equiv N_i^L\left[\mathcal{H}_{\Delta h}|D_{i-1}^L\rangle-\sum_{j=0}^{i-1}|D_{i-1}^L\rangle\langle D_{i-1}^L|\mathcal{H}_{\Delta h}|D_j^L\rangle\right], \tag{8}$$

where N_i^L is a normalization factor. The transition matrix then reads

$$T_L=\cfrac{N_L^{-2}}{E-E_R+\frac{1}{2}i\Gamma-\mathcal{H}_{00}^L-\cfrac{(\mathcal{H}_{01}^L)^2}{E-E_R+\frac{1}{2}i\Gamma-\mathcal{H}_{11}^L-\cfrac{(\mathcal{H}_{12}^L)^2}{E-E_R+\frac{1}{2}i\Gamma-\mathcal{H}_{22}^L-\cdots}}}\,, \qquad \mathcal{H}_{ij}^L\equiv\langle D_i^L|\mathcal{H}_{\Delta h}|D_j^L\rangle\,. \tag{9}$$

Note that the leading term in the continued fraction conforms to the phenomenological doorway prescription [Eq. (7)]. The microscopic calculations proceed by constructing the first few states in the doorway basis $|D_i^L\rangle$ and evaluating the relevant expectation values in Eq. (9). An important result is that the continued fraction converges to very high accuracy with retention of only a few states in the doorway basis.[4,5] For example, the result for T_L with only one state (i.e., the doorway) already is accurate to about 15%, while three states give accuracies better than 1%. This rapid convergence offers some justification for the phenomenological doorway prescription for the partial wave amplitudes. However, it will be seen below that the matrix elements of $\mathcal{H}_{\Delta h}$ are strongly L dependent, so that angular distributions are far from scaled versions of the impulse approximation. Further, it must be remarked that this rapid convergence need not hold true for inelastic reactions.

Before leaving this section, it is worthwhile to examine the behavior of the doorway expectation value $\mathcal{H}_{00}^L$, since this gives a reasonable approximation to the partial wave amplitude. In particular, it is interesting to look at the expectation values of the various dynamical ingredients in $\mathcal{H}_{\Delta h}$. The imaginary part is shown in Fig. 4 for the nuclear 1^+ and 4^- partial waves. (The spreading interaction is the same as that used in Ref. 4.) Note the considerable variation with L. For the

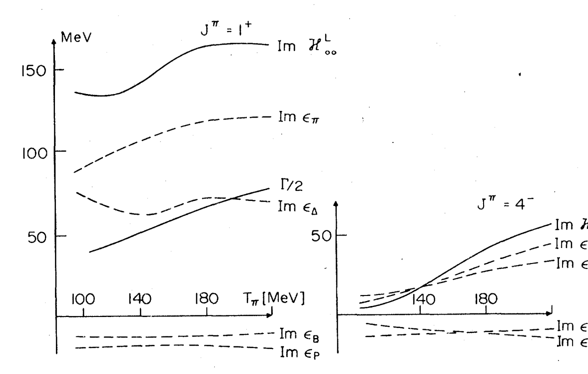

FIG. 4. Contributions to imaginary part of $\mathcal{H}_{00}^L \equiv\langle D_0^L|\mathcal{H}_{\Delta h}|D_0^L\rangle$. Im ϵ_π: elastic width due to $H_{DD}^\dagger$; Im ϵ_Δ: broadening from spreading potential and Δ propagation; Im ϵ_P: Pauli quenching; Im ϵ_B: nucleon binding effects. (a) $L^\pi=1^+$ state. (b) $L^\pi=4^-$ state.

1^+, the pion "exchange" and spreading terms, ϵ_π and ϵ_Δ respectively, dominate; the former ensures strong collectivity for the doorway state, while the latter greatly increases the ratio of inelastic to elastic scattering. For the 4^-, the binding and Pauli terms, ϵ_B and ϵ_F respectively, become comparatively much more important.

III. COHERENT π^0 PHOTOPRODUCTION

The structure of the coherent π^0 photoproduction amplitude in the isobar-hole formalism is clear from Fig. 5. The isobar-hole propagator is identical to that calculated for elastic scattering [see Fig. 2(b)], and one needs to compute only the photon coupling to the Δ-hole states. This feature unifies the coherent (γ, π^0) and elastic scattering calculation. Formally, we have

$$T_{\gamma\pi^0} = H_{PD}\frac{1}{E - E_R + \frac{1}{2}i\Gamma - H_{DD} - H_{DD}^{\prime} - H_{DD}^{\prime}} H_{D\gamma}\,. \tag{10}$$

The coupling $H_{D\gamma}$ is generated by the $\gamma N \Delta$ vertex, which reads in the Δ rest frame

$$G_{\gamma N\Delta} \equiv g_\gamma \hat{\epsilon}(\vec{k}, \lambda)\cdot\vec{k}\times\vec{S}T_3\,, \tag{11}$$

where $\hat{\epsilon}$ is the photon polarization vector and $\vec{k}$ is the photon wave number. The $\pi N\Delta$ spin and isospin transition operators are denoted by $\vec{S}$ and $\vec{T}$, respectively. The coupling constant $g_\gamma = 0.165f$ is chosen to fit the experimental $M_{1+}(\frac{3}{2})$ multipole at resonance.[10] Defining the isobar doorway state for photoexcitation $|D_\gamma^L\rangle = N_L^\gamma H_{D\gamma}|(\vec{k})_L, 0\rangle$, we have an obvious generalization of Eq. (6):

$$T_L^{\gamma\pi^0} = (N_L^\gamma N_L)^{-1}\langle D_0^L|G_{\Delta h}|D_\gamma^L\rangle$$

$$= (N_L^\gamma N_L)^{-1}\sum_i \langle D_0^L|G_{\Delta h}|D_i^L\rangle\langle D_i^L|D_\gamma^L\rangle\,. \tag{12}$$

In the isobar-hole formalism, the "new" physics probed in the (γ, π^0) reaction are the off-diagonal expectation values of the Δ-hole propagator. In optical potential language, this is equivalent to the fact that elastic scattering determines the asymptotic pion wave function (i.e., phase shifts), while a production reaction is sensitive to the wave function inside the nucleus. The extent to which the (γ, π^0) reaction is sensitive to the pion-nucleus off-shell elastic amplitude is determined by the overlap of the photon doorway state with the doorway states $|D_i^L\rangle$. Comparison of Eq. (6) and (12) makes clear the sense in which a unified treatment of elastic scattering and inelastic reactions is provided in the isobar-hole approach.

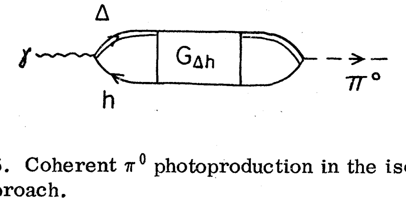

FIG. 5. Coherent π^0 photoproduction in the isobar-hole approach.

In Fig. 6, we show the total coherent $^{16}O(\gamma, \pi^0)^{16}O$ cross section as a function of energy. Compared to the impulse approximation, the cross section in the resonance region is reduced by about a factor of 7, which is comparable to the reduction obtained by Saunders.[1] The angular distribution for incident photon energy corresponding to outgoing pion kinetic energy $T_{\pi^0} = 140$ MeV is shown in Fig. 7. The same reduction factor with respect to the impulse approximation is evident here. Recalling the experimental result[2] mentioned in the Introduction, this is in sharp disagreement with the Davidson data for the $^{12}C(\gamma, \pi^0)^{12}C$ reaction. This disagreement will be addressed below.

As remarked earlier, very few doorway states were needed to obtain the elastic transition matrix.[3-5] We show in Figs. (6) and (7) the results obtained by saturating the sum in Eq. (12) with only the doorway state $|D_0^L\rangle$ [this is analogous to Eq. (7)]. This is remarkably close to the full result, demonstrating again the very compact description of pion-nucleus interactions in the Δ-hole approach.

It is important to understand how the convergence in the doorway expansion comes about. Table I

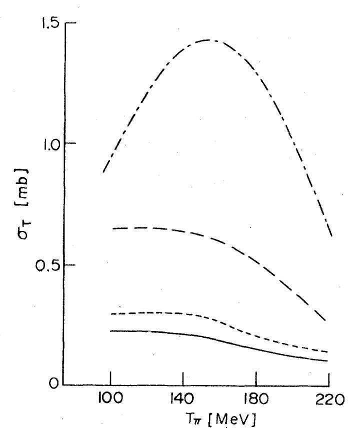

FIG. 6. Total coherent π^0 photoproduction cross section from ^{16}O: solid line shows results of full calculation; short-dashed line, of the one-doorway approximation calculation; and dot-dash line, of the plane wave impulse approximation calculation. The long-dashed line is the result of a DWIA calculation using the full isobar-hole pion-nucleus transition matrix; many-body effects in the production operator [see Eq. (16)] are not included. The kinetic energy of the produced pion is denoted by T_π.

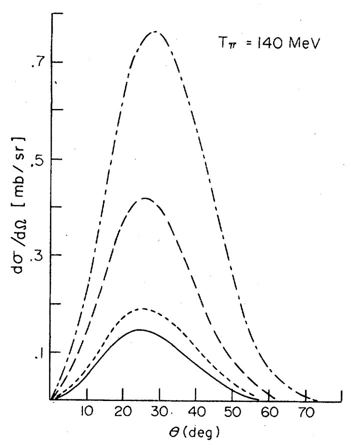

FIG. 7. Coherent $^{16}O(\gamma, \pi^0)^{16}O$ differential cross section for $T_\pi = 140$ MeV. Labeling of curves same as for Fig. 6.

gives the overlap of the photon doorway state with the Δ-hole state $|D_i^L\rangle$ and the off-diagonal matrix elements of the propagator $\langle D_0^L | G_{\Delta h} | D_i^L\rangle$ for $i = 0, \ldots, 5$ and $L^\pi = 1^+$ and 4^-. These quantities determine the (γ, π^0) transition operator given in Eq. (12). For the semiperipheral partial wave $L^\pi = 4^-$, the states $|D_0^L\rangle$ and $|D_\gamma^L\rangle$ have very good overlap so that the partial wave contribution is given almost entirely by the diagonal matrix element $\langle D_0^L | G_{\Delta h} | D_0^L\rangle$. Since the latter is determined in elastic scattering, the (γ, π^0) results for these partial waves are essentially model independent in that all reasonable theories providing the same elastic phase shift will give the same (γ, π^0) transition matrix element within about 10%. On the other hand, the overlap $\langle D_0^L | D_\gamma^L\rangle$ is very small for the central partial wave $L = 1^+$. In fact, for the central partial waves, convergence of the sum in Eq. (12) comes primarily from the off-diagonal matrix elements $\langle D_0^L | G_{\Delta h} | D_i^L\rangle$. Therefore, the central partial waves are quite model dependent and may differ markedly even for reasonable elastic phase-equivalent theories. Of course most of the (γ, π^0) cross section comes from the semiperipheral and peripheral partial waves, so that microscopic theories providing equivalent descriptions of elastic scattering will not differ greatly in predictions of the coherent (γ, π^0) cross section.

The marked difference in the overlap $\langle D_0^L | D_\gamma^L\rangle$ for different L values can be understood (somewhat formally) by examining the impulse approximation. We have

$$\langle 0; \vec{q} | T_{\gamma\pi^0} | 0; \vec{k}\rangle = \sum_{\substack{L\geq 1\\ \lambda=\pm 1}} Y_{L\lambda}^*(\hat{q}) P_{L,\lambda}\,, \tag{13}$$

$$P_{L,\lambda} = i\lambda (4\pi)^{3/2} \frac{g_\pi g_\gamma}{\sqrt{2}} \frac{h(q^2)}{E - E_R + i\Gamma/2} \left(\frac{L(L+1)}{2L+1}\right)^{1/2} \times [\rho_{L+1}(k,q) - \rho_{L-1}(k,q)]\,,$$

$$\rho_L(k,q) \equiv \int_0^\infty dr\, r^2 j_L(kr) j_L(qr) \rho(r)\,, \tag{14}$$

where the incoming photon momentum $\vec{k}$ is chosen as the polar axis, $g_\pi h(q^2)$ describes the $\pi N \Delta$ vertex, and $\rho(r)$ is the nuclear ground state density. For comparison, the pion elastic scattering impulse approximation is given by

$$\langle 0; \vec{q}\,' | T_\pi | 0; \vec{q}\rangle = \sum_{L\geq 0} Y_{L0}(\hat{q}) T_L\,, \tag{15}$$

$$T_L = (4\pi)^{3/2} \frac{g_\pi^2 h^2(q^2)}{E - E_R + \frac{1}{2} i\Gamma} \frac{1}{(2L+1)^{1/2}} \times [(L+1)\rho_{L+1}(q,q) + L\rho_{L-1}(q,q)]\,.$$

The photon and pion momenta $\vec{k}$ and $\vec{q}$ are almost equal in the resonance region, so that the primary difference between Eqs. (15) and (13) is the relative sign of the two moments of the nuclear density ρ_{L+1} and ρ_{L-1}. This is a direct reflection of the different isobar couplings: the pion couples lon-

TABLE I. Overlap of photon doorway state $|D_\gamma^L\rangle$ with pion doorway basis states $|D_i^L\rangle$, defined in Eq. (8), and matrix elements of the Δ-h propagator for $L^\pi = 1^+$ and 4^- and for kinetic energy $T_\pi = 140$ MeV.

	$L^\pi = 1^+$		$L^\pi = 4^-$	
i	$\langle D_i^L \| D_\gamma^L\rangle$	$\langle D_0^L \| G_{\Delta h} \| D_i^L\rangle$ (GeV^{-1})	$\langle D_i^L \| D_\gamma^L\rangle$	$\langle D_0^L \| G_{\Delta h} \| D_\gamma^L\rangle$ (GeV^{-1})
0	-0.29	$0.07 + 4.7i$	-0.91	$6.1 + 11.0i$
1	$-0.31 + 0.24i$	$2.5 - 0.96i$	$0.07 + 0.15i$	$2.1 - 1.5i$
2	$0.64 + 0.17i$	$0.15 - 0.87i$	$0.12 + 0.16i$	$-0.06 - 0.89i$
3	$-0.02 + 0.16i$	$0.05 - 0.05i$	$-0.27 + 0.15i$	$0.17 - 0.21i$
4	$0.06 - 0.12i$	$-0.06 - 0.01i$	$0.24 + 0.34i$	$0.01 - 0.03i$
5	$-0.07 - 0.27i$	$0.03i$	$0.09 + 0.10i$	-0.02

gitudinally, $\vec{q}\cdot\vec{S}$, while the photon couples transversely, $\hat{\epsilon}\times\vec{k}\cdot\vec{S}$. Figure 8 shows $\rho_L(q,q)$ as a function of L and pion kinetic energy. For small L and high energy, $\rho_{L+1}\approx\rho_{L-1}$, so that the overlap $\langle D_0^L | D_\gamma^L\rangle$ is very small. For larger L, $\rho_{L-1}\ll\rho_{L+1}$, so that the two (normalized) doorway states are virtually identical. This is the origin of the behavior demonstrated in Table I.

It is interesting to compare the expression for $T^{\gamma\pi^0}$ in the isobar-hole approach with that in a more conventional calculation using pion distorted waves. Using Eqs. (1)–(3) and (10), the photoproduction operator can be written as

$$T_{\gamma\pi^0}=(T_{PP}G_P+1)\times\left\{H_{PD}\frac{1}{E-E_R+\frac{1}{2}i\Gamma-H_{DD}-H_{DD}^{\dagger}}H_{D\gamma}\right\}. \tag{16}$$

This is very similar to the DWIA form, with the quantity in parentheses generating the pion distorted wave in the optical potential and that in curly brackets describing π^0 production through the isobar. However, a very important difference is that the production operator contains the doorway space interactions H_{DD} and $H_{DD}^{\dagger}$, corresponding to a production operator modified by the nuclear medium. These can be understood as multistep contributions in the conventional distorted-wave approach and are connected directly to the elastic scattering spreading interaction in the isobar formalism. Consistent treatment of the production operator and of the distorted wave is very important. This can be seen by replacing the production operator in Eq. (16) by the free $\gamma+N\to\pi^0+N$ transition operator. This corresponds to the DWIA approach with the distortion effected by the full optical potential including, for example, the spreading interaction. These DWIA results are also shown in Figs. 6 and 7. Clearly, the spreading interaction in the production operator suppresses the calculated cross section considerably. It is important to stress that a DWIA calculation using the first order pion optical potential gives results quite similar to those obtained in the full isobar-hole calculation. The interpretation of these results is clear. The full pion optical potential leads to considerably less attenuation of the pion wave function than does the first order optical potential; this accounts for the substantial increase of the DWIA cross section when the full distortion is used. This increase is compensated largely by reduction of the cross section owing to the modified production operator in Eq. (16). Consequently, one sees the importance of a consistent treatment of the production operator and of the distorting optical potential for the pion wave function.

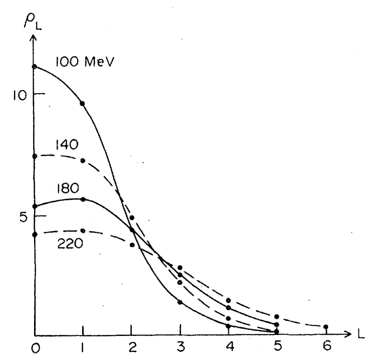

FIG. 8. Moments of the nuclear density $\rho_L(q,q)$ [see Eq. (14)] in arbitrary units for several pion kinetic energies T_π.

This leaves for discussion the Davidson result that the $^{12}C(\gamma,\pi^0)^{12}C$ cross section is almost equal to the impulse approximation result (see Fig. 1). In light of our results above, we strongly urge that the experiment be repeated. The Davidson experiment had extremely poor energy resolution: The photon detectors had energy resolution ±12% and ±20% and a bremsstrahlung beam was used. The main argument used to support the *coherence* of the measured π^0 spectrum was the angular dependence. However, the cross section at large angles failed to show the expected diffraction structure, indicating that at least some of the cross section was incoherent. Of course, in the extreme that a totally inclusive reaction is measured, and excluding true pion absorption and charge exchange, the nuclear cross section would be large and approximately equal to the $\gamma+N\to\pi^0+N$ cross section times the number of nucleons. We do not mean to imply that this is sufficient to explain the discrepancy; indeed the data appear to be too high even for an incoherent process with pion absorption and charge exchange included.

One last topic of discussion involves application of the phenomenological isobar doorway model[7] to coherent π^0 photoproduction. This has been done by Woloshyn.[9] His results are closer to the data than are either the results of Saunders[1] or the results of our microscopic isobar-hole approach. However, Woloshyn's calculation has the very serious shortcoming (recognized by the author) that the phenomenological widths and shifts chosen for the nuclear amplitude [the analog to Eq. (7) for the (γ,π^0) reaction] are independent of L. We have seen that, even in the one-doorway approximation, these should be strongly L dependent (see

Fig. 4). Furthermore, the one-doorway approximation upon which the phenomenological form rests is very poor for the central partial waves. One may expect a drastically incorrect angular distribution. Indeed, Woloshyn finds a large cross section for backward angles, more or less in agreement with the Davidson data (see Fig. 1). This cross section is merely the result of the approximations made and is not a success of the model. As argued above, we expect that the backward angle data reflects incoherent production.

IV. CONCLUDING REMARKS

We have discussed coherent π^0 photoproduction in the resonance region using the isobar-hole formalism. Our expression Eq. (12) is written in terms of matrix elements of the isobar-hole propagator, which has been used previously to describe elastic pion scattering, and in terms of the overlap between photon and pion generated doorway states. In the more usual optical potential language, the calculation differs from standard DWIA calculations by including higher order terms, consistently, both in the distorting optical potential and in the production operator. The many-body modification $H_{DD}^{\dagger}$ of the production operator corresponds to a summation of multistep processes in terms of the Δ hole spreading interaction fit to elastic scattering.[4,5]

Our isobar-hole results for $^{16}O(\gamma, \pi^0)^{16}O$ are qualitatively similar to DWIA results using a simple first order optical potential. Both calculations give cross sections roughly an order of magnitude less than that given by the impulse approximation. This rough agreement is due to a competition between two effects: While the pion wave function is considerably less damped in the isobar-hole calculation, the π^0 photoproduction is concomitantly suppressed. Ignoring the many-body modification of the photoproduction operator would lead to a much larger cross section, apparently in closer agreement with the old Davidson data (compare Figs. 1 and 7). This large effect points to the importance of a microscopic understanding of the optical potential used in DWIA calculations.

As in elastic scattering, rapid convergence is achieved in terms of Δ-hole doorway states, so that elastic scattering and (γ, π^0) can be "unified" with only a few parameters in each partial wave. However, the phenomenological isobar doorway approach using L-independent parameters[9] leads to serious errors in the evaluation of Eq. (10).

For the peripheral and semiperipheral partial waves, the overlap $\langle D_0^L | D_\gamma^L \rangle$ is close to one. Consequently, the corresponding (γ, π^0) amplitudes are determined almost completely by the pion-nucleus elastic scattering phase shifts. Since these partial waves dominate the cross section, discrimination between different theories is likely to require fairly accurate absolute normalization of experimental cross sections and rather good angular distributions.

This work was supported in part by the Foundation for Fundamental Research (FOM) and the Netherlands Organization for the Advancement of Pure Research (ZWO) and supported in part through funds provided by the U.S. Department of Energy (DOE) under Contract No. EY-76-C-02-3069.

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