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Analysis of the High-Energy Photonucleon Emission from O^{16} and C^{12} in a Direct Reaction Model

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High-energy photonucleon emission probabilities in O^{16} and C^{12} have been studied by (a) the standard direct-emission model and (b) the standard direct-emission model modified to incorporate the ground-state correlation. The calculation has been made for the 90° differential cross section, integrated cross section, angular distributions at a few energies, and 45° polarization. Inclusion of the ground-state correlation improves the agreement between the theoretical results and the observations. This, along with the sensitivity of the absolute magnitude of the photonucleon cross section to the extent of the correlation, may be exploited to study the composition of the wave functions of the target and the residual nucleus.

I. INTRODUCTION

THE experimental data of the photoproton and photoneutron emission from light nuclei indicate considerable cross section beyond the giant-dipole-resonance region.[1] As is well known, such a high cross section is not compatible with a compound-nucleus process but is indicative of direct emission.[2–4] The partial success of an intermediate-coupling model in providing an empirical fit to the fine structure of the giant-dipole resonance also implies a considerable strength for the direct transition probability in the giant-dipole region.[5] This has been conjectured earlier.[3,6] At energies higher than the giant-dipole-resonance energies, the direct emission probabilities are expected to be dominant,[7] and the purpose of this paper is to investigate to what extent the observed yield of the photonucleons is consistent with the direct transition probabilities. We examine this within the framework of (a) the standard direct-emission model, and (b) the standard direct-emission model modified to incorporate the ground-state correlation. This treatment of the direct emission differs from the one used by Shklyarevskii[8] to the extent that the bound-state wave functions used here are eigenfunctions in an appropriate Woods-Saxon potential with a spin-orbit interaction, as opposed to the simple-harmonic-oscillator functions used by him, and the ground-state correlation has been included here for the first time.

II. THEORY

In the direct-reaction model, the incident photon interacts with a bound nucleon and lifts it directly from one of the bound shells into the continuum without sharing energy with other nucleons. For such a model, the differential cross section for the emission of a nucleon with wave vector **k** by an incident photon

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[1] C.-P. Wu, F. W. K. Firk, and T. W. Phillips, Phys. Rev. Letters **10**, 1182 (1968).

[2] E. D. Courant, Phys. Rev. **82**, 703 (1951).

[3] G. E. Brown and J. S. Levinger, Proc. Phys. Soc. (London) **A71**, 733 (1958).

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[5] C. B. Duke and F. B. Malik, in *Proceedings of the International Conference on Nuclear Physics, Gatlinberg, Tenn., 1966*, edited by R. L. Becker (Academic Press Inc., New York, 1967); C. B. Duke, F. B. Malik, and F. W. K. Rirk, Phys. Rev. **157**, 879 (1967).

[6] D. H. Wilkinson, Physica **22**, 1039 (1956).

[7] M. G. Mustafa and F. B. Malik, Bull. Am. Phys. Soc. **14**, 607 (1969).

[8] G. M. Shklyarevskii, Zh. Eksperim. i Teor. Fiz. **36**, 1492 (1959); **41**, 451 (1962) [English Transls.: Soviet Phys.—JETP **9**, 1057 (1959); **14**, 324 (1962)].

of wave vector $\mathbf{k}_\gamma$ is given by[2,4,9,10]

$$d\sigma(\mathbf{k}_\gamma,\mathbf{k})/d\Omega=\tfrac{16}{9}\pi(E_\gamma/\hbar v)(k^2/\hbar c)[2(2I_A+1)]^{-1}\times[\sum_{\nu\nu'}M(E1)M^*(E1)X_{1\nu}(\hat{k}_\gamma)\cdot X_{1\nu'}{}^*(\hat{k}_\gamma)], \quad (1)$$

where E_γ, v, and I_A are, respectively, the energy of the incident γ ray, the velocity of the outgoing nucleon, and the total spin of the target nucleus. The summation also includes the sum over final-state orbital and total angular momenta. $X_{1\nu}$ is the vector spherical harmonic and is defined by

$$X_{1\nu}(\hat{k}_\gamma)=\sum_{\rho\epsilon}(11\rho\epsilon\mid 1\nu)Y_{1\rho}(\hat{k}_\gamma)\chi_{1\epsilon}, \quad (2)$$

where $\chi_{1\epsilon}$ is the spin function orthonormalized to

$$\langle\chi_{1\epsilon}\mid\chi_{1\epsilon'}\rangle=\delta_{\epsilon\epsilon'}.$$

$M(E1)$ is the matrix element of the electric dipole operator between the initial state Ψ_i and the final state Ψ_f of a system of A nucleons and is given by

$$M(E1)=\langle\Psi_f\mid O\mid\Psi_i\rangle. \quad (3)$$

The operator O is, in principle, a sum of single-particle operators,

$$O=\sum_{n=1}^{A}O(\hat{r}_n)F(r_n), \quad (4)$$

where for the electric dipole transition $O(\hat{r}_n)=e'\tau_3Y_{1\nu}(\hat{r}_n)$, and $F(r_n)=r_n$. Here e' is the effective charge on the nucleons, which is $+eN/A$ for protons and $-eZ/A$ for neutrons. τ_3 is the z component of the isospin and has eigenvalues $\hbar/2$ for protons and $-\hbar/2$ for neutrons.

Restricting ourselves to the photoemission of a nucleon from a particular closed shell, we note that (3) reduces to N integrals with the same radial part where N is the number of nucleons in the closed shell.

The initial and the final wave functions are approximated in terms of a product of the $(A-1)$ core nucleon wave function Φ and the participating nucleon wave function ϕ:

$$\Psi_i=\sum_{M_1M_2}(I_1I_2M_1M_2\mid I_AK)\Phi_i(\text{core},I_1M_1)\phi_i(I_2M_2), \quad (5)$$

where I_1, I_2, and I_A are, respectively, the initial angular momenta of the core, of the participating nucleon, and of the system as a whole. M_1, M_2, and K are their respective z components. Since we are concerned with even-even target nuclei the only type of Clebsch-Gordan coefficient permitted is $(I\,I\,M\,-M\mid 00)$, and

$$\Psi_f=\sum_{M_0m}(I_0jM_0m\mid JM)\Phi_f(\text{core},I_0M_0)\phi_f(jm), \quad (6)$$

where I_0, j, and J are, respectively, the angular momenta of the residual nucleus, of the outgoing nucleon, and of the final system. M_0, m, and M are their respective z components.

The initial and the final wave functions of the participating nucleon can be decomposed in terms of its orbital angular momentum and intrinsic spin as follows:

$$\phi_i(I-M)=\sum_{\sigma\Lambda}(L\tfrac{1}{2}\Lambda\sigma\mid I-M)i^LY_{L\Lambda}(\hat{r})\times u_{LI\tau}(E_B,r)r^{-1}\chi_{1/2,\sigma}\chi_{t\tau}, \quad (7)$$

where $r^{-1}u_{LI\tau}$ is the radial wave function and E_B is the binding energy of the nucleon in consideration, and

$$\phi_f(jm)=(4\pi/k)\sum_{\lambda\mu}(l\tfrac{1}{2}\lambda\mu\mid jm)Y_{l\lambda}{}^*(\hat{k})\times[\sum_{\lambda'\mu'}(l\tfrac{1}{2}\lambda'\mu'\mid jm)i^lY_{l\lambda'}(\hat{r})u_{lj\tau}(E,r)\chi_{1/2,\mu'}\chi_{t\tau}], \quad (8)$$

where $\mathbf{k}$ is the wave vector of the outgoing nucleon, and $r^{-1}u_{lj\tau}$ is the radial wave function. This can be written in terms of an incoming wave function $u_{lj\tau}{}^{(-)}$, an outgoing wave function $u_{lj\tau}{}^{(+)}$, and the phase shift $\delta_{lj\tau}$. Explicitly,

$$u_{lj\tau}=[e^{2i\delta}u^{(+)}-u^{(-)}/2i]_{lj\tau}{}^*. \quad (9)$$

In the direct-reaction model, $u^{(\pm)}$ are solutions of an appropriate optical-model equation.[11]

Using Eqs. (4)–(8), the transition matrix element reduces to

$$M(E1)=-[e'(4\pi)^{1/2}/k]\sum_{M_0\lambda\mu}(I_0jM_0m\mid 1\nu)(l\tfrac{1}{2}\lambda\mu\mid jm)\times Y_{l\lambda}(\hat{k})E_1(ILjl\tau)\langle\Phi_f(\text{core},I_0M_0)\mid\Phi_i(\text{core},IM)\rangle, \quad (10)$$

where

$$E_I(ILjl\tau)=\tau i^{L-l}1/\sqrt{3}[(2j+1)(2I+1)]^{1/2}(-)^{j+1/2}\times\tfrac{1}{2}[1+(-)^{l+L+1}](Ij-\tfrac{1}{2}\tfrac{1}{2}\mid 1\,0)\times\int_0^\infty dr\,u_{lj\tau}{}^*(E,r)r\,u_{LI\tau}(E_B,r) \quad (11)$$

and

$$\langle\chi_{t\tau}\mid\tau_3\mid\chi_{t\tau}\rangle=\tau.$$

The differential cross section is thus given by

$$d\sigma(\mathbf{k}_\gamma,\mathbf{k})/d\Omega=12e'^2(E_\gamma/\hbar c)(\hbar v)^{-1}[2(2I_A+1)]^{-1}\times\sum_{jj'll'Q}A_{fi}[(2j+1)(2j'+1)]^{1/2}\tfrac{1}{2}[1+(-)^{l+l'+Q}]\times(-)^{I+1/2}(jj'\tfrac{1}{2}-\tfrac{1}{2}\mid Q0)(1100\mid Q0)W(j1j'1;IQ)\times E_1(ILjl\tau)E_1{}^*(ILj'l'\tau)P_Q(\cos\theta), \quad (12)$$

where A_{fi} represents the square of the core overlap integral, $\langle\Phi_f(\text{core},I_0M_0)\mid\Phi_i(\text{core},IM)\rangle^2$, θ is the angle

[9] J. S. Levinger, *Nuclear Photodisintegration* (Oxford University Press, London, 1960).

[10] B. Buck and A. D. Hill, Nucl. Phys. **A95**, 271 (1967).

[11] E. H. Auerbach, N. C. Francis, D. T. Goldman, and C. R. Lubitz Knolls Atomic Power Laboratory Report No. KAPL-3020, Schenectady, N.Y., 1964 (unpublished).

TABLE I. Potential parameters used in the optical model, and single-particle energies in O^{16} and C^{12}.

		Protons				Neutrons			
Target	States	Energy (MeV)	$V_{re}(0)$ (MeV)	$V_{im}(0)$ (MeV)	$V_{s.o.}(0)$ (MeV)	Energy (MeV)	$V_{re}(0)$ (MeV)	$V_{im}(0)$ (MeV)	$V_{s.o.}(0)$ (MeV)
O^{16}	$1P_{1/2}$	−12.11	57.95	···	9.89	−15.65	57.39	···	9.64
	$1P_{3/2}$	−18.44	57.95	···	9.89	−21.81	57.39	···	9.64
	$d_{3/2}, d_{5/2}$	10–40	54.94	5.0	5.27	10–40	54.94	5.0	5.27
	$s_{1/2}$	10–40	54.94	5.0	···	10–40	54.94	5.0	···
C^{12}	$1P_{3/2}$	−15.96	57.5	···	9.7	−18.72	57.0	···	9.7
	$d_{3/2}, d_{5/2}$	10–40	55.0	5.0	5.27	10–40	55.0	5.0	5.27
	$s_{1/2}$	10–40	55.0	5.0	···	10–40	55.0	5.0	···
	$a=0.51$ F,	$b=0.81$ F,	$r_0=1.2$ F,	$R=R_c=r_0(A-1)^{1/3}$ F					

between $\mathbf{k}_\gamma$ and $\mathbf{k}_n$, and $W(abcd;\, ef)$ is a Racah coefficient. The cross section integrated over the angle θ is given by

$$\sigma(\mathbf{k}_\gamma, \mathbf{k}) = \tfrac{16}{3}\pi e'^2 (E_\gamma/\hbar c)(\hbar v)^{-1}[2(2I_A+1)]^{-1} \times \sum_{jl} A_{fi}\,|E_1(ILjl\tau)|^2, \quad (13)$$

where $E_1(ILjl\tau)$ is given by (11).

The expressions (12) and (13) reduce to the standard expressions of the direct emission if $A_{fi}=1$. This core overlap factor, however, gives a measure of that part of the interaction of the outgoing nucleon with the nucleons of the residual nucleus, which is not included in the average static potential. It represents the extent of the purity of the single-particle state of the emitted nucleon in the target nucleus. So long as the process is a direct one, this overlap factor does not depend on the energy of the reaction and therefore must be included throughout the whole energy range. The absolute magnitude of the cross sections is influenced by this factor.

A. Polarization

Choosing the system of reference in which the xy plane is the reaction plane and the z axis is the direction of the vector $(\mathbf{k}_\gamma \times \mathbf{k})$, where $\mathbf{k}_\gamma$ is the incident photon wave vector and $\mathbf{k}$ is the outgoing nucleon wave vector, the expression for the polarization reduces to[12]

$$\begin{aligned} P(\theta) = {} & [D(\theta)]^{-1}[(3\times 5^{1/2})/2]\sin 2\theta \sum_{jj'll'} i(-)^{I+j+l} \\ & \times[(2l+1)(2l'+1)(2j+1)(2j'+1)]^{1/2}(ll'00\,|\,20) \\ & \times(1100\,|\,20)W(1111;12)W(j1j'1;I2) \\ & \times\begin{Bmatrix} 2 & j & j' \\ 2 & l & l' \\ 1 & \frac{1}{2} & \frac{1}{2}\end{Bmatrix} E_1(ILjl\tau)E_1^*(ILj'l'\tau), \end{aligned} \quad (14)$$

[12] W. Czyz and J. Sawicki, Nuovo Cimento 3, 864 (1956).

where $D(\theta)$ is given by the right-hand side of Eq. (12) without the numerical factors before the summation sign, and the quantity in the bracket is the 9-j symbol of Wigner.

The final-state continuum wave function is calculated using an optical-model potential

$$-V_{\text{opt}} = V_{\text{re}}(r) + V_{\text{im}}(r) + V_{\text{s.o.}}(r)\,\text{l.s.} - V_{\text{Coul}}(r) \quad (15)$$

with

$$\begin{aligned} V_{\text{re}}(r) &= V_{\text{re}}(0)/\{1+\exp[(r-R)/a]\}, \\ V_{\text{im}}(r) &= V_{\text{im}}(0)\exp\{-[(r-R)/b]^2\}, \\ V_{\text{s.o.}}(r) &= V_{\text{s.o.}}(0)[\hbar/m_\pi c]^2[V_{\text{re}}(0)]^{-1}r^{-1}\,|\,dV_{\text{re}}(r)/dr\,| \\ V_{\text{Coul}}(r) &= (Z-1)e^2/r, \qquad r>R_c \\ &= [(Z-1)e^2/2R_c](3-r^2/R_c^2) \qquad r<R_c. \end{aligned}$$

The parameters used are listed in Table I. For the O^{16} case, they are similar to those in Ref. 10.

III. APPLICATIONS AND RESULTS

A. O^{16} Target

1. *No Ground-State Correlation*

The ground state of O^{16} nucleus is assumed to be described by a simple nuclear shell model without any correlation. Therefore, depending on the incident photon energy, direct transitions can take place from the $1S_{1/2}$, $1P_{1/2}$, or $1P_{3/2}$ bound levels to the continuum. However, the excitation energy required for transitions from the $1S_{1/2}$ level is about 50 MeV. Since we are interested in the region below the excitation energy of 50 MeV, we have neglected transitions from this level.

The electric dipole interaction involves the transition between the following single-particle states (i) $1P_{1/2}\to s_{1/2}$, $1P_{1/2}\to d_{3/2}$, and (ii) $1P_{3/2}\to s_{1/2}$, $1P_{3/2}\to d_{3/2}$, $1P_{3/2}\to d_{5/2}$. The first set of transitions is pertinent to the $O^{16}{}_{\text{g.s.}}$ and the $O^{16}(\gamma,\, p_0)N^{15}{}_{\text{g.s.}}$ reactions, while

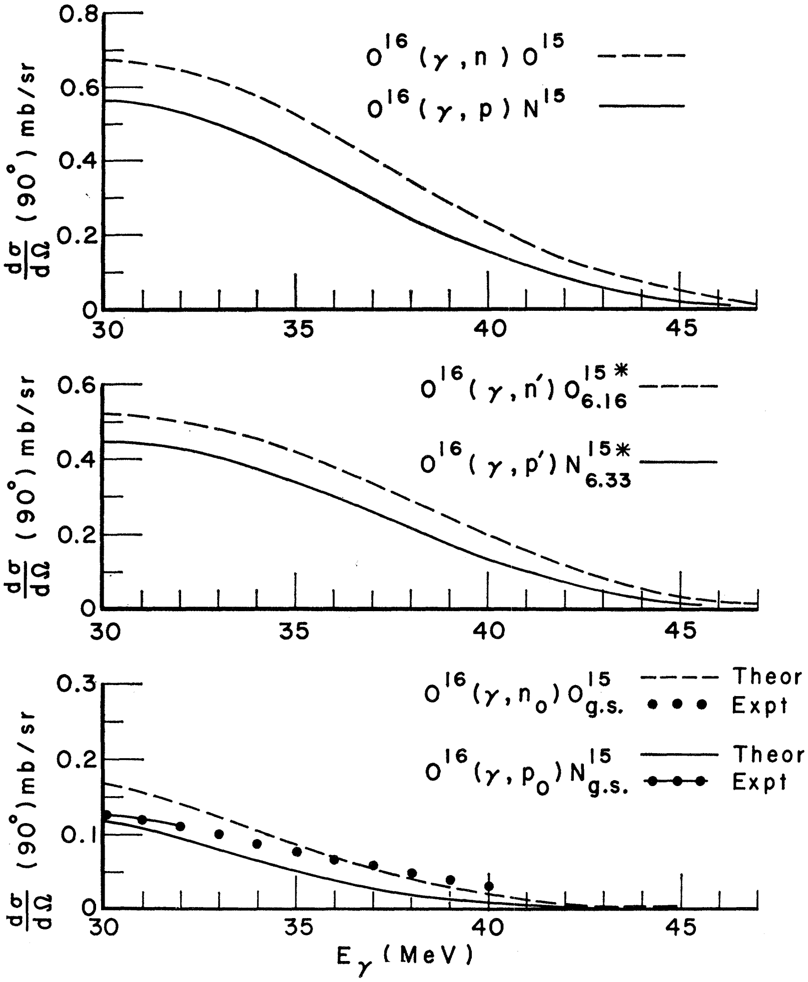

FIG. 1. The 90° differential cross section of the $O^{16}(\gamma, n_0)O^{15}_{g.s.}$ and the $O^{16}(\gamma, p_0)N^{15}_{g.s.}$; and the $O^{16}(\gamma, n)O^{15*}_{6.16}$ and the $O^{16}(\gamma, p')N^{15*}_{6.33}$ reactions. The (γ, n_0) experimental data are from the work of Wu *et al.*,[1] and the (γ, p_0) data (only up to $E_\gamma = 32$ MeV) are from Morrison *et al.*[14] The (γ, n) is the sum of (γ, n_0) and (γ, n') reactions and, similarly, (γ, p) is the sum of (γ, p_0) and (γ, p') reactions.

the second set involves the $O^{16}(\gamma, n')O^{15*}_{6.16}$ and the $O^{16}(\gamma, p')N^{15*}_{6.33}$ reactions.

To compute the theoretical cross section, the exact locations of energies of the bound states are necessary. The location of the $1P_{1/2}$ states can easily be obtained from the mass difference. Unfortunately, the $(1P_{3/2})$ state is not very well known experimentally, but the values -21.81 MeV for neutrons and -18.44 MeV for protons are consistent with the interpretation that the $\frac{3}{2}^-$ states at 6.16 MeV in O^{15*} and at 6.33 MeV in N^{15*} are holes in the $(1P_{3/2})$ shell. These values are also used in most recent calculations.[10] A summary of these energies and the potential parameters used to obtain the bound and the continuum wave functions is presented in Table I. It may be noted that the choice of the imaginary part of the optical potential used to describe the continuum wave function is important to get the correct order of magnitude of the cross section. In all the calculations, the imaginary potential is kept at 5 MeV. Similar values have been used by Lutz *et al.*[13] in the optical-model analysis of neutron scattering in light nuclei.

Using the core overlap factor $A_{fi} = 1$, the cross sections are calculated from Eqs. (12) and (13). The 90° differential cross section and the integrated cross section are shown in Figs. 1 and 2, respectively. The 90° cross section is compared with the experimental findings of Wu *et al.*[1] and Morrison *et al.*[14] There seems to be reasonable agreement with the experimental conclusions that the high-energy cross section, above the giant-dipole region, does not show any structure and that it falls off smoothly as the incident photon energy increases.

Angular distributions of the neutrons and protons from the $O^{16}(\gamma, n_0)O^{15}_{g.s.}$ and the $O^{16}(\gamma, p_0)N^{15}_{g.s.}$

[13] H. F. Lutz, J. B. Mason, and M. D. Karvelis, Nucl. Phys. **47**, 521 (1963).

[14] R. C. Morrison, J. R. Stewart, and J. S. O'Connell, Phys. Rev. Letters **15**, 509 (1965); R. C. Morrison, thesis, Yale University, 1965 (unpublished).

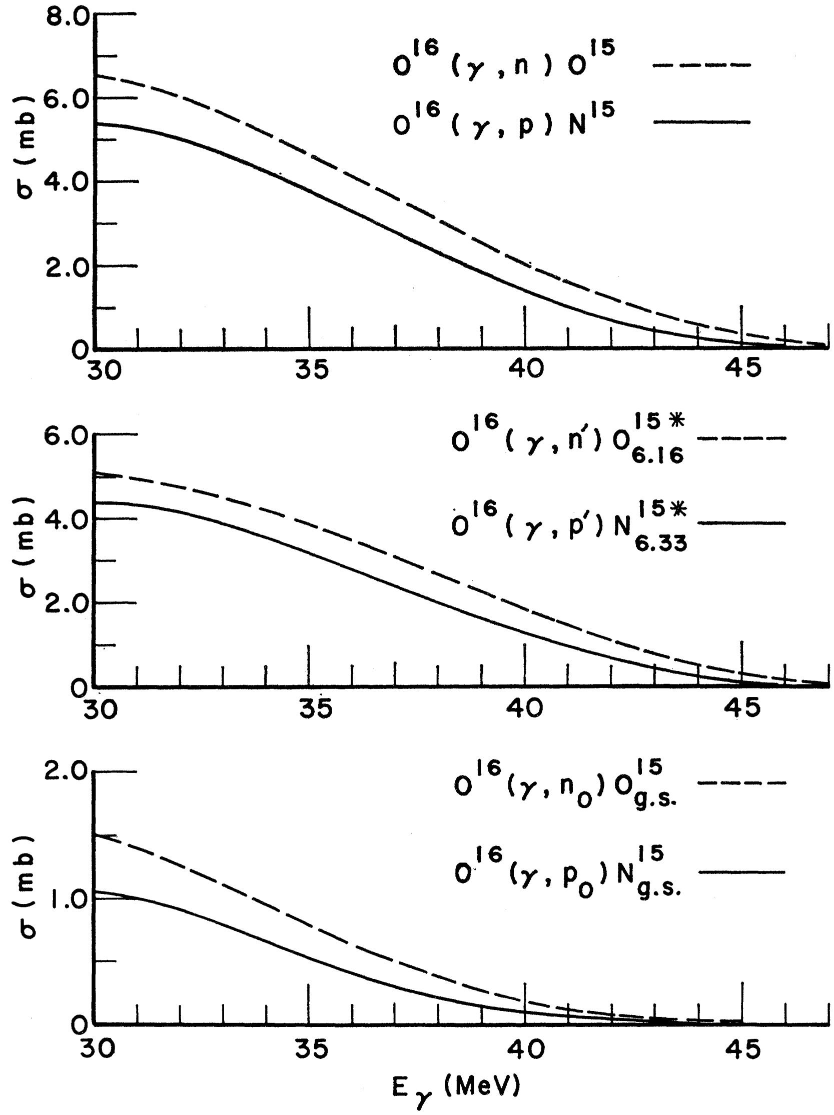

FIG. 2. The integrated cross section of the $O^{16}(\gamma, n_0)O^{15}{}_{g.s.}$ and the $O^{16}(\gamma, p_0)N^{15}{}_{g.s.}$, and the $O^{16}(\gamma, n')O^{15*}{}_{6.16}$ and the $O^{16}(\gamma, p')N^{15*}{}_{6.33}$ reactions. The (γ, n) and (γ, p) have the same meaning as in Fig. 1.

reactions are shown in Figs. 3 and 4, respectively. The computed distribution is normalized to unity at $\theta_{c.m.}=90°$ and is symmetrical about 90°. No definite experimental information is yet available in this energy region. However, the asymmetry that has been observed in the giant-resonance region, might also be present in this energy region. There are some indications that this may be the case.[15] A part of this asymmetry may be attributed to the interference between dipole and higher multipole transitions.

The expression for the polarization is given in Eq. (14). It is clear from this formula that the 90° polarization of the emitted photonucleons, due to electric dipole ($E1$) interaction, is zero. The polarization at 45° in the $O^{16}(\gamma, n_0)O^{15}{}_{g.s.}$ and the $O^{16}(\gamma, p_0)N^{15}{}_{g.s.}$ reactions are shown in Fig. 5. Very little experimental information on polarization is available in the high-energy region. However, the experimental data of Bertozzi *et al.*[16] and Cole *et al.*[17] for neutron emission around 30 MeV seem to indicate that the calculated value is of the right magnitude and has the correct sign.

2. *With Ground-State Correlation*

The ground state of O^{16} (or O^{15}) can be described with respect to the C^{12} core and four (or three) par-

[15] D. E. Frederick and A. Daniel Sherick, Phys. Rev. **176**, 1177 (1968).

[16] W. Bertozzi, P. Demos, F. Hanser, S. Kowalski, C. Sargent, W. Turchinetz, R. Fullwood, and J. Russell, *Comptes Rendus du Congrès International de Physique de Nucleaire,* edited by P. Gugenberger (Centre National de la Recherche Scientifique, Paris, 1964), Vol. II, p. 1026.

[17] G. W. Cole, F. W. K. Firk, and T. W. Phillips, Phys. Letters **30B**, 91 (1969).

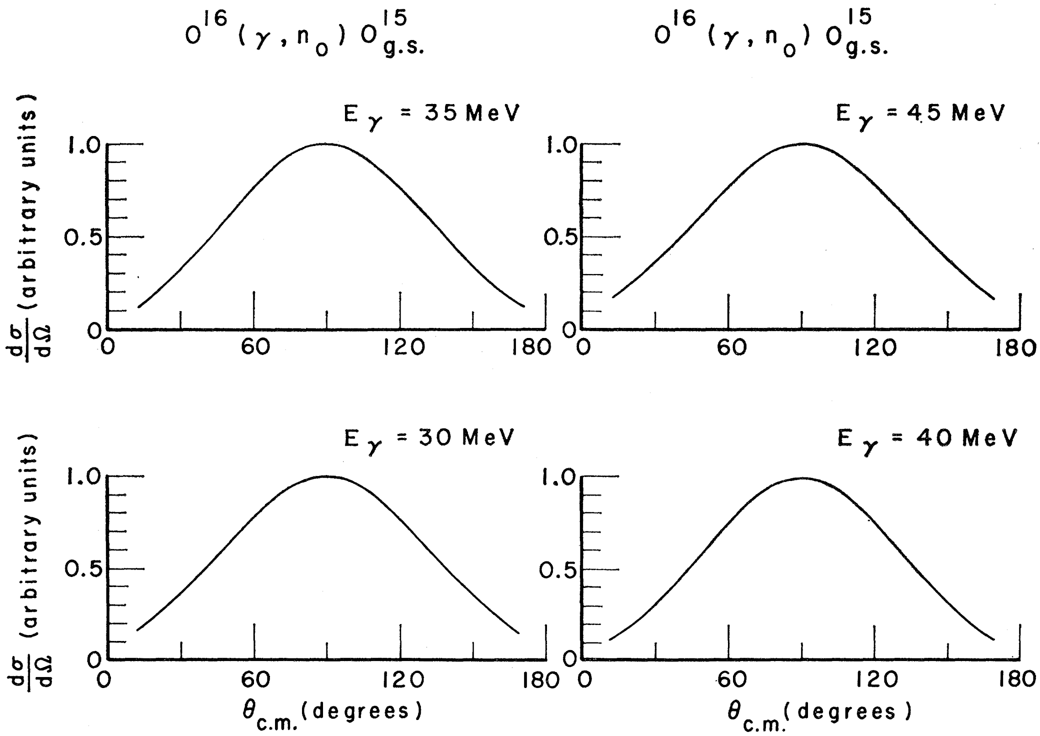

FIG. 3. Angular distribution (normalized to unity at $\theta_{c.m.}=90°$) of photoneutrons from the $O^{16}(\gamma, n_0)O_{g.s.}$ reaction.

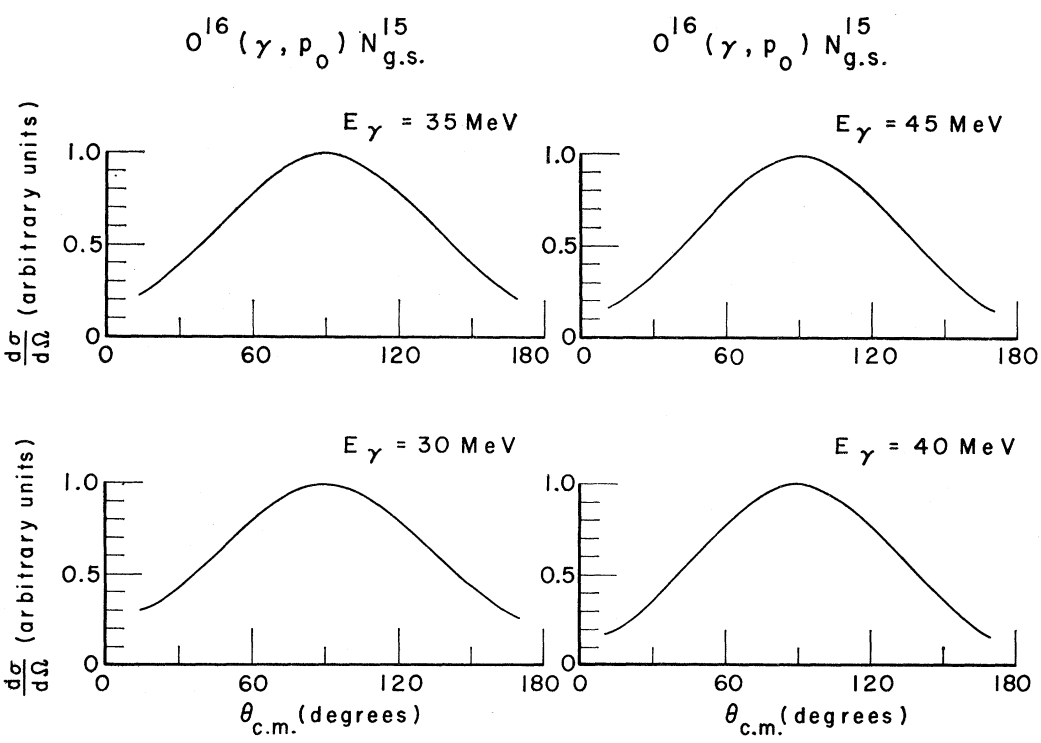

FIG. 4. Angular distribution (normalized to unity at $\theta_{c.m.}=90°$) of photoprotons from the $O^{16}(\gamma, p_0)N^{15}{}_{g.s.}$ reaction.

ticles moving outside that core. Zuker, Buck, and McGrory[18] used such an approach to study the spectrum of O^{16}. We have taken their wave functions to see how it affects the photonucleon yield. The wave functions are

$$\Psi(O^{16}{}_{\text{g.s.}})=[0.71P_{1/2}{}^4+0.58d_{5/2}{}^2P_{1/2}{}^2]_{I^\pi=0^+,T=0}, \quad (16a)$$

$$\Psi(O^{15}{}_{\text{g.s.}})=[0.83P_{1/2}{}^3+0.51d_{5/2}{}^2P_{1/2}{}^1]_{I^\pi=1/2^-,T=1/2}. \quad (16b)$$

In these wave functions, amplitudes smaller than 0.28 have been neglected. With these wave functions the direct interaction cross section is modified as follows:

$$(d\sigma/d\Omega)_{\text{correlated}}=0.78(d\sigma/d\Omega)_{\text{uncorrelated}}. \quad (17)$$

The uncorrelated cross section is given by Eq. (12). The same reduction factor also applies to the integrated cross section. We may therefore conclude that this kind of correlation reduces the magnitude but does not affect the shape of the excitation curve. A comparison between the correlated and uncorrelated excitation function of the reaction $O^{16}(\gamma, n_0)O^{15}{}_{\text{g.s.}}$ is shown in Fig. 6.

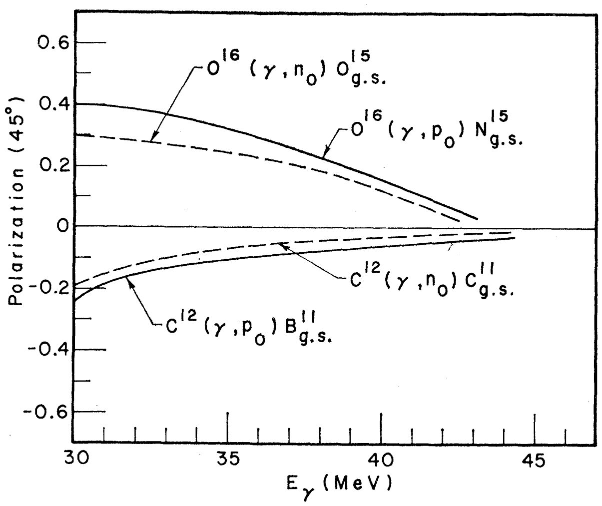

FIG. 5. Photoneutron and photoproton polarization at $\theta=45°$ from the $O^{16}(\gamma, n_0)O^{15}{}_{\text{g.s.}}$ and the $O^{16}(\gamma, p_0)N^{15}{}_{\text{g.s.}}$, and the $C^{12}(\gamma, n_0)C^{11}{}_{\text{g.s.}}$ and the $C^{12}(\gamma, p_0)B^{11}{}_{\text{g.s.}}$ reactions.

To compute the $O^{16}(\gamma, p_0)N^{15}{}_{\text{g.s.}}$ cross section using the ground-state correlation we require the wave function of the N^{15}. From the observed similarity of the low-lying spectra of the N^{15} and O^{15}, it is reasonable to assume that the N^{15} ground-state wave function is very similar to that of O^{15}. Therefore, we have used (16b) to describe $N^{15}{}_{\text{g.s.}}$, and the result is presented in Fig. 6. Again the incorporation of the ground-state correlation is important to get the correct magnitude.

[18] A. P. Zuker, B. Buck, and J. B. McGrory, Phys. Rev. Letters **21**, 39 (1968).

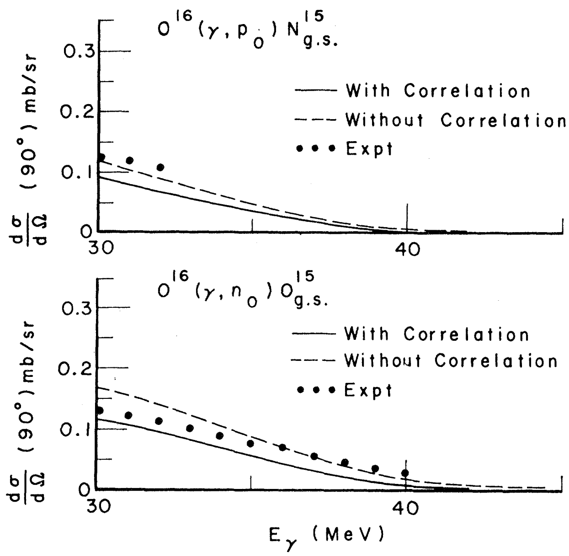

FIG. 6. The 90° differential cross section of the $O^{16}(\gamma, n_0)O^{15}{}_{\text{g.s.}}$ and the $O^{16}(\gamma, p_0)N^{15}{}_{\text{g.s.}}$ reactions with and without the ground-state correlation. The experimental data are from Refs. 1 and 14.

B. C^{12} Target

1. *No Ground-State Correlation*

According to the elementary shell model, there should be single-particle transitions from the $1P_{3/2}$ and the $1S_{1/2}$ bound levels. Since the experimental 90°

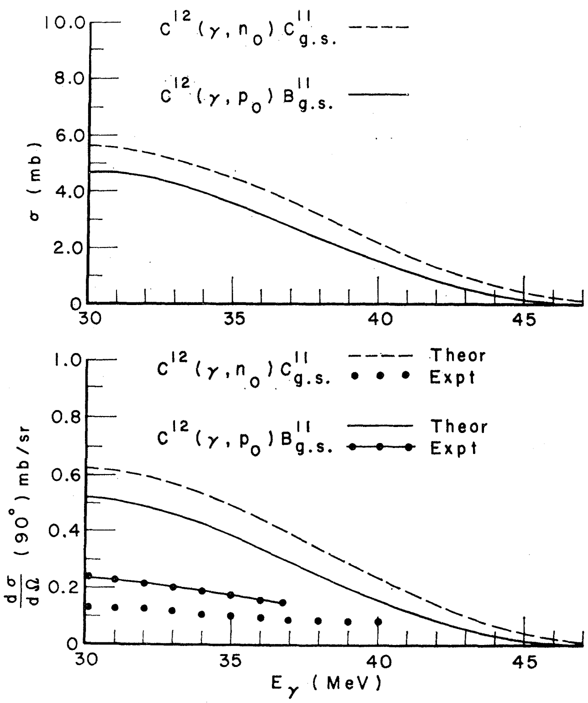

FIG. 7. The 90° differential and the integrated cross sections of the $C^{12}(\gamma, n_0)C^{11}{}_{\text{g.s.}}$ and the $C^{12}(\gamma, p_0)B^{11}{}_{\text{g.s.}}$ reactions. The (γ, n_0) experimental data are from Wu *et al.*,[1] and the (γ, p_0) data are from Barssard *et al.*[19]

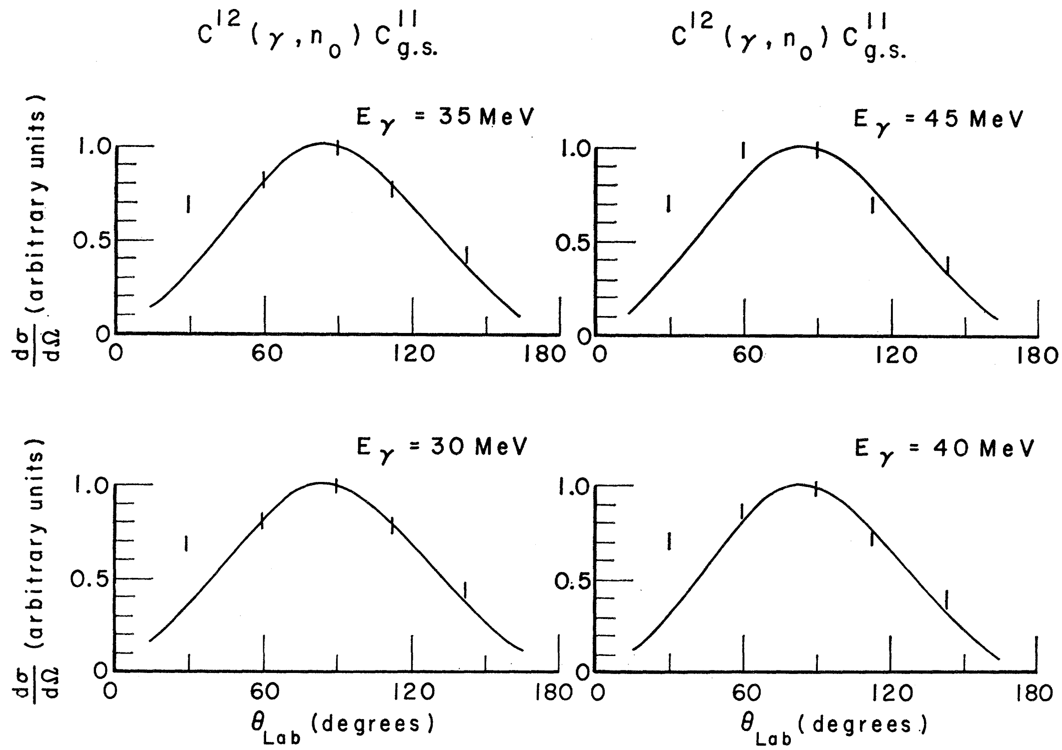

FIG. 8. Angular distribution (normalized to unity at $\theta_{lab}=90°$) of photoneutrons from the $C^{12}(\gamma, n_0)C^{11}{}_{g.s.}$ reaction. The experimental data are from Rawlins *et al.* (Ref. 20).

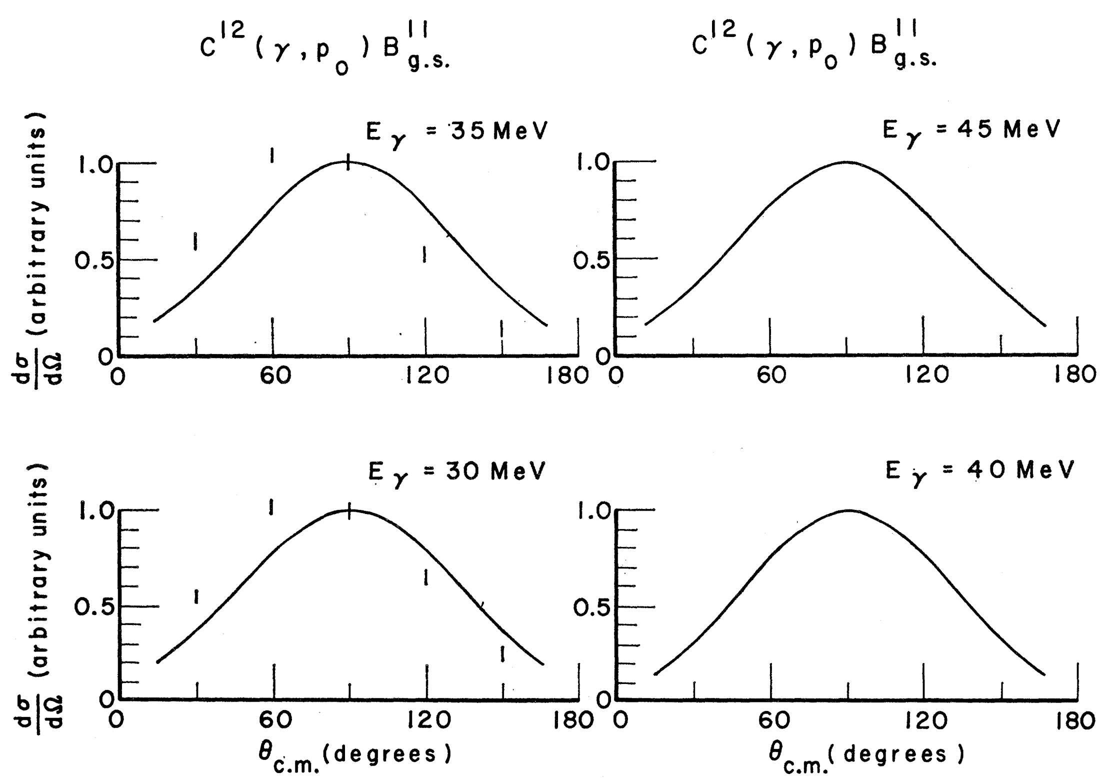

FIG. 9. Angular distribution (normalized to unity at $\theta_{c.m.}=90°$) of photoprotons from the $C^{12}(\gamma, p_0)B^{11}{}_{g.s.}$ reaction. The experimental data are from Brassard *et al.* (Ref. 21).

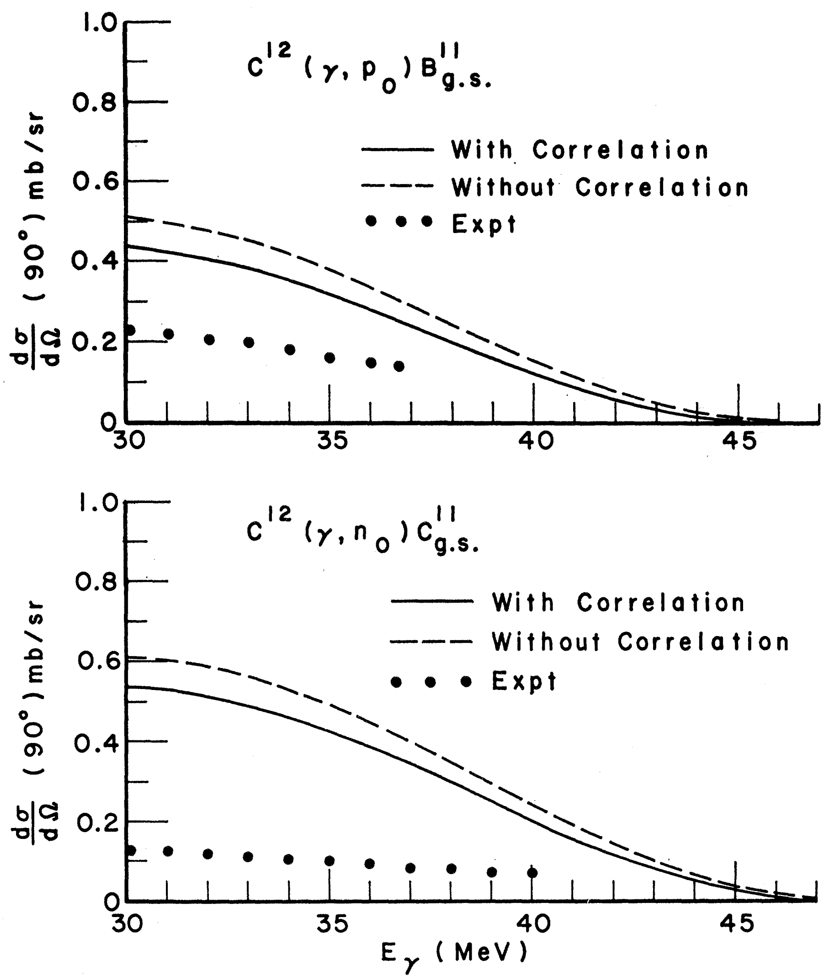

FIG. 10. The 90° differential cross section of the $C^{12}(\gamma, n_0)C^{11}_{g.s.}$ and the $C^{12}(\gamma, p_0)B^{11}_{g.s.}$ reactions with and without the ground-state correlation. The experimental data is from Refs. 1 and 19. The data of Ref. 22 lie somewhat lower.

cross section is available only for the ground-state transitions, we have neglected the transitions from $1S_{1/2}$ level. Hence, for electric dipole interaction, the transitions involved in the calculation are $1P_{3/2}\rightarrow s_{1/2}$, $1P_{3/2}\rightarrow d_{3/2}$, and $1P_{3/2}\rightarrow d_{5/2}$.

We have used the same type of potential as in O^{16} case. However, the potential parameters are slightly different. The binding energy used for the $1P_{3/2}$ level and the potential parameters are given in Table I. The calculated cross sections are shown in Fig. 7. The 90° differential cross section for the (γ, n_0) and the (γ, p_0) reactions are compared with the experimental data of Wu *et al.*[1] and Brassard *et al.*,[19] respectively. We can draw the same conclusion as in the case of O^{16}, that the cross section does not show any structure above the giant resonance and that it falls off smoothly with increasing incident photon energy.

Angular distributions of neutrons and protons in the reactions $C^{12}(\gamma, n_0)C^{11}_{g.s.}$ and $C^{12}(\gamma, p_0)B^{11}_{g.s.}$ are shown in Figs. 8 and 9, respectively. Here also the calculated distributions are symmetrical about 90°. The neutron distributions are compared with the experimental findings of Rawlins *et al.*[20] and the proton distributions with the findings of Brassard *et al.*[21] It is clear that electric dipole interaction cannot account for the observed distributions. However, the experimental results do not separate out the contribution from the cross sections of other multipole interactions such as $M1$ and $E2$. The presence of higher multipole transition is indicated in the analysis of the angular distribution at 22.1-, 25.3-, 30.7-, 38.2-, and 48.2-MeV photon energies done by Penner and Leiss.[22]

The polarization at 45° for the ground-state transition is shown in Fig. 5 and seems to have opposite sign compared to the O^{16} case.

2. *With Ground-State Correlation*

Just like O^{16}, the actual ground-state wave function of C^{12} contains considerable correlation. We have used

[19] C. Brassard, W. Scholz, and D. A. Bromley, J. Phys. Soc. Japan Suppl. **24**, 139 (1968).

[20] J. A. Rowlins, C. Glavina, S. H. Ku, and Y. M. Shin, Nucl. Phys. **A122**, 128 (1968).

[21] C. Brassard, H. D. Shay, J. P. Coffin, W. Scholz, and D. A. Bromley (private communications).

[22] S. Penner and J. E. Leiss, Phys. Rev. **114**, 1101 (1959).

the wave function of Amit and Katz,[23] who allowed the mixture of the $1P_{3/2}$ and the $1P_{1/2}$, in the intermediate-coupling shell-model calculation. Their wave functions are

$$\Psi(C^{12}{}_{g.s.})=[0.90P_{3/2}{}^{8}P_{1/2}{}^{0}+0.42P_{3/2}{}^{6}P_{1/2}{}^{2}]_{I^{\pi}=0^{+},T=0}, \tag{18a}$$

$$\Psi(B^{11}{}_{g.s.})=[0.77P_{3/2}{}^{7}P_{1/2}{}^{0}+0.26P_{3/2}{}^{6}P_{1/2}{}^{1} +0.55P_{3/2}{}^{5}P_{1/2}{}^{2}]_{I^{\pi}=3/2,T=1/2}, \tag{18b}$$

where amplitudes smaller than 0.2 have been neglected. Using such wave functions, the differential cross section reduces to

$$(d\sigma/d\Omega)_{\text{correlated}}=0.85(d\sigma/d\Omega)_{\text{uncorrelated}}. \tag{19}$$

The same relation holds for the integrated cross section. Here again, the cross section is reduced due to correlation. Assuming the ground-state wave functions of C^{11} and B^{11} are very similar, roughly the same reduction in cross section can also be obtained in the $C^{12}(\gamma, n_0)C^{11}{}_{g.s.}$ reaction. The result is shown in Fig. 10. It may be noted that the matrix element involving the $0.42P_{3/2}{}^{6}P_{1/2}{}^{2}$ term of the C^{12} wave function and the $0.26P_{3/2}{}^{6}P_{1/2}{}^{1}$ term of the B^{11} wave function does not contribute to the cross section because it contains an integral of the form

$$\langle[P_{3/2}{}^{6}P_{1/2}{}^{1}]_{I=3/2}\,|\,[P_{3/2}{}^{6}P_{1/2}{}^{1}]_{I=1/2}\rangle=\delta_{I=3/2,I=1/2}=0.$$

For the $C^{12}(\gamma, n_0)C^{11}{}_{g.s.}$ and the $C^{12}(\gamma, p_0)B^{11}{}_{g.s.}$ reactions, the computed cross section with the consideration of the ground-state correction is higher than the observed ones. However, the absolute magnitude of the computed cross section is quite sensitive to the degree of the correlation present in the ground state. For example, in Eq. (18), if one changes the $0.90p_{3/2}{}^{8}p_{1/2}{}^{0}$ component of C^{12} to $0.8p_{3/2}{}^{8}p_{1/2}{}^{0}$, and the $0.77p_{3/2}{}^{7}p_{1/2}{}^{0}$ component of B^{11} to $0.7p_{3/2}{}^{7}p_{1/2}{}^{0}$, then the correlated cross section is reduced to 0.6 times the uncorrelated one, and the computed cross section in the $C^{12}(\gamma, p_0)B^{11}{}_{g.s.}$ is in agreement with experiment. This also lessens the disagreement in the $C^{12}(\gamma, n_0)C^{11}{}_{g.s.}$ reaction. In fact, if the C^{11} ground-state wave function which has been assumed to be the same as that of B^{11}, is somewhat different, the remaining disagreement may be accounted for.

This sensitivity of the cross section to the ground-state correlation clearly suggests that if one can experimentally separate out the pure $E1$ transition in a photonuclear process and if the direct emission predominates, the study of the absolute magnitude of the high-energy photonucleon yield provides an interesting tool to determine the ground-state correlation. This would be extremely important information for the nuclear spectroscopy and structure. A check of the direct emission can be made by sorting out only the $E1$ component and then looking into angular distributions.

[23] D. Amit and A. Katz, Nucl. Phys. **58**, 388 (1964).

IV. CONCLUSION

In the first approximation, both the structure and the yield of the high-energy photonucleons beyond the giant-dipole region are consistent with the simple picture of the direct emission used here. Our conclusion concurs with that of Kaushal, Winhold, Yergin, Medicus, and Auguston.[24] They observed sizeable yields of photoneutrons above 10 MeV in a number of nuclei from Li to U at 67°. However, more experimental data on the angular distributions, sorting out the yield due to the $E1$ transitions only, will be helpful in refining this picture. In evaluating the direct-reaction photoemission, the incorporation of a good bound-state wave function is important. Instead of the bound-state wave functions in a Woods-Saxon well, if a simple harmonic-oscillator wave function is used, the computed yield will be too low and structureless. Because of the presence of r in the matrix element, the tail of the bound-state wave function (which is absent in the pure oscillator case) provides significant contribution. The effect of the ground-state correlation is to reduce the magnitude of the cross section and, in particular, in cases of the $C^{12}(\gamma, p_0)B^{11}{}_{g.s.}$, $C^{12}(\gamma, n_0)C^{11}{}_{g.s.}$, and $O^{16}(\gamma, n_0)O^{15}{}_{g.s.}$ reactions, the incorporation of the ground-state correlation improves the agreement. The actual magnitude of the cross section is sensitive to the extent of the correlation in the wave functions of the target and the residual nucleus, suggesting that the determination of the absolute magnitude of the high-energy photonucleon emission may provide useful information on the composition of these wave functions.

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[24] N. N. Kaushal, E. J. Winhold, P. F. Yergin, H. A. Medicus, and R. H. Augustson, Phys. Rev. **175**, 1330 (1968).