Photopion production from polarized nuclear targets

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Recently nonlocal distorted wave impulse approximation (DWIA) calculations for $A(\gamma,\pi N)B$ photopion production reactions have been reported. In this paper we examine the factorized DWIA formulation to help clarify some of the underlying physics content inherent in such measurements. We also extend the formulation to the case of polarized targets and to the measurement of the emitted nucleon polarization. We show that with reasonable approximations one can formally write the three-body reaction spin observables in terms of products of the two-body $\gamma+N\rightarrow\pi+N$ spin observables, the polarization of the struck nucleon which depends on nuclear structure, and an effective polarization created by the distortions. Calculations are presented for two nuclear targets for which adequate polarization has been obtained to utilize them in experiments; viz., ^{13}C and ^{19}F. These calculations provide typical predictions which might be compared to experimental data to search for evidence of medium modifications of the pion production amplitudes.

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I. INTRODUCTION

In a recent paper [1] Li, Wright, and Bennhold (LWB) presented a theoretical study of exclusive photopion production on nuclei using a distorted wave impulse approximation (DWIA) formalism. Their calculations were carried out in a momentum representation, requiring integration over six dimensions, but permitting them to include, without approximation, nonlocal effects arising from the production operator. Calculations were presented for $A(\gamma,\pi p)B$ on p-shell nuclei and compared to available experimental data. LWB found reasonable agreement between the calculations and data, and proposed new experiments emphasizing the measurement of photon asymmetry, which might lead to an improved understanding of the photopion production process and Δ propagation. The photon asymmetry measurements appear to have least sensitivity to the various uncertainties in the calculations.

In this paper we wish to examine in somewhat more detail the sensitivity of the photopion production reaction and to expand upon the work of LWB by calculating additional spin observables which might also lead to an improved understanding of the process. In particular we consider polarized nuclear targets which have come into use over the past few years, as well as measurements of the emitted nucleon polarization. To do so we have carried out factorized calculations in coordinate space, thereby evaluating the elementary photopion production amplitude at the asymptotic momenta. Although this is an additional approximation, we believe that our approach enables us to highlight more clearly some of the physics which might be accessible in such three-body studies. Furthermore the studies of LWB show that a local approximation is rather good, particularly for regions of phase space near those populated in the elementary photoproduction process [quasifree angles in which the residual nucleus in the $A(\gamma,\pi p)B$ reaction is left with zero or low recoil momentum]. Thus, although a proper treatment of the nonlocality will be needed for detailed comparison with experimental data, the formalism and calculations presented here provide the experimentalist with excellent guidance in designing experiments to elucidate the underlying physics.

In Sec. II we very briefly present the DWIA formalism, and in Sec. III examine the qualitative features of the reaction by considering a polarized nuclear target, and a coplanar geometry. In Sec. IV we present specific calculations for ^{13}C$(\gamma,\pi^-p)^{12}$C and ^{19}F$(\gamma,\pi^+n)^{18}$O. The conclusions and summary are presented in Sec. V.

II. DWIA FORMALISM

In several previous papers [2–4] we have discussed various aspects of distorted wave effects in pion-induced reactions, and specifically have considered polarized nuclear targets for both nucleon-knockout reactions $\vec{A}(\pi,\pi p)B$ and two-nucleon absorption $\vec{A}(\pi^+,\ 2p)B$. The present work very closely follows that of Ref. [2]. As in that case, since we are considering spin-dependent effects, particularly the outgoing nucleon polarization, we have included the possibility of spin-dependent distortion for the emitted nucleon.

We consider the reaction $\vec{A}(\gamma,\pi N)B$ where N represents a nucleon and the target $A=B+N$. We note that the incoming photon has spin one with two possible polarization states. Thus, in DWIA the cross section must have nearly the same form as that for $\vec{A}(\pi,\pi N)B$, the major change being that the incoming photon has two possible spin projections. Assuming a linearly polarized photon beam, the two polarization states enter as an incoherent sum in the expression for the unpolarized cross section. Therefore, generalizing Eqs. (2) and (3) of Ref. [2] we can write the cross section for photopion production $(\gamma,\pi N)$ by a photon with polarization ϵ on a nucleus A with total angular momentum J_A (projection M_A), leading to a final nucleus with total angular momentum J_B (projection M_B), and a nucleon with spin projection ρ_d'' as

$$\sigma_{BA}(M_A,\rho_d'',\epsilon)=\frac{2\pi}{\hbar c}\omega_B\sqrt{A}\sum_{M_B}\left|\sum_{\lambda\sigma m\tau\sigma_d''}^{lj}\mathscr{T}_{AB}(nlsjt)\;(J_BM_Bjm|J_AM_A)(T_BN_Bt\tau|T_AN_A)(l\lambda s\sigma|jm)T^{l\lambda}_{\sigma_d''\rho_d''}\langle\sigma_d'';\mathbf{k}_\pi'|t|\sigma;\epsilon\mathbf{k}_\gamma\rangle\right|^2, \tag{1}$$

where ω_B is the energy density of final states. For the struck nucleon, N, the quantum number j (projection m), represents the total angular momentum which is composed of angular momentum l (λ) and spin $s=1/2$ (σ). The corresponding isospin quantum numbers for N are t (τ), and for nucleus i are T_i (N_i). The quantities in brackets are vector coupling coefficients. The spectroscopic amplitude $\mathscr{T}_{AB}$ is related to the conventional single nucleon spectroscopic factor through $S=A\times\mathscr{T}_{AB}^2$, where A is the number of target nucleons. The matrix element $\langle\sigma_d'';\mathbf{k}_\pi'|t|\sigma;\epsilon\mathbf{k}_\gamma\rangle$ is the amplitude for photopion production from the free nucleon, and for these calculations is taken from the work of Blomqvist and Laget [5]. The effects of distortion are contained in

$$T^{l\lambda}_{\sigma_d''\rho_d''}=\int\chi_c^{(-)*}\chi^{(-)*}_{\sigma_d''\rho_d''}\phi_{l\lambda}\chi^{(+)}d^3r, \tag{2}$$

where the $\chi^{(-)}$'s represent distorted waves for the emitted pion and nucleon, and $\chi^{(+)}$ represents the incoming photon plane wave. The spatial part of the struck nucleon wave function is $\phi_{l\lambda}$.

It is the above equations which have been coded to carry out our DWIA calculations. From these calculations we may compute a variety of spin observables, in addition to the unpolarized cross section. Defining a sum of partial cross sections as

$$S^{3b}=\sum_{M_A,\rho_d'',\epsilon}\sigma_{BA}(M_A,\rho_d'',\epsilon), \tag{3}$$

we can write the photon asymmetry as

$$\Sigma^{3b}=\sum_{M_A,\rho_d''}[\sigma_{BA}(M_A,\rho_d'',\epsilon=\perp)-\sigma_{BA}(M_A,\rho_d'',\epsilon=\|)]/S^{3b}. \tag{4}$$

Considering only polarization components along the axis of quantization we can write the polarization of the emitted nucleon as

$$P_N^{3b}=\sum_{M_A,\epsilon}[\sigma_{BA}(M_A,\rho_d''=+\tfrac{1}{2},\epsilon)-\sigma_{BA}(M_A,\rho_d''=-\tfrac{1}{2},\epsilon)]/S^{3b} \tag{5}$$

and the target analyzing powers are a combination of cross sections for different substates M_A, summed over ρ_d'' and ϵ. For $J_A=1/2$ we have only a vector analyzing power

$$A_T^{3b}=\sum_{\rho_d'',\epsilon}[\sigma_{BA}(M_A=+\tfrac{1}{2},\rho_d'',\epsilon)-\sigma_{BA}(M_A=-\tfrac{1}{2},\rho_d'',\epsilon)]/S^{3b}. \tag{6}$$

For larger values of J_A tensor analyzing powers can be calculated. One may also generate various spin correlation observables, but those are not considered in the present paper.

III. QUALITATIVE FEATURES OF THE DWIA CALCULATIONS

The above equations permit the calculation of the various spin observables for photopion production on nuclei. However, as in our previous papers [2–4], it is instructive to examine the DWIA formalism with certain simplifying approximations. These approximations bring to light much of the underlying physics of the reaction, but are not so drastic as to invalidate the conclusions for the full calculation. First, we assume that only one bound nucleon orbital (i.e., single values of l and j) contributes to the transition. We will also ignore the spin dependence of the emitted nucleon distorted wave so that $T^{l\lambda}_{\sigma_d''\rho_d''}$ becomes diagonal in the emitted nucleon spin, σ_d'',ρ_d''. This approximation is likely to have greatest impact on the calculation of the outgoing nucleon polarization and is examined in the calculations which follow. Next, we restrict our considerations to a coplanar geometry for the reaction, and assume that the target nucleus A is polarized along a direction normal to the scattering plane, certainly a common experimental arrangement. With this restriction in geometry it is natural to choose the axis of quantization normal to the reaction plane. With these approximations Eqs. (1) and (2) can be rewritten as

$$\sigma_{BA}(M_A,\rho_d'',\epsilon)=\frac{2\pi}{\hbar c}\omega_B\sqrt{A}\sum_{M_B}\left|\sum_{\lambda\sigma m\tau}\mathscr{T}_{AB}(nlsjt)\;(J_BM_Bjm|J_AM_A)(T_BN_Bt\tau|T_AN_A)(l\lambda s\sigma|jm)T^{l\lambda}\langle\rho_d'';\mathbf{k}_\pi'|t|\sigma;\epsilon\mathbf{k}_\gamma\rangle\right|^2 \tag{7}$$

and

$$T^{l\lambda}=\int \chi_c^{(-)*}\chi_d^{(-)*}u_{lj}(r)Y_{l\lambda}(\theta,\phi)\chi^{(+)}d^3r, \tag{8}$$

where we have written the bound nucleon wave function explicitly in terms of the radial wave function and the spherical harmonic.

As in Refs. [2–4], with the coplanar geometry and the choice of axis of quantization, we may now use the symmetry of the spherical harmonics to show that the amplitude $T^{l\lambda}$ is nonzero only if $l+\lambda$ is even. Since the sum over M_B is incoherent and M_A is fixed, the sum over m is incoherent. As a result, the summations over λ and σ become incoherent since, for each value of m, $\lambda=m\pm 1/2$ and terms of the form $T^{l\lambda}T^{l\lambda'*}$ must vanish if $\lambda\neq\lambda'$. Thus we arrive at essentially the same expression as that given in Ref. [2] for $A(\pi,\pi N)B$ with the addition of the polarization of the incoming photon. Omitting spin-independent factors we obtain

$$\sigma_{BA}(M_A,\epsilon,\rho_d'')\sim\sum_{M_B\lambda}(J_BM_Bjm|J_AM_A)^2(l\lambda s\sigma|jm)^2|T^{l\lambda}|^2 \times|\langle\rho_d'';\mathbf{k}_\pi'|t|\sigma;\epsilon\mathbf{k}_\gamma\rangle|^2. \tag{9}$$

In the following two subsections we use the above expression to examine two special cases that clarify the physics. In particular we consider the removal of s-state ($l=0$) and p-state ($l=1$) nucleons from a target nucleus with $J_A=1/2$. To relate the spin observables in the three body nuclear reactions to two-body photopion production, we first need to define some spin observables for the two-body case $\vec{\gamma}+\vec{N}\rightarrow\pi+\vec{N}$. Rather than using the conventional notation used for photopion production [6], we prefer the more descriptive notation of Ref. [7] for nucleon-nucleon scattering, in which observables for a reaction $a+b\rightarrow c+d$ are denoted by X_{cdab} with the subscripts used to denote the polarization measurement. (The conventional notation is also included in the following definitions.) We define the normal to the scattering plane as $\vec{n}=\vec{k_a}\times\vec{k_c}$ and use partial cross sections σ_{cdab} in which the subscripts specify the spin or polarization coordinate and a zero is used to imply summation over the corresponding observable. Then the unpolarized two-body cross section is

$$\sigma^{2b}=\sigma_{0000}/4, \tag{10}$$

the two-body photon asymmetry is

$$\Sigma=A^{2b}_{00\epsilon 0}=(\sigma_{00\perp 0}-\sigma_{00\|0})/\sigma_{0000}, \tag{11}$$

the two-body target analyzing power is

$$T=A^{2b}_{000n}=(\sigma_{000+}-\sigma_{000-})/\sigma_{0000}, \tag{12}$$

and the two-body spin correlation coefficient is

$$C^{2b}_{00\epsilon n}=(\sigma_{00\perp +}-\sigma_{00\perp -}-\sigma_{00\| +}+\sigma_{00\| -})/\sigma_{0000}. \tag{13}$$

For our discussion of emitted nucleon polarization, we shall need the two-body nucleon polarization,

$$P=P^{2b}_{0n00}=(\sigma_{0+00}-\sigma_{0-000})/\sigma_{0000}, \tag{14}$$

and the polarization transfer,

$$D^{2b}_{0n0n}=(\sigma_{0+0+}-\sigma_{0+0-}-\sigma_{0-0+}+\sigma_{0-0-})/\sigma_{0000}. \tag{15}$$

A. Transitions with $l=0$

Using Eq. (9) we now examine specific transitions. The simplest and most transparent case is that of s-wave nucleon removal, l=0, j=1/2. This is a case of some interest for the polarized targets mentioned. The ground state to ground state transition for $^{19}\mathrm{F}(\gamma,\pi^+n)^{18}\mathrm{O}$ results from the removal of a $2s_{1/2}$ proton. For ^{13}C the removal of a deeply bound $1s_{1/2}$ nucleon leads to 0^- and 1^- states. For $l=0$ the cross section in Eq. (9) may be written

$$\sigma_{BA}(M_A,\epsilon,\rho_d'')\sim\sum_{M_B}(J_BM_B\tfrac{1}{2}\sigma|J_AM_A)^2|T^{00}|^2 \times|\langle\rho_d'';\mathbf{k}_\pi'|t|\sigma;\epsilon\mathbf{k}_\gamma\rangle|^2. \tag{16}$$

Thus the effects of distortions factor out of the spin dependence, leading to a change (primarily a reduction) in the overall cross section with no change in the spin-observables for an unpolarized target (in which we can also sum over M_A). In the case of $l=0$ nucleon removal from a polarized target, the only modification is that the polarized nuclear target simply provides a polarized *nucleon* target with an effective polarization P_s determined by the nuclear structure through the Clebsh-Gordon coefficient. Specifically, for $l=0$ we can write

$$P_s^{M_A}=\frac{\Sigma_{M_B}(J_BM_B\tfrac{1}{2}+\tfrac{1}{2}|J_AM_A)^2-\Sigma_{M_B}(J_BM_B\tfrac{1}{2}-\tfrac{1}{2}|J_AM_A)^2}{\Sigma_{M_B}(J_BM_B\tfrac{1}{2}+\tfrac{1}{2}|J_AM_A)^2+\Sigma_{M_B}(J_BM_B\tfrac{1}{2}-\tfrac{1}{2}|J_AM_A)^2}. \tag{17}$$

As in Ref. [2] for $(\pi,\pi N)$, we can write the cross section for target substate M_A and incoming photon polarization ϵ as

$$\sigma(M_A,\epsilon=\pm)=\sigma_0^{3b}(1\pm A^{2b}_{00\epsilon 0}+A^{2b}_{000n}P_s^{M_A}\pm C^{2b}_{00\epsilon n}P_s^{M_A}), \tag{18}$$

where, for convenience, we have changed the notation so that $\epsilon=\pm$ is associated with photon polarization perpendicular ($\perp$) and parallel ($\|$) to the reaction plane, respectively.

Considering the special cases of ^{13}C and ^{19}F with $J_A=1/2$, we calculate the target vector analyzing powers and photon asymmetries in terms of the two-body spin observables. For the case of a spin-1/2 target the $P_s^{M_A}$ simply reverse in sign for $M_A=\pm 1/2$, so that we can write $P_s^{M_A}=\pm P_s$. Values of P_s for the various transitions are given in Table I. For the three body reaction on a spin-1/2 target, the photon asymmetry for $l=0$ is given by

TABLE I. Values of the effective polarization parameters for $l=0$ and $l=1$ transitions for target nuclei with $J_A=1/2$.

lj	$J_A \;\; J_B$	P_s	P_o	α
$s_{1/2}$	$\frac{1}{2} \rightarrow 0$	1	0	0
	$\frac{1}{2} \rightarrow 1$	$-\frac{1}{3}$	0	0
$p_{1/2}$	$\frac{1}{2} \rightarrow 0$	-1	1	-1
$p_{1/2}$	$\frac{1}{2} \rightarrow 1$	$\frac{1}{3}$	$-\frac{1}{3}$	-1
$p_{3/2}$	$\frac{1}{2} \rightarrow 1$	$\frac{2}{3}$	$\frac{5}{6}$	$\frac{1}{2}$
$p_{3/2}$	$\frac{1}{2} \rightarrow 2$	$-\frac{2}{5}$	$-\frac{1}{2}$	$\frac{1}{2}$

$$\Sigma^{3b}=A^{2b}_{00\epsilon 0} \tag{19}$$

and the target asymmetry by

$$A^{3b}_T=A^{2b}_{00n0}P_s. \tag{20}$$

Additionally, examining Eq. (16) we note that the emitted nucleon polarization defined by Eq. (5) is simply the two-body polarization

$$P^{3b}_N=P^{2b}_{0n00}. \tag{21}$$

Thus we see that for this $1/2\rightarrow 0$ transition the spin observables are simply related to the two-body spin observables, the only nuclear structure dependence being the factor P_s (involving only initial and final nuclear spins) contained in A^{3b}_T. Deviations due to various medium modifications, such as a spin-dependent Δ spreading potential, should be easily identified with the $l=0$ transition, and the measurement of all three spin observables would provide a stringent test of any reaction model.

The preceding discussion makes measurements of the spin observables for $l=0$ transitions on spin-1/2 targets prime candidates for studies of the modification of the photopion production mechanism in the nuclear medium — for example modifications due to the propagation of the Δ in the medium. As will be seen in the next section the $l=0$ case is to be preferred to the $l=1$ transitions for which distortion effects contribute significantly to the spin observables.

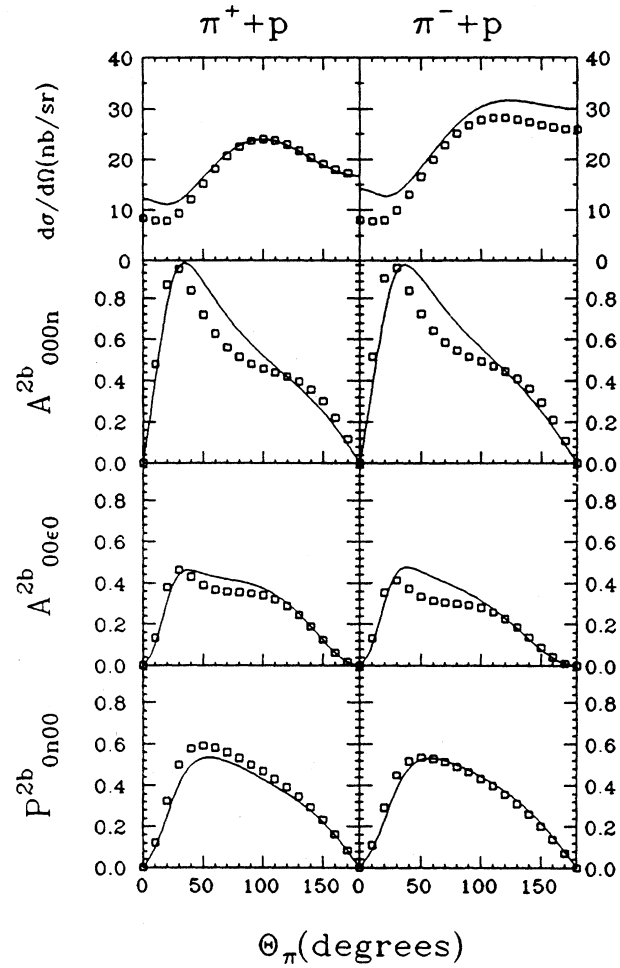

FIG. 1. Observables for charged pion photoproduction from nucleons at 300 MeV. The squares are calculated with the code SAID using the solution SM95 phase shifts [6]. The curves are from the work of Blomqvist and Laget [5] using the parameters for which the transition operator is unitary.

B. Transitions with $l=1$

We now turn to the somewhat more complicated case of $l=1$ nucleon removal. As noted above, for $l=1$ the distorted wave amplitude $T^{l\lambda}$ is nonzero only if $\lambda=\pm 1$. Again we follow $(\pi,\pi p)$ first introducing an effective polarization in orbital angular momentum induced by the distortion effects (Maris effect [8] or Newns' polarization [9]) as

$$P=\frac{|T^{11}|^2-|T^{1-1}|^2}{|T^{11}|^2+|T^{1-1}|^2}. \tag{22}$$

We also define various angular momentum projection probabilities for spin

$$C_\pm=\sum_{M_B m\lambda=\pm 1}(J_BM_Bjm|J_AM_A)^2(1\lambda\tfrac{1}{2}\pm\tfrac{1}{2}|jm)^2, \tag{23}$$

for orbital angular momentum

$$D_\pm=\sum_{M_B m\sigma}(J_BM_Bjm|J_AM_A)^2(1\pm 1\tfrac{1}{2}\sigma|jm)^2, \tag{24}$$

and for parallel versus antiparallel λ and σ

$$E_\pm=\sum_{M_B m}(J_BM_Bjm|J_AM_A)^2(1+1\tfrac{1}{2}\pm\tfrac{1}{2}|jm)^2 + (J_BM_Bjm|J_AM_A)^2(1-1\tfrac{1}{2}\mp\tfrac{1}{2}|jm)^2. \tag{25}$$

Note that $C_+ + C_- = D_+ + D_- = E_+ + E_-$. Using these probabilities we define the "spin polarization" $P^{M_A}_s$ as a generalization of the expression already given for $l=0$ transitions. Similarly, we can define an "orbital polarization" $P^{M_A}_o$, and a parameter α^{M_A} which converts the effective polarization in **l** to an effective polarization in the nucleon spin as follows:

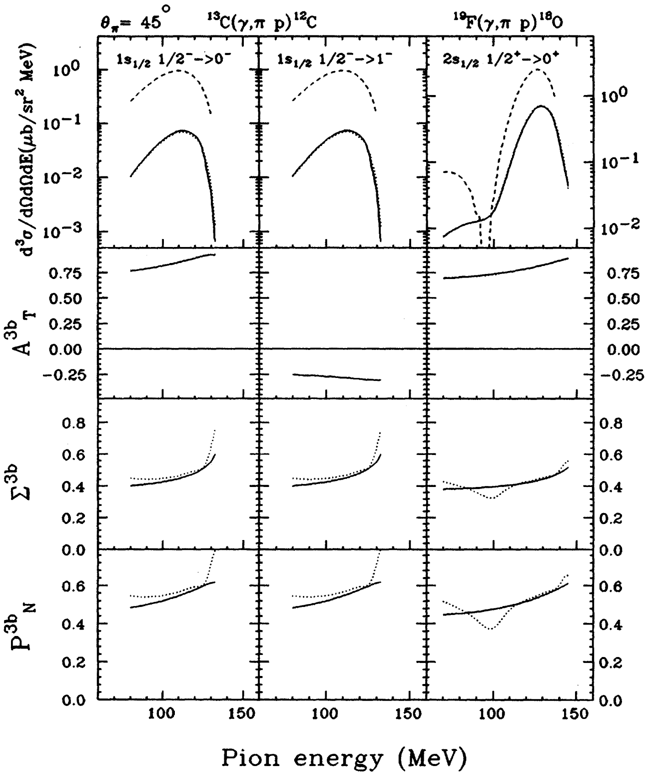

FIG. 2. DWIA energy sharing calculations of $l=0$ nucleon removal for $^{13}\vec{\mathrm{C}}(\gamma,\pi^-p)^{12}\mathrm{C}$ to 0^- and 1^- final states and $^{19}\vec{\mathrm{F}}(\gamma,\pi^+n)^{18}\mathrm{O}$ to the 0^+ ground state at an incident photon energy of 300 MeV and angles $\theta_\pi=45°$ and $\theta_N=-51°$. The top panels show the cross section assuming a spectroscopic factor of $C^2S=1.0$ for both PWIA (dashed) and DWIA (solid). The lower panels show the target vector analyzing power A_T^{3b}, the photon asymmetry Σ^{3b}, and the recoil nucleon polarization P_N^{3b} with (dotted) and without (solid) the emitted nucleon spin-orbit potential.

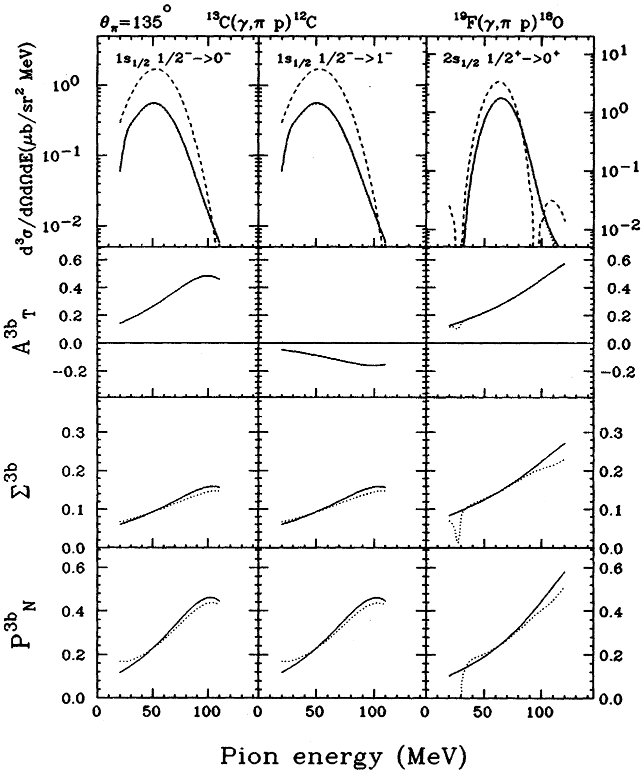

FIG. 3. DWIA energy sharing calculations of $l=0$ nucleon removal for $^{13}\vec{\mathrm{C}}(\gamma,\pi^-p)^{12}\mathrm{C}$ to 0^- and 1^- final states and $^{19}\vec{\mathrm{F}}(\gamma,\pi^+n)^{18}\mathrm{O}$ to the 0^+ ground state at an incident photon energy of 300 MeV and angles $\theta_\pi=135°$ and $\theta_N=-15°$. The top panels show the cross section assuming a spectroscopic factor of $C^2S=1.0$ for both PWIA (dashed) and DWIA (solid). The lower panels show the target vector analyzing power A_T^{3b}, the photon asymmetry Σ^{3b}, and the recoil nucleon polarization P_N^{3b} with (dotted) and without (solid) the emitted nucleon spin-orbit potential.

$$P_s^{M_A}=\frac{(C_+-C_-)}{(C_++C_-)},\quad P_o^{M_A}=\frac{(D_+-D_-)}{(D_++D_-)},$$

$$\alpha^{M_A}=\frac{(E_+-E_-)}{(E_++E_-)}. \tag{26}$$

Carrying out the sum over the emitted nucleon spin, ρ_d'', it is now possible to write a simplified expression for the cross section for a target substate M_A and incoming photon polarization $\epsilon=\pm$ using the two-body analyzing powers as

$$\sigma(M_A,\epsilon=\pm)=\sigma_0^{3b}(1\pm A_{00\epsilon0}^{2b}+A_{000n}^{2b}P_s^{M_A}\pm A_{00\epsilon n}^{2b}P_s^{M_A}$$
$$+PP_o^{M_A}\pm A_{00\epsilon0}^{2b}PP_o^{M_A}+\alpha^{M_A}PA_{000n}^{2b}$$
$$\pm\alpha^{M_A}PA_{00\epsilon n}^{2b}). \tag{27}$$

For $J_A=1/2$ the coefficients defined above take simple forms as $M_A=\pm1/2$. Specifically, $P_s^{M_A}=\pm P_s$, $P_o^{M_A}=\pm P_o$, and $\alpha_s^{M_A}=\alpha$. Values of P_s, P_o, and α are given in Table I for the various transitions. Thus we can write the photon asymmetry

$$\Sigma^{3b}=\frac{A_{00\epsilon0}^{2b}+\alpha PC_{00\epsilon n}^{2b}}{1+\alpha PA_{000n}^{2b}}, \tag{28}$$

and the target vector analyzing power

$$A_T^{3b}=\frac{A_{000n}^{2b}P_s+PP_o}{1+\alpha PA_{000n}^{2b}}. \tag{29}$$

In similar fashion, we can sum over the values of M_A and ϵ to isolate the dependence on the emitted nucleon spin

$$\sigma(\rho_d''=\pm)=\sigma_0^{3b}(1\pm P_{0n00}^{2b}+\alpha PA_{000n}^{2b}\pm\alpha PD_{0n0n}^{2b}) \tag{30}$$

and thus

$$P_N^{3b}=\frac{P_{0n00}^{2b}+\alpha PD_{0n0n}^{2b}}{1+\alpha PA_{000n}^{2b}}. \tag{31}$$

Clearly the situation with $l=1$ is more complicated with the addition of the distortion dependent effective polarization P which mixes the two-body observables. Thus it will be more difficult to separate effects due to distortion from intrinsic modifications of the fundamental process. Moreover, additional complications arise from the spin dependent part of the emitted nucleon potential.

IV. DWIA CALCULATIONS

Having examined the qualitative features of the $\vec{A}(\gamma,\pi N)B$ reaction within the DWIA, we now present some typical calculations for specific nuclear transitions. As a sample set of calculations we present results for 300 MeV incident photons, the peak of the Δ resonance, and an energy regime in which our choices of potentials and amplitudes should be appropriate. As input to the code we need the various optical model potentials and bound state wave functions. In this regard we have used the same phenomenolgical potentials as those used in our previous publications. The proton potentials are from the global fits of Nadasen *et al.* [10], and the pion potentials are of the Kisslinger type with parameters taken from the work of Cottingame and Holtkamp [11]. The bound state wave functions are eigenfunctions of a Woods-Saxon potential with geometrical parameters from the paper of Elton and Swift [12] and the well depth obtained by searching on the nucleon separation energy. The two-body photoproduction amplitudes were taken from the (unitary) parametrization of Blomqvist and Laget [5] assuming an initial energy prescription. While a better treatment of the two-body production process may exist, we have just shown that for the coplanar geometry considered here and neglecting the emitted nucleon spin orbit potential, the three-body spin observables depend solely on the two-body spin observables. Thus as long as the Blomqvist-Laget parametrization reproduces the two-body observables reasonably well, its use in our calculations should be adequate. We show this to be the case in Fig. 1 where we compare the Blomqvist-Laget observables to those calculated with the phase shift analysis code SAID [6] using solution SM95. Future theoretical investigations which examine medium modifications to the two-body amplitudes will probably need to utilize other production models.

In Figs. 2 and 3 we present DWIA calculations for $l=0$ nucleon removal for $^{13}\vec{C}(\gamma,\pi^-p)^{12}C$ to final states with $J^\pi=0^-$ and 1^- and for $^{19}\vec{F}(\gamma,\pi^+n)^{18}O$ to the ground state with $J^\pi=0^+$. The π-meson angles are 45° (Fig. 2) and 135° (Fig. 3) and the proton angle was chosen so that it is kinematically allowed to leave the residual nucleus at rest (quasifree angle pairs). The top panels show both the PWIA and DWIA cross sections. We observe that the dominant effect of the distortion is primarily attenuation, reducing the cross section by a factor of four to ten for the deeply bound $1s$ orbital of ^{12}C and a factor of two to four for the $2s$ proton in ^{19}F. In the lower panels we present the target analyzing power A_T^{3b}, the photon asymmetry Σ^{3b}, and the recoil nucleon polarization P_N^{3b}, assuming pure shell model configurations. Calculations were done both with (dotted curves) and without (solid curves) the spin-orbit potential for the emitted nucleon. As expected from Eqs. (18)–(21), the solid curves (no spin orbit) are simply the two-body analyzing powers, where in the case of A_T^{3b}, it is scaled by the effective spin polarization listed in Table I. The variation with emitted pion energy simply reflects the change in the two-body kinematics due to the Fermi motion of the bound nucleons. The differences between ^{13}C and ^{19}F are simply due to the differing Q values which changes the kinematics for the two-body prescription.

The effects of the emitted nucleon spin-orbit potential on the cross section and target analyzing power A_T^{3b} are very small, typically a few percent, and can basically be ignored. For $\theta_\pi=135°$ and near the peak of the cross section for $\theta_\pi=45°$, the effects of the spin orbit potential on the photon asymmetry and recoil nucleon polarization are also small. However, in regions where the cross sections are small, the spin orbit effects are significant, particularly for ^{19}F near the node in the $2s$ wave function. In spite of this the DWIA predicts an average behavior of the spin observables which is basically the same as the two-body values, even with the inclusion of the spin-orbit potential, and changes in analyzing power due to medium modifications should be readily apparent in the data.

We now consider the cases of p-shell nucleon removal. As noted early for these $l=1$ transitions the distortion effects introduce an effective polarization P defined in Eq. (22). In Fig. 4 we present calculations of P for $^{13}C(\gamma,\pi^-p)^{12}C$. We see that the effective polarization can be quite large, and that it changes sign at the quasifree point. Clearly the distortions can play a significant role in determining the three-body spin observables.

In Figs. 5 and 6 we present DWIA calculations for $^{13}C(\gamma,\pi^-p)^{12}C$. Although the transitions to the 1^+ states can have contributions from both $j=1/2$ and $j=3/2$ nucleon removal, to simplify the interpretation we will present results only for pure j transitions. Thus we present calculations for $1p_{1/2}$ removal to 0^+ and 1^+ states, and $1p_{3/2}$ removal to 1^+ and 2^+ states. The calculations are presented in the same format as in Figs. 2 and 3. Due to the polarization induced by distortions, a variety of contributions arise to obscure the simple one-to-one correspondence between three-body and two-body spin observables which we observed for $l=0$. For example, Eqs. (28)–(31) show contributions from the two-

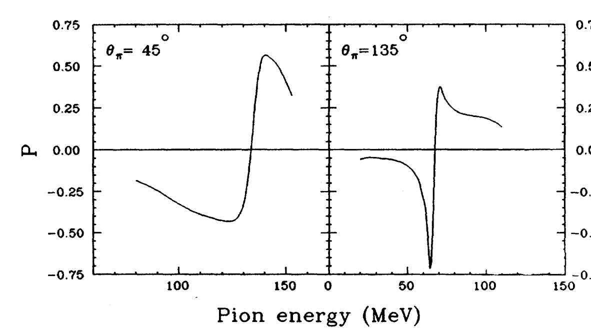

FIG. 4. DWIA calculations of the effective polarization P for $^{13}\vec{C}(\gamma,\pi^-p)^{12}C$ for $l=1$ removal. The incident photon energy is 300 MeV and the emitted pion angles are $\theta_\pi=45°$ and $\theta_\pi=135°$. The kinematical conditions are those appropriate to the calculations presented in Figs. 4 and 5.

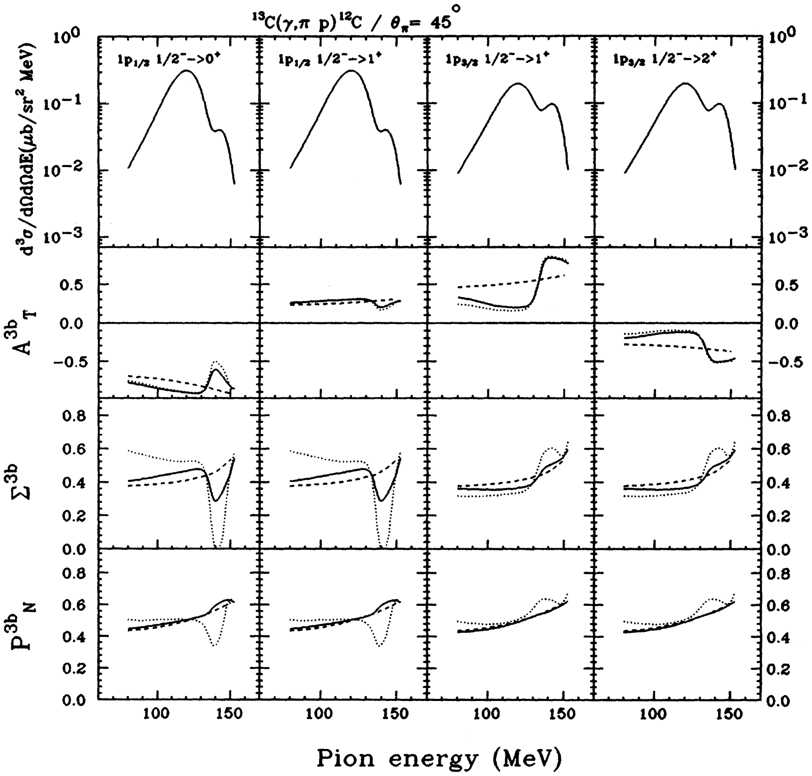

FIG. 5. DWIA energy sharing calculations of $^{13}\vec{C}(\gamma,\pi^- p)^{12}C$ for $l=1, j=1/2$ nucleon removal to 0^+ and 1^+ states, and $l=1, j=3/2$ nucleon removal to 1^+ and 2^+ states. The incident photon energy is 300 MeV and the emitted particle angles are $\theta_\pi=45°$ and $\theta_N=-51°$. The top panels show the cross section assuming a spectroscopic factor of $C^2S=1.0$. The lower panels show the target vector analyzing power A_T^{3b}, the photon asymmetry Σ^{3b}, and the recoil nucleon polarization P_N^{3b} for PWIA (dashed) and DWIA with (dotted) and without (solid) the emitted nucleon spin orbit potential.

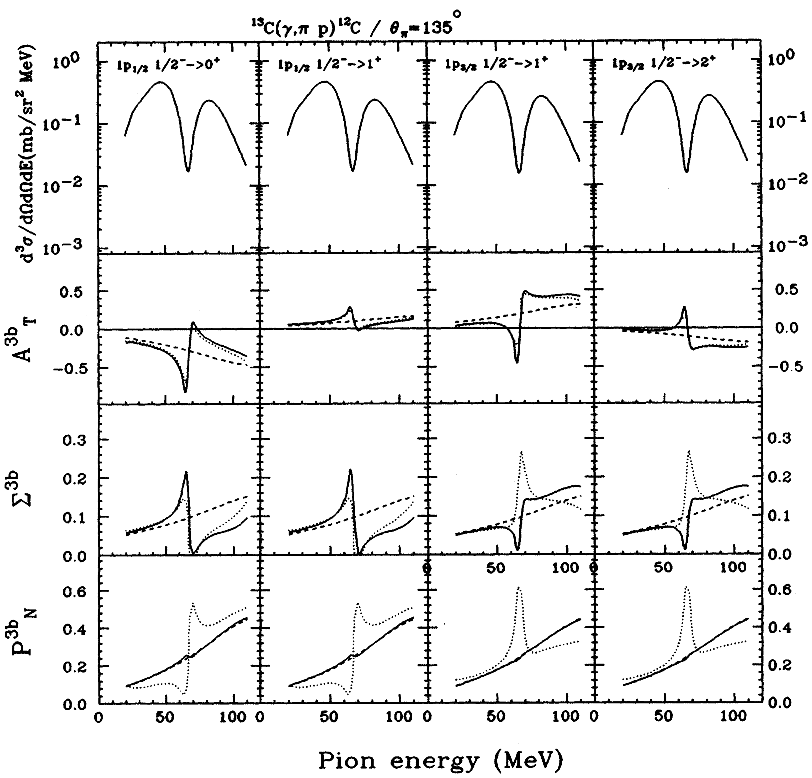

FIG. 6. DWIA energy sharing calculations of $^{13}\vec{C}(\gamma,\pi^- p)^{12}C$ for $l=1, j=1/2$ nucleon removal to 0^+ and 1^+ states, and $l=1, j=3/2$ nucleon removal to 1^+ and 2^+ states. The incident photon energy is 300 MeV and the emitted particle angles are $\theta_\pi=135°$ and $\theta_N=-15°$. The top panels show the cross section assuming a spectroscopic factor of $C^2S=1.0$. The lower panels show the target vector analyzing power A_T^{3b}, the photon asymmetry Σ^{3b}, and the recoil nucleon polarization P_N^{3b} for PWIA (dashed) and DWIA with (dotted) and without (solid) the emitted nucleon spin orbit potential.

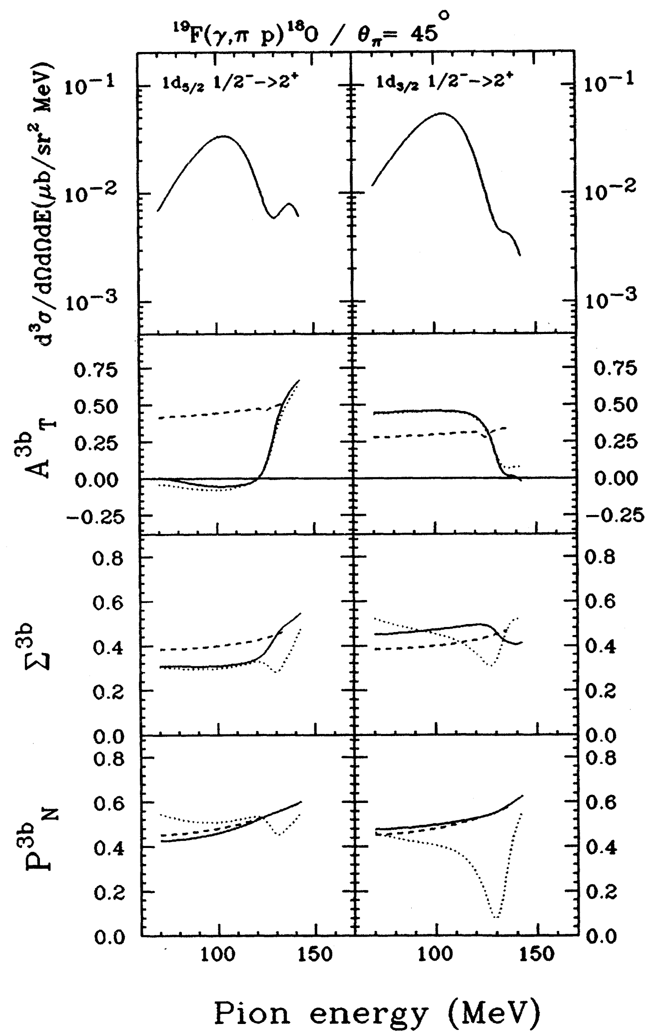

FIG. 7. DWIA energy sharing calculations of $^{19}\vec{F}(\gamma,\pi^+n)^{18}O$ for $l=2, j=3/2$ nucleon removal and $l=2, j=5/2$ nucleon removal to 2^+ states. The incident photon energy is 300 MeV and the emitted particle angles are $\theta_\pi=45°$ and $\theta_N=-51°$. The top panels show the cross section assuming a spectroscopic factor of $C^2S=1.0$. The lower panels show the target vector analyzing power A_T^{3b}, the photon asymmetry Σ^{3b}, and the recoil nucleon polarization P_N^{3b} for PWIA (dashed) and DWIA with (dotted) and without (solid) the emitted nucleon spin orbit potential.

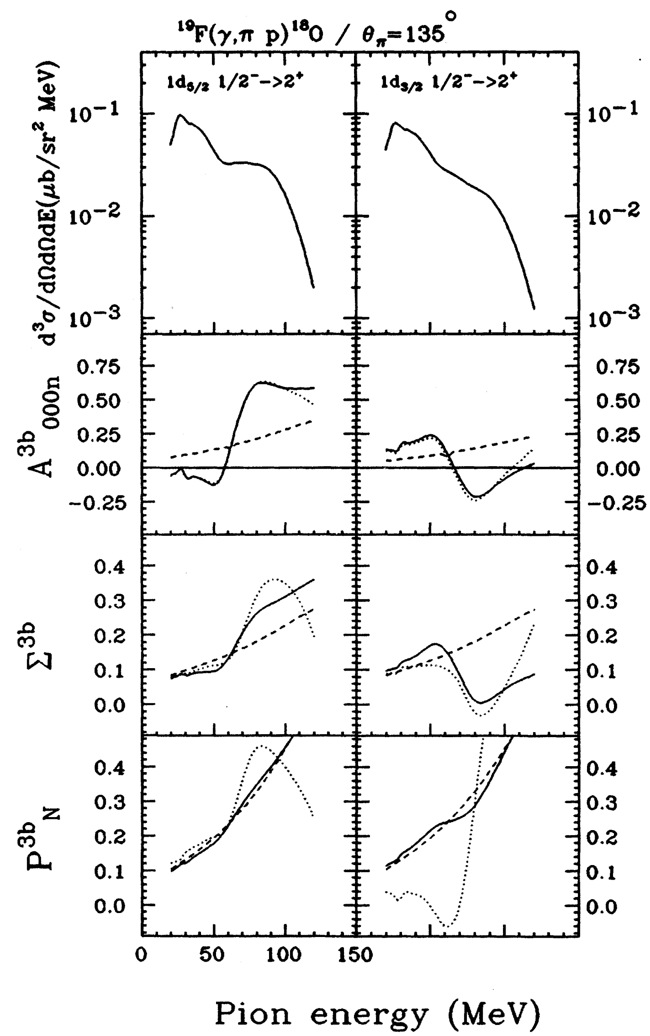

FIG. 8. DWIA energy sharing calculations of $^{19}\vec{F}(\gamma,\pi^+n)^{18}O$ for $l=2, j=3/2$ nucleon removal and $l=2, j=5/2$ nucleon removal to 2^+ states. The incident photon energy is 300 MeV and the emitted particle angles are $\theta_\pi=135°$ and $\theta_N=-15°$. The top panels show the cross section assuming a spectroscopic factor of $C^2S=1.0$. The lower panels show the target vector analyzing power A_T^{3b}, the photon asymmetry Σ^{3b}, and the recoil nucleon polarization P_N^{3b} for PWIA (dashed) and DWIA with (dotted) and without (solid) the emitted nucleon spin orbit potential.

body spin correlation and spin transfer coefficients. To show the role of these additional terms, we also present the spin observables with the effective polarization P set equal to zero—basically a PWIA calculation of the spin observables. These calculations are indicated by the dashed curves in Figs. 5 and 6.

Clearly the additional terms due to distortion play a major role in the determination of the target analyzing powers and photon asymmetry. The effects are generally most pronounced in the region of the quasifree point (corresponding to the minimum in the cross section) where P changes sign, and are also of opposite sign for the spin orbit partners, $p_{1/2}$ and $p_{3/2}$, due to a change in the sign of α. The observable least affected by the emitted nucleon spin orbit potential is the target analyzing power, A_T^{3b}. However, the effects of spin independent distortions through P are significant. The most robust spin observable in terms of insensitivity to distortion effects would appear to be the recoil nucleon polarization (bottom panels of Figs. 5 and 6). In this case the differences between the no spin orbit DWIA and PWIA are less than 10%. However, this result is quite misleading. The corrections due to the additional terms in the numerator (αPD_{0n0n}^{2b}) and the denominator (αPA_{000n}^{2b}) can be quite large, changing each by as much as a factor of two compared to the PWIA case. In this energy range these corrections tend to cancel. The same statement applies, but to a lesser extent, to the target analyzing power and photon asymmetry, so that the effects of the terms involving P do not appear to be pronounced. The cancellation is basically a consequence of the fact that the correction terms in the numerator and denominator are of the same sign, so that the overall effect of the distortions on the spin observables is reduced. Moreover, the role of the spin orbit potential can be significant in both cases. Clearly to use these observables as a probe of in-medium modifications of the two-body interaction is a dangerous course which could lead to misleading conclusions.

To summarize the results for $l=1$ we note that distortion effects play an important role. Thus modifications to the spin observables can arise not simply from more complicated reaction dynamics, but also from the more mundane effects of distortions in the initial and final state. For this reason it will be essential to have a proper treatment of distortion effects in comparing to experimental data for $l=1$ transitions. It will clearly be desirable to also have data for $l=0$ transitions in which distortion effects play a minimal role (through the spin orbit potential of the emitted nucleon).

Finally for completeness for ^{19}F we present results for d-shell removal to the 2^+ state of ^{18}O in Figs. 7 and 8. Again we have assumed single j value transitions to simplify the interpretation. For these cases the effective polarization is more complicated since there are three $T^{l\lambda}$ amplitudes which are nonzero leading to tensor polarization. To show the effects of distortion we again include a PWIA calculation of the spin observables in the figures. The calculations show effects similar to the p-shell case with an effective vector polarization which is of opposite sign for $d_{3/2}$ and $d_{5/2}$ and changes sign at the zero recoil momentum point. Clearly, distortion effects play an important role.

V. CONCLUSIONS

We have examined the photopion production process in nuclei $A(\gamma,\pi N)B$ using a factorized DWIA formulation. Assuming a coplanar geometry and neglecting the spin-orbit potential for the emitted nucleon, we obtained simplified expressions for the photon asymmetry Σ^{3b}, the target asymmetry A_T^{3b}, and the emitted nucleon polarization P_N^{3b}. These observables can be written as sums of products of the two-body pion photoproduction observables, the nuclear structure dependent struck nucleon polarization arising from the polarization of the target, and an effective polarization created by the distortion effects. For s-wave nucleon removal the distortion effects are zero, and in the case of nuclear targets with $J_A=1/2$ particularly simple relationships to the two-body observables are found. For higher angular momentum states distortion effects contribute and explicit expressions were obtained for p-wave nucleon removal. Again for $J_A=1/2$ rather simple expressions were obtained.

This discussion was followed by numerical calculations of the observables for $J_A=1/2$ nuclei and various final states. These calculations show the importance of distortion effects for $l=1$ and that near the quasifree point the effects of the emitted nucleon spin orbit potential are small.

The expressions developed in this paper provide a powerful tool with which to design experiments to examine the photopion production process in the nuclear medium. As we show, for s-wave removal distortion effects are essentially negligible in the quasifree region, making such transitions ideal for the study of medium effects. Of the observables examined here only the target asymmetry depends on the nuclear structure, and the dependence merely involves the target and residual nucleus total angular momenta. Thus it is trivial to calculate the struck nucleon polarization from Eq. (17). The removal of nucleons with $l>0$ depends on the distortion through the effective polarization P which can mask medium effects. However such measurements provide excellent tests of the reaction model and provide more extensive tests of the production process, since the observables contain contributions from the spin correlation coefficient and the polarization transfer coefficient. The size of these contributions to the three-body spin observables can be estimated from our expressions.

We note that we have not examined in any detail the effects of the nuclear structure uncertainties, since we have considered pure single orbital transitions. Fortunately numerous such transitions exist in nature. However, it is clear from our calculations that the structure can have a significant effect on the spin observables, particularly for $l>0$. For example, transitions $J_A=(1/2)^-$ to $J_B=1^+$ will have contributions from both $p_{1/2}$ and $p_{3/2}$ orbitals. Since various coefficients change sign (see Table I) for the two orbitals, the result will depend on the relative mixture of the two contributions. The effect of these admixtures can be explored using the equations provided here, but a good estimate will require DWIA calculations since the reaction is coherent in j.

ACKNOWLEDGMENT

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[1] Xiaodong Li, L. E. Wright, and C. Bennhold, Phys. Rev. C **48**, 816 (1993).
[2] N. S. Chant and P. G. Roos, Phys. Rev. C **38**, 787 (1988).
[3] Mohammad G. Khayat, N. S. Chant, P. G. Roos, and T.-S. H. Lee, Phys. Rev. C **46**, 2415 (1992).
[4] Mohammad G. Khayat, N. S. Chant, and P. G. Roos, Phys. Rev. C **51**, 776 (1995).
[5] I. Blomqvist and J. L. Laget, Nucl. Phys. **A280**, 405 (1977).
[6] Zhujun Li, Richard A. Arndt, L. David Roper, and Ron L. Workman, Phys. Rev. C **47**, 2759 (1993); R. A. Arndt, R. L. Workman, Z. Li, and L. D. Roper, *ibid.* **42**, 1853 (1990).
[7] J. Bystricky, F. Lehar, and P. Winternitz, J. Phys. (Paris) **39**, 1 (1978).
[8] P. Kitching, W. J. McDonald, Th. A. Maris, and C. A. Z. Vas-

concellos, in *Advances in Nuclear Physics,* edited by J. W. Negele and E. Vogt (Plenum, New York, 1985), Vol. 15, p. 43.
[9] H. C. Newns, Proc. Phys. Soc. London Sect. A **66**, 477 (1953); H. C. Newns and M. Y. Refai, *ibid.* **71**, 627 (1958).
[10] A. Nadasen, Ph.D. thesis, Indiana University, 1977; 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] W. B. Cottingame and D. B. Holtkamp, Phys. Rev. Lett. **45**, 1828 (1980).
[12] L. R. B. Elton and A. Swift, Nucl. Phys. **A94**, 52 (1967).