Role of the $\Delta(1232)$ in the (γ,p) reaction

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The magnetic Born and intermediate-Δ contributions to the (γ,p) reaction are calculated for ^{4}He and ^{16}O targets, within a framework which includes pion rescattering and proton distortions. The results are discussed in light of experimental data, which show no obvious resonant energy dependence for these nuclei. The calculations are also compared to several published theoretical approaches, which disagree with each other regarding the importance of the intermediate-Δ contribution to the (γ,p) reaction at intermediate energies.

I. INTRODUCTION

The (γ,p) reaction at intermediate energies has been the subject of considerable theoretical and experimental interest.[1] One of the original intentions was to use this reaction to learn about the high momentum behavior of nucleons in nuclei. In a simple plane-wave impulse approximation (PWIA), the (γ,p) differential cross section is proportional to $|\psi(\mathbf{q})|^2$, where $\psi(\mathbf{q})$ is the bound proton momentum-space wave function evaluated at the momentum $\mathbf{q}$ acquired by the bound nucleon, and $\mathbf{q}=\mathbf{p}-\mathbf{k}$ for outgoing proton momentum $\mathbf{p}$ and photon momentum $\mathbf{k}$. Naively, therefore, one would expect a (γ,p) experiment to be a direct measure of the high-momentum content of single-nucleon wave functions. However, a comparison of ^{16}O$(\gamma,p)^{15}$N and ^{16}O$(\gamma,n)^{15}$O data at $E_\gamma \simeq 80$ MeV clearly indicates that the PWIA is inadequate.[2,4] In the PWIA description of (γ,N) reactions, the contribution from the nucleon convective current dominates over that from the magnetic current (except for extreme forward or backward scattering angles) for $E_\gamma \lesssim 100$ MeV. One therefore expects the (γ,p) cross section to be much larger than the (γ,n) cross section. However, to the contrary, the experimental data show that the two cross sections are approximately equal. This strongly suggests that reaction mechanisms involving two or more nucleons must be important. Recent theoretical investigations[2,3] do confirm that for $E_\gamma \gtrsim 100$ MeV, the two-nucleon current contribution can be sizable, if not completely dominant. At $E_\gamma \simeq 300$ MeV, the photon can also excite a nucleon to a $\Delta(1232)$.[5] Given the large amplitude for this process, one expects it to enhance the two-body current in this energy region.

Despite the substantial effort invested in the (γ,p) reaction to date, our present understanding of this process is not yet clear. While it is generally agreed that the two-body current gives a large contribution for photon energies of several hundred MeV, its important elements are still a matter of debate. For example, Londergan and Nixon[2] have evaluated the intermediate-Δ contribution to ^{16}O$(\gamma,p)^{15}$N$_{\text{g.s.}}$, and found it to be dominant in the Δ-resonance energy region. However, Gari and Hebach[3] found that one-body and nonresonant exchange currents are very important, even for $E_\gamma \simeq 300$ MeV. They also claimed that the Δ current does not dominate, and its inclusion is not needed to explain the general feature of the data. For the energy dependent differential cross section, the data show a characteristic Δ peak at $E_\gamma \simeq 300$ MeV only for the case of $\gamma d \rightarrow np$,[6] and not for heavier nuclei. The absence of an obvious peak means that one must turn to a calculation to examine quantitatively the effect of the Δ. Regrettably, it is precisely this effect about which the latter two previous approaches disagree.

In light of the present status of both theory and experiment, we have in this paper a specific objective in mind, namely, to reexamine the contribution of the intermediate-Δ term to the intermediate energy (γ,p) reaction, using an approach which was recently applied to the (p,π) reaction at comparable energies.[7] The main ingredients of the calculation are described in Sec. II. In Sec. III, we compare calculated results to data for ^{16}O(γ,p) ^{15}N$_{\text{g.s.}}$ and ^{4}He$(\gamma,p)^3$H. We find that the Δ is an important element of the two-nucleon current, but does not dominate the energy dependence of the cross section. We also compare our approach to the calculations of Refs. 2 and 3. The possibilities for improved calculations are discussed in Sec. IV.

II. INGREDIENTS OF THE MODEL

Our calculation of the (γ,p) reaction is summarized graphically by the diagrams shown in Fig. 1. In addition to the one-nucleon Born term [Fig. 1(a)], there are two-nucleon intermediate-Δ contributions, which we call direct [Fig. 1(b)] and exchange [Fig. 1(c)].

The dynamical picture is exactly the same as that used by Keister and Kisslinger to describe the (p,π) reaction at intermediate energies. The details of that approach (including angular momentum decomposition, etc.) are described in Ref. 7. The adaptation to (γ,p) reactions consists of replacing a distorted-wave pion with a plane-wave photon, and changing the $\pi N\Delta$ vertex to a $\gamma\pi\Delta$ vertex. In the following paragraphs, we outline the salient features of our approach, particularly as they pertain to comparisons either with experimental data or to other calculations.

A. Overview

The (γ,p) differential cross section can be decomposed as

$$\frac{d\sigma}{d\Omega}=\left.\frac{d\sigma}{d\Omega}\right|_{\rm el}+\left.\frac{d\sigma}{d\Omega}\right|_{\rm mag}, \tag{2.1}$$

corresponding to contributions from the electric and magnetic portions of the nuclear current. The one-nucleon current appearing in the Born diagram has a convective electric part,

$$H_{\rm el}^{(1)}=\frac{e}{2M}(1+\tau_3)\mathbf{p}\cdot\boldsymbol{\epsilon}\ , \tag{2.2}$$

where M is the nucleon mass, $\mathbf{p}$ the nucleon momentum, and $\boldsymbol{\epsilon}$ the photon polarization vector; and a magnetic part,

$$H_{\rm mag}^{(1)}=i\frac{e}{2M}(\mu_S+\mu_V\tau_3)\boldsymbol{\sigma}\cdot\mathbf{k}\times\boldsymbol{\epsilon}\ , \tag{2.3}$$

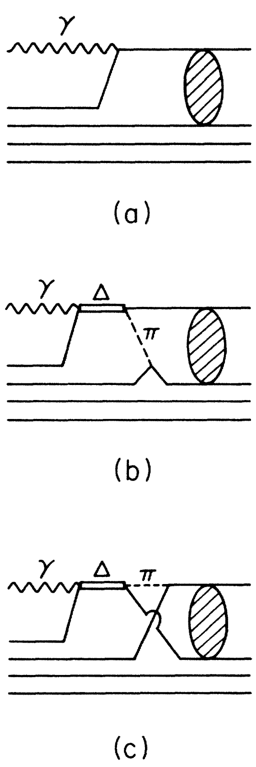

FIG. 1. (γ,p) reaction mechanisms considered in this work. (a) Born term; (b) direct-Δ term; (c) exchange-Δ term.

where μ_S (μ_V) is the nucleon isoscalar (isovector) magnetic moment and $\mathbf{k}$ the photon momentum. In the intermediate-Δ diagrams, the photon couples to the Δ via the magnetic interaction

$$H_{\rm mag}^{(\Delta)}=i\frac{e^*}{2M_\Delta}T_3\boldsymbol{\Sigma}\cdot\mathbf{k}\times\boldsymbol{\epsilon}\ , \tag{2.4}$$

where $M_\Delta=1232$ MeV, e^* is an effective charge, and $\Sigma(T)$ is a $\frac{1}{2}\rightarrow\frac{3}{2}$ spin (isospin) transition operator normalized such that

$$\langle\tfrac{3}{2}||\Sigma||\tfrac{1}{2}\rangle=\langle\tfrac{3}{2}||T||\tfrac{1}{2}\rangle=1$$

using the convention of Brink and Satchler.[8] We let $e^*=6.6e$, so as to reproduce the Δ-resonance portion of the $\gamma N\rightarrow\pi N$ amplitude.[5]

In our calculations, we retain the *magnetic* contributions only. We do this for two reasons. First, while the electric one-nucleon current dominates at low energies, it is much smaller than the (magnetic) intermediate-Δ terms at $E_\gamma\simeq300$ MeV, where our primary interest lies. Second, while the magnetic terms are individually gauge invariant, the electric terms are not. There exist ways of rendering the electric terms gauge invariant, but there are also other accompanying problems, as discussed in Sec. IV, and we therefore do not include these terms.

The intermediate-Δ terms involve an integral over the virtual pion momentum. In principle, the Δ subenergy (i.e., the energy at which the $\gamma N\rightarrow\pi N$ amplitude is to be evaluated) depends upon this momentum. In our calculations, the Δ subenergy is fixed at a value determined by external kinematics. The result is that our transition amplitude does not reflect certain nonlocality properties of the Δ which are contained in the correct momentum integral. The main advantage of factorizing the Δ subenergy is that the calculations become simpler and can be performed in coordinate space. Proton distortions can then be included by using an optical potential without introducing any extra integrals.

The propagation of the intermediate pion requires further discussion. Since relativistic kinematics are combined with nonrelativistic wave functions, there is an ambiguity in the choice of the fourth component of the exchanged pion momentum, k_π^0. We choose *spectator* kinematics in which the intermediate nucleus is placed on its mass shell. In the limit of static, heavy nuclei, the fourth component is

$$k_\pi^0(\text{direct})\simeq0\ , \tag{2.5}$$

$$k_\pi^0(\text{exchange})\simeq E_\gamma\ . \tag{2.6}$$

Since the static spectator core of the nucleus absorbs no energy, the active nucleon [the uppermost line to the left of the shaded oval in Figs. 1(b) and (c)] must absorb the photon energy, which is carried by the pion in the exchange case, but not in the direct case. For the direct term, pion propagation is always spacelike, but for the exchange term, propagation can be spacelike or timelike, depending on whether E_γ is less than or greater than m_π. Physically, this reflects the presence of the threshold for real pion photoproduction from nuclei.

B. Physical input

As mentioned earlier, our approach follows closely that of Ref. 7: a (p,π) reaction is changed to a (γ,p) reaction by replacing a distorted wave pion with a plane-wave photon. In the following, we indicate the physical input and parameters used in specific calculations performed.

For the reaction $^{16}O(\gamma,p)^{15}N_{g.s.}$, we assume ^{16}O to be a closed $1p$-shell nucleus, and take intermediate mass-15 and mass-14 states to be pure configurations in the $1p$ shell. Woods-Saxon wave functions for the bound nucleons are obtained using the parameters of Elton and Swift[9] which give a fit to elastic electron scattering from ^{16}O. The p-^{15}N optical potential is taken from a recent analysis of p-^{16}O scattering in the 200 MeV region, in which an energy-dependent parametrized optical potential is obtained.[10]

For the case $^4He(\gamma,p)^3H$, we consider only pure $1s$ shell configurations. The single-nucleon wave functions are obtained with a harmonic oscillator potential adjusted to fit elastic electron scattering data for ^4He.[11] The distortions in the p-^3H final state are ignored.

III. RESULTS AND DISCUSSION

A. Comparison with experiment

Our predicted energy distributions for the $^{16}O(\gamma,p)^{15}N_{g.s.}$ reaction are shown for θ_p=45°, 90°, and 135° in Figs. 2—4. The three curves illustrate the contri-

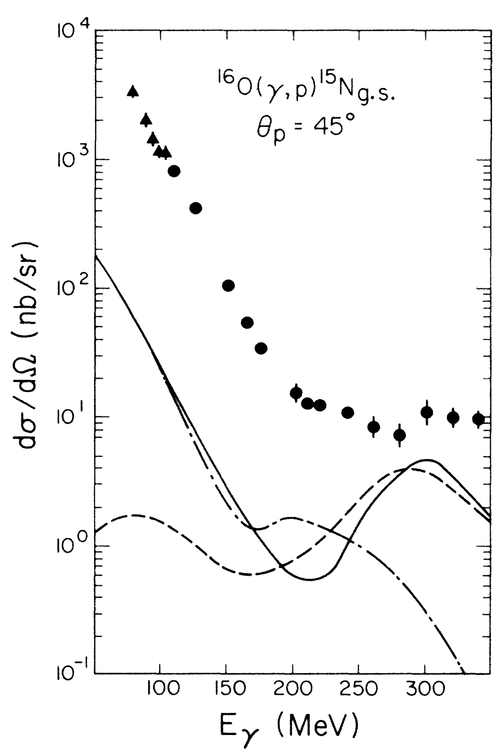

FIG. 2. Energy distributions for $^{16}O(\gamma,p)^{15}N_{g.s.}$ at θ_p(lab)=45°. Dashed-dotted curve: Born contributions; dashed curve: Δ(direct + exchange) contribution; solid curve: total (Born + Δ) contribution. The data are from Ref. 12 (triangles) and Ref. 13 (circles).

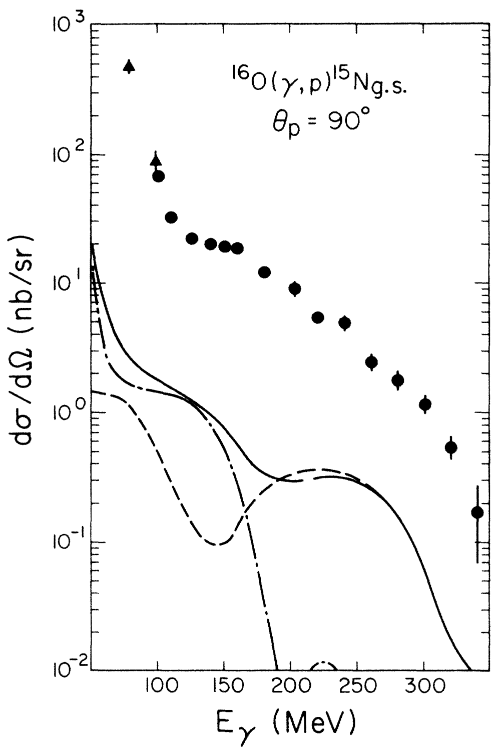

FIG. 3. Same as in Fig. 2, except that θ_p(lab)=90°.

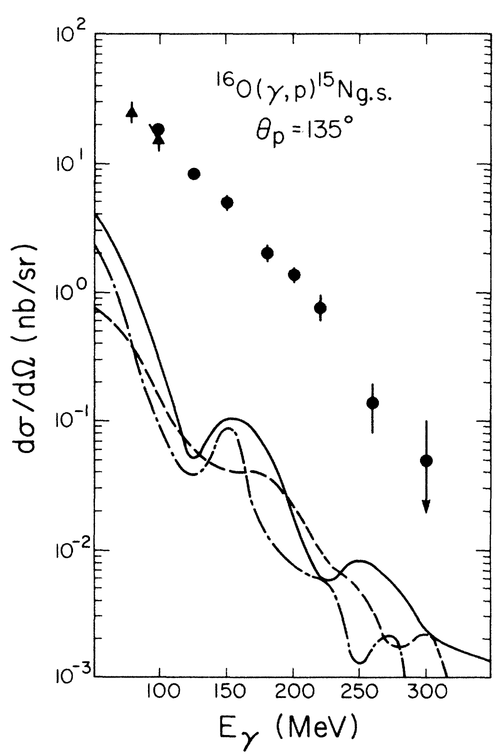

FIG. 4. Same as in Fig. 2, except that θ_p(lab)=135°.

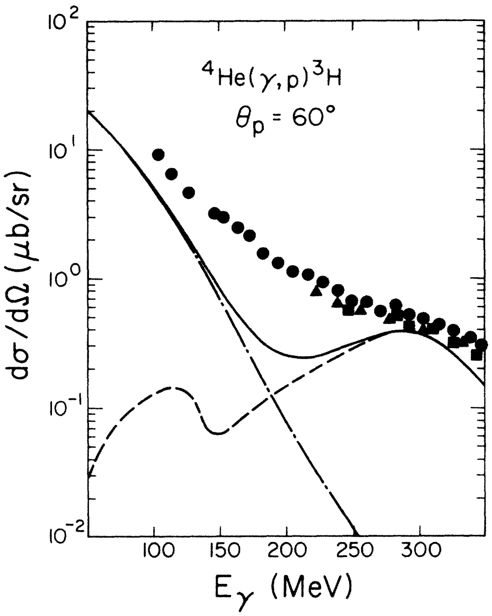

FIG. 5. Energy distributions for ^{4}He(γ,p)^{3}H at θ_p(c.m.)$=60°$. Dash-dotted curve: Born contribution; dashed curve: Δ(direct + exchange) contribution; solid curve: total (Born + Δ) contribution. The data are from Ref. 14 (circles), Ref. 15 (triangles), and Ref. 16 (squares).

butions from the Born term [Fig. 1(a)], and the two rescattering terms [Figs. 1(b) and (c)], as well as their coherent sum. The calculated differential cross section falls below the data in the entire energy range of 50—350 MeV. For an incident proton angle $\theta_p=45°$, the theoretical prediction is about a factor of 2 smaller than the data at the resonance energy of $E_\gamma \simeq 300$ MeV. The discrepancy increases as one goes to larger proton angles. Our calculated differential cross sections are far below the data at lower energies.

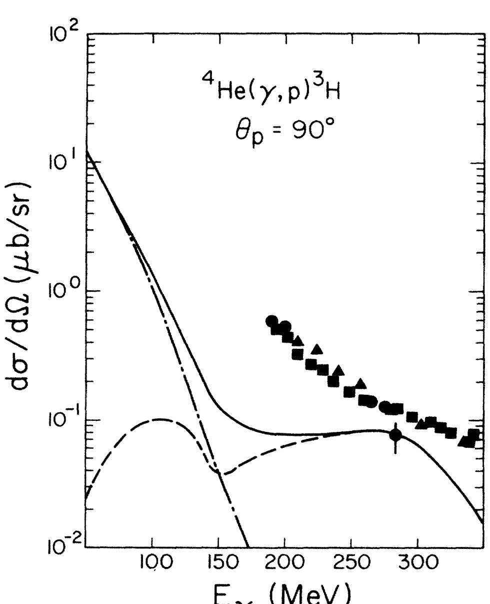

FIG. 6. Same as in Fig. 5, except that θ_p(c.m.)$=90°$.

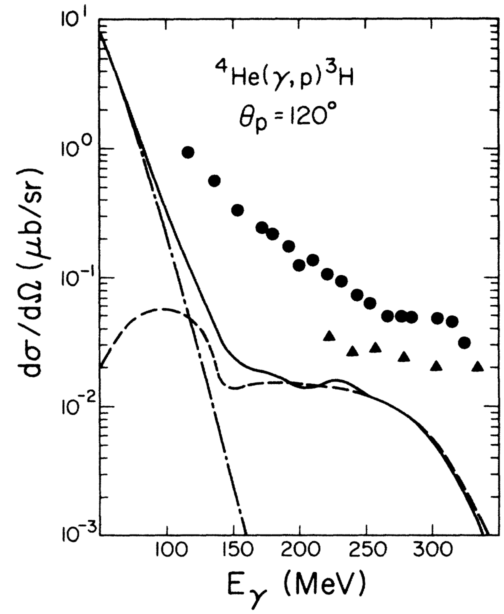

FIG. 7. Same as in Fig. 5, except that θ_p(c.m.)$=120°$.

Calculated energy distributions for ^{4}He(γ,p)^{3}H for $\theta_p=60°$, 90°, and 120° are shown in Figs. 5—7. Here the agreement with the data is qualitative in the resonance energy region, for $\theta_p=60°$ and $\theta_p=90°$. However, for $\theta_p=120°$, the theoretical curve is substantially below the data. Again, at lower energies, our predicted differential cross sections are too low. We also note that, at $\theta_p=120°$, the data of Refs. 14 and 15 are in disagreement with each other. This is discussed further in Ref. 14.

As mentioned in Sec. II, we have left out the convective one-body current, the inclusion of which would substantially enhance the theoretical predictions at low energies, and hence improve the agreement with experiment. Nevertheless, one can already draw certain conclusions from a comparison between the data and our calculations as they stand. Specifically, at photon energies of a few hundred MeV and proton angles $\theta_p \geq 45°$:

(1) the one-body current is insufficient for explaining the data at all angles;

(2) the Δ current is not insignificant; but

(3) the Δ current falls short of resolving the discrepancy between the one-body current and the data.

B. Comparison with other calculations

Our one-nucleon current contributions can be compared to the results of Londergan and Nixon,[2] and of Boffi *et al.*,[19] who have investigated sensitivities to the choices of bound nucleon wave functions and final-state proton-nucleus interactions. In particular, we observe that our one-nucleon results are substantially lower than those of Ref. 19 at photon energies around 100 MeV. While we do not understand this large difference completely, there are several possible explanations for it. First, since we only include the one-nucleon magnetic current, the one-nucleon cross section is missing the one-nucleon convective electric current, which dominates by factors of 2—4 in the lower region, as noted in Ref. 2. Second, at $E_\gamma=100$ MeV and $\theta_p \geq 45°$, the calculation becomes quite sensitive to the choice of bound nucleon wave functions, as is also noted in Ref. 2. This can be seen particularly in Fig. 4, where the calculated (unphysical) structure at $E_\gamma=150$ MeV is a consequence of evaluating single-nucleon wave functions at momenta $q \gtrsim 375$ MeV/c. (Note that, at these values of q, even the Δ contribution is sensitive to the high-momentum tails of the wave functions.) Finally, as noted in Ref. 19, the one-nucleon calculation is also quite sensitive to the choice of proton-nucleus optical potential. In particular, these authors note a strong dependence upon the imaginary part of the spin-orbit potential, which is surface peaked. Our optical potential, which is taken from Ref. 10, does not contain a surface-peaked *central* term of the sort used in Ref. 19, and this may affect the cross section magnitude even further.

The primary purpose of this paper is to compare our Δ contributions to the other results presently available. Our approach is closest to that of Londergan and Nixon (LN),[2] differing primarily in the treatment of Δ nonlocality and proton distortions. The results of our calculations agree with theirs on the point that the one-body current is negligible in the Δ-resonance energy region, at least for $\theta_p \geq 45°$. However, the size of the intermediate-Δ contribution we find is smaller than theirs, sometimes by an order of magnitude.

In LN, proton distortions were approximated by employing an effective complex proton momentum. For the one-nucleon term, these authors stated that this approximation agreed well with an optical potential calculation. However, the Fourier transform of the ^{16}O wave function has zeros (or sharp minima) at certain values of momentum.[2] These zeros will appear in (γ,p) differential cross sections in a PWIA calculation of the ^{16}O(γ,p)$N_{g.s.}$ reaction. We find that the wave function zeros are filled in by our proper inclusion of the proton distortions, so that only shallow minima or shoulders remain. In the case of LN, because proton distortions were introduced via an effective complex proton momentum, sharp minima still appear in their calculated differential cross sections.

For the case of the Δ term, no direct check was made by LN regarding the use of effective complex momenta in lieu of an optical potential. We have therefore examined explicitly the effects of proton distortions in our intermediate-Δ terms. Figures 8–10 display angular distributions for $E_\gamma=50$, 200, and 350 MeV [as well as data at 200 MeV (Ref. 20)], with and without a proton optical potential. The primary effect at the higher energies is an overall reduction of the differential cross section suggesting the appropriateness of the use of a complex effective proton momentum. At 50 MeV, the distortions have an angle-dependent effect, and the use of an effective momentum may not be reasonable, as anticipated by LN. Of course, at 50 MeV, the Δ contribution alone is certainly not dominant, and must be added coherently to the other one- and two-nucleon terms with distortions included.

For the ^{16}O wave function, LN used an oscillator expansion of the Negele density-dependent Hartree-Fock wave function,[21] while we have used the Elton-Swift wave function.[9] Both the Negele and Elton-Swift wave func-

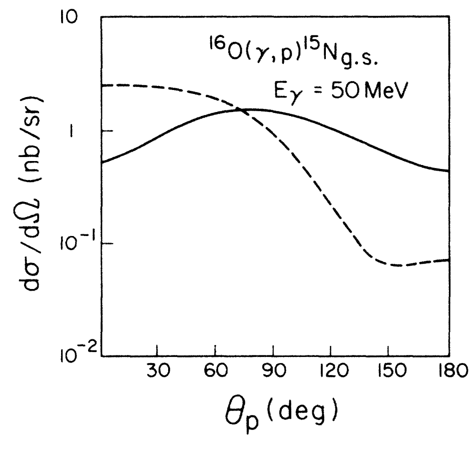

FIG. 8. Angular distributions for ^{16}O(γ,p)$^{15}N_{g.s.}$ at $E_\gamma=50$ MeV (Δ contribution only). Solid curve: with proton distortions; dashed curve: without proton distortions.

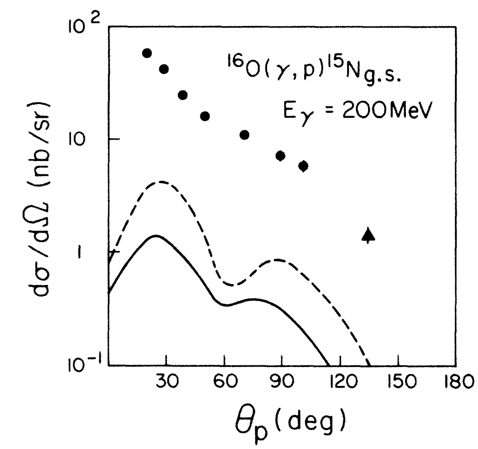

FIG. 9. Same as in Fig. 8, except that $E_\gamma=200$ MeV, and the data are from Ref. 13 (triangle) and Ref. 19 (circles).

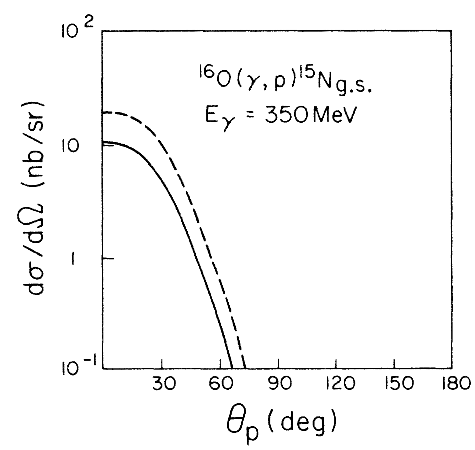

FIG. 10. Same as in Fig. 8, except that $E_\gamma=350$ MeV.

tions fit binding energies and elastic electron scattering data at 420 MeV. However, the high momentum contents of the two functions are quite different, since they are not constrained by the experimental data. Also it is shown in LN that, compared with the original Negele wave function, the oscillator expansion does not contain sufficient large momentum components for $q \gtrsim 700$ MeV/c.

A shortcoming of our calculation vs LN is the neglect of Δ nonlocality, which arises both from recoil and medium effects. However, the quantitative implications of this approximation may not be so large. The LN paper included (γ,p) calculations in which the Δ width was artificially reduced from 115 to 50 MeV: the difference is very small. We would expect recoil, Fermi motion, and medium effects to be connected to the sensitivity to the width of the Δ, and are therefore encouraged to see that these effects may be small for the (γ,p) process. Nevertheless, it is quite possible that, apart from varying the width of the Δ, other nonlocal effects in the LN calculation could explain why our cross sections have a pronounced Δ peak and theirs do not. Another difference between the two approaches comes from our choice of k_π^0 in the exchange term in Eq. (2.6), which gives rise to different angular and energy dependence of the cross sections than when $k_\pi^0=0$, as in LN.

For the specific case of ^{4}He(γ,p)^{3}H (Figs. 5—7), our calculations of the Δ contribution exhibit anomalous behavior at the pion production threshold not seen in the data. This behavior stems directly from our factorization of the momentum integral. A more careful evaluation of these integrals would amount to Fermi averaging this threshold, and its effect would be much less pronounced, though the choice of k_π^0 might still affect the results.

Finally, we have also neglected the Δ contribution with crossed pion and photon lines which was included in LN. This contribution corresponds to the process in which pion emission precedes photon absorption. However, because of its large energy denominator, this term should be small compared with the terms with uncrossed lines [Figs. 1(b) and (c)].

For the case of ^{16}O(γ,p)^{15}N$_{\rm g.s.}$, the one-nucleon-current contribution calculated by Gari and Hebach (GH)[3] is much larger than those obtained by us and LN in the region of $E_\gamma \gtrsim 250$ MeV. This could be due to the fact that the ^{16}O shell-model Woods-Saxon potential used by GH has parameters different from those of Elton and Swift.[9] Moreover, they have neglected the spin-orbit potential.

We find a Δ contribution comparable to that obtained by GH. Although it is difficult to make a direct comparison of the two approaches, there are some specific observations which can be made. First, their γNΔ coupling was obtained by fitting to the exchange current effect in np→dγ at thermal energies, and may well differ from ours. Second, GH claimed that a major reduction in their intermediate-Δ term comes from an extra retardation factor [see Eq. (3.29) of Ref. 3]. The only retardation effect which we obtain stems from the choice of fourth component of the virtual pion momentum; we are unable to reproduce their retardation factor which, according to them, is a crucial element in the size of their Δ contribution. It was also claimed by GH that nucleon-nucleon correlations are important in a proper evaluation of an intermediate-Δ diagram. Like LN, we have neglected the correlation effects. However, the πNN and πNΔ vertex form factors used by us and LN are "soft," so that the importance of the short-range correlations are expected to be greatly reduced.[22]

The calculations of Finjord[17] for ^{4}He employed Chew-Goldberger-Low-Nambu (CGLN) γN→πN amplitudes[22] for the two-nucleon-current term. The important nonresonant contributions to photoproduction of charged pions described in Sec. III A are thus included. However, there remain questions of gauge invariance when the CGLN amplitudes are taken off shell, and of overcounting the crossed photon-pion contribution to pion photoproduction, which is normally considered to be included in the single nucleon wave function contributions to the electromagnetic current.

IV. SUMMARY AND ANALYSIS

In this paper, we have investigated the role of the Δ(1232) in the (γ,p) reaction at intermediate energies. The calculation follows an approach which was recently developed for the (p,π) reaction.[7] Both one-nucleon and two-nucleon intermediate-Δ contributions have been included. In order to focus on the Δ contribution, we have included only the magnetic portions of the nuclear current, which are individually gauge invariant. Some momentum-dependent effects, including the Δ nonlocality, have been neglected, and the calculations have been performed in coordinate space. It is then straightforward to describe the proton final-state interactions with an optical potential.

Although there is much room for improvement, the following features stand out in our calculations. Our predicted differential cross sections are generally below the data. We find that the one-nucleon contribution, while dominant at $E_\gamma < 100$ MeV, is negligible at proton angles greater than 45° for $E_\gamma \simeq 300$ MeV. In the Δ-resonance energy region, the intermediate-Δ contributions are comparable to the data for the ^{16}O(γ,p)^{15}N$_{\rm g.s.}$ reaction at $\theta_{\rm p}=45°$, and for the ^{4}He(γ,p)^{3}H reaction at $\theta_{\rm p}=60°$ and 90°. For the other cases investigated, the Δ contribution is about a factor of 10 below the data. Our findings thus imply that, depending on the combined conditions of target nucleus, energy and angle, the intermediate-Δ contribution can be an important contribution to the (γ,p) reaction.

Experiments do show a bump in the energy distribution at $E_\gamma \simeq 300$ MeV for the γd→pn reaction, but not for the (γ,p) reaction in heavier nuclei. [Earlier data indicating a resonant structure in the ^{4}He(γ,p) cross section have not been reproduced in more recent experiments; see Ref. 17.] One would like to understand theoretically why this is the case. First, the deuteron case is unique because the final state consists of free nucleons. In heavier nuclei, the virtual pion emitted at the πNΔ vertex must be absorbed by a bound nucleon, so that, in the calculation, there are two extra single particle wave functions which tend to damp out and broaden the resonance. Second, the deuteron is unique because the initial pair of nucleons has *only* isospin zero. The $T=\frac{3}{2}$ and $T=\frac{1}{2}$ amplitudes for physical pion photoproduction from a single nucleon are comparable at intermediate energies.[5] Indeed, the Δ dominates only in the case of π^0 photoproduction because the nonresonant contribution is suppressed due to the Kroll-Ruderman theorem.[18] Since all pion charge states contribute to photon absorption by two nucleons, the nonresonant $T=\frac{1}{2}$ $\gamma N \rightarrow \pi N$ amplitude is involved as well as the $T=\frac{3}{2}$ amplitude. For two nucleons coupled to isospin zero, such as in the deuteron, the $T=\frac{3}{2}$ amplitude has a large weighting factor compared to the $T=\frac{1}{2}$ amplitude. This is not the case for two nucleons coupled to isospin one. Thus, for all target nuclei except the deuteron, the signature of the Δ contribution is diluted relative to the nonresonant background.

Our calculated Δ contributions are smaller than those obtained in LN,[2] but appear comparable to those of GH (Ref. 3) at angles where a comparison can be made. From the comparison of our results with the data, we can also draw the conclusion that other exchange currents of nonresonant nature must be important. This is in agreement with the findings of GH. Regarding the one-nucleon contribution, GH found a much larger effect than those obtained by us and LN, for $E_\gamma \gtrsim 250$ MeV, in the $^{16}O(\gamma,p)^{15}N_{g.s.}$ reaction.

The primary missing ingredient in our calculation is the contribution from the *electric* portion of the nuclear current, including both one-nucleon and (nonresonant) two-nucleon terms. While it is clear from the data that these contributions, particularly the two-nucleon term, are important, if not dominant, at intermediate energies, it is also difficult to obtain a model independent, gauge-invariant result for these terms. We mention briefly some of the difficulties involved.

First, as mentioned earlier, the approach of Finjord[17] which employs CGLN pion photoproduction amplitudes[23] is gauge invariant when the nucleons are on the mass shell, but not off the mass shell. Furthermore, the crossed photon-pion term should be subtracted out, since it is already included in the proton single-particle wave function.

Second, the approach of GH ensures gauge invariance of the two-nucleon contribution, at least in the zero-photon-energy limit, by replacing the nonresonant part of the two-nucleon current with a commutator between an effective two-nucleon potential and the long-wavelength limit of the electric multipole operator. However, it is not obvious that this procedure is justified in an energy region where neglected terms of higher order in $(E_\gamma R)$ are not small (R is the nuclear radius). The issue of gauge invariance has been examined for the case of deuteron photodisintegration for $E_\gamma < 100$ MeV by Arenhövel[24] and by Hwang and Miller.[25] In this case, the use of Siegert's theorem[26] gives stable results because the charge density is actually calculated from the two-nucleon potential, and because the commutator $[V,\rho]$ is approximately independent of potentials adjusted to fit NN phase shifts. However, for heavier nuclei, the two-nucleon potential is not directly related to the charge density. Indeed, it has been shown[27] that a calculation which uses Siegert's theorem is uncertain by an amount related to the lack of orthogonality between bound state and continuum (scattering) state single-particle wave functions as they are normally calculated. On the other hand, as was done by GH, one could guarantee orthogonality by calculating the continuum state wave function in the shell model potential. However, it is also well known that this procedure would give an unrealistic description of the proton final state distortions in the energy region of this work. We conclude, therefore, that these calculations of the nonresonant two-nucleon current contribution to the (γ,p) reaction at intermediate energies, while achieving reasonable fits to some of the data, contain steps which can only be checked by a calculation which goes well beyond the scope of any paper yet written on this reaction.

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[1]See, e.g., Refs. 2 and 3, and references quoted therein.
[2]J. T. Londergan and G. D. Nixon, Phys. Rev. C **19**, 998 (1979).
[3]M. Gari and H. Hebach, Phys. Rep. **72**, 1 (1981).
[4]J. T. Londergan, in *Pion Production and Absorption in Nuclei 1981—(Indiana University Cyclotron Facility)*, Proceedings of the Conference on Pion Production and Absorption in Nuclei, AIP Conf. Proc. No. 79, edited by R. D. Bent (AIP, New York, 1982).
[5]See, e.g., I. Blomqvist and J. M. Laget, Nucl. Phys. **A280**, 405 (1977).
[6]J. Arends, H. J. Gassen, A. Hegerath, B. Mecking, G. Nöldeke, P. Prenzel, T. Reichelt, A. Voswinkel, and W. W. Sapp, Nucl. Phys. **A412**, 509 (1984).
[7]B. D. Keister and L. S. Kisslinger, Nucl. Phys. **A412**, 301 (1984).
[8]D. M. Brink and G. R. Satchler, *Angular Momentum*, 2nd ed. (Oxford University, Oxford, 1968).
[9]L. R. B. Elton and A. Swift, Nucl. Phys. **A94**, 52 (1967).
[10]A. Nadasen, P. Schwandt, P. P. Singh, W. W. Jacobs, A. D. Bacher, P. T. Debevec, M. D. Kaitchuck, and J. T. Meek, Phys. Rev. C **23**, 1023 (1981).
[11]R. F. Frosch, J. S. McCarthy, R. E. Rand, and M. R. Yearian, Phys. Rev. **160**, 874 (1967).
[12]D. J. S. Findlay and R. O. Owens, Nucl. Phys. **A279**, 385 (1977).
[13]M. J. Leitch, J. L. Matthews, W. W. Sapp, C. P. Sargent, S. A.

Wood, D. J. S. Findlay, R. O. Owens, and B. L. Roberts, Phys. Rev. C **31**, 1633 (1985).

[14]R. A. Schumacher, J. L. Matthews, W. W. Sapp, R. S. Turley, G. S. Adams, and R. O. Owens, Phys. Rev. C (in press).

[15]J. Arends, J. Eyink, A. Hegerath, H. Hartmann, B. Mecking, G. Nöldeke, and H. Rost, Nucl. Phys. **A322**, 253 (1979).

[16]P. E. Argan, G. Audit, N. de Botton, J. L. Faure, J. M. Laget, J. Martin, C. G. Schuhl, and G. Tamas, Nucl. Phys. **A237**, 447 (1975).

[17]J. Finjord, Nucl. Phys. **A274**, 495 (1976).

[18]N. M. Kroll and M. A. Ruderman, Phys. Rev. **93**, 233 (1954).

[19]S. Boffi, C. Giusti, and F. D. Pacati, Nucl. Phys. **A359**, 91 (1981).

[20]R. S. Turley, E. R. Kinney, J. L. Matthews, W. W. Sapp, E. J. Scheidker, R. A. Schumacher, S. A. Wood, G. S. Adams, and R. O. Owens, Phys. Lett. **157B**, 19 (1985).

[21]J. W. Negele, Phys. Rev. C **1**, 1260 (1970).

[22]See, e.g., J. M. Eisenberg, J. Hüffner, and E. J. Moniz, Phys. Lett. **47B**, 381 (1973).

[23]G. F. Chew, M. L. Golderger, F. E. Low, and Y. Nambu, Phys. Rev. **106**, 1345 (1957).

[24]H. Arenhövel, Z. Phys. A **302**, 25 (1981).

[25]W. Y. P. Hwang and G. A. Miller, Phys. Rev. C **22**, 968 (1980).

[26]A. J. F. Siegert, Phys. Rev. **52**, 787 (1937).

[27]J. M. Lafferty and S. R. Cotanch, Nucl. Phys. **A373**, 363 (1982).