Multinucleon mechanisms in $(\gamma,N)$ and $(\gamma,NN)$ reactions

Jan Ryckebusch, Lars Machenil, Marc Vanderhaeghen, Veerle Van der Sluys and Michel Waroquier

Laboratory for Theoretical Physics
Proeftuinstraat 86
B-9000 Gent
Belgium
(March 31, 2022)

Abstract

The similarities in the experimental indications for multinucleon mechanisms in $(\gamma,p)$ and $(e,e'p)$ processes are pointed out. For both types of reactions, the substantial role of two-nucleon emission processes for transitions to high excitation energies in the residual nucleus is stressed. A microscopic model for the calculation of the two-body knockout contributions to the inclusive $(\gamma,N)$ reaction is presented. It is based on an unfactorized formalism for the calculation of electromagnetically induced two-nucleon emission cross sections. The model is shown to yield a reasonable description of the overall behaviour of the $^{12}\text{C}(\gamma,p)$ and $^{12}\text{C}(\gamma,n)$ data at high excitation energies in the residual nucleus. In the calculations, effects from non-resonant and resonant pion exchange currents are included. Photoabsorption on these currents are predicted to produce the major contributions to the exclusive $^{16}\text{O}(\gamma,n_0)^{15}\text{O}$ process at photonenergies above the pion threshold. Double differential cross sections for photon induced $pp$ and $pm$ emission from $^{16}\text{O}$ are calculated and compared with the data.
I. INTRODUCTION

The apparent success of the one-body picture in explaining the quasi-elastic \((e, e')\) results was put in a different perspective when the longitudinal and transverse response functions were experimentally separated. In the one-body picture the virtual photon is assumed to couple with the individual nucleons in the target nucleus, the whole process exhibiting hardly any medium dependence. Notwithstanding the extensive amount of work which has been devoted to a theoretical understanding of the separated \((e, e')\) data, a full explanation of both response functions in a consistent model has not yet been accomplished [1].

Recent coincidence \((e, e'p)\) measurements have established the significant role played by multinucleon mechanisms in the quasi-elastic [2,3], dip [1] and \(\Delta_{33}\)-excitation domain [3]. The existence of multinucleon mechanisms was evidenced through a rise in the measured transverse \((e, e'p)\) response functions at high missing energies \(E_m=\omega-T_p\) in the \((A-1)\) system [3]. This excess strength has been shown to be unlikely due to rescattering effects, since a comparable rise of the longitudinal strength would then be expected, an effect which has not been observed experimentally. It is worth mentioning that recent measurements of the separate \((e, e'p)\) structure functions indicate that a similar situation seems to occur as for the inclusive \((e, e')\) cross sections: whereas the one-body picture gives a fair account of the complete quasi-elastic \((e, e'p)\) cross sections for reactions in which the residual nucleus is created in a hole state [4], discrepancies turn up when it comes to comparing the separated structure functions [3]. In a recent paper we have illustrated the particular sensitivity of the longitudinal-transverse \((e, e'p)\) response function to multihadron currents related to pion-exchange and \(\Delta_{33}\) excitation [1].

Whereas the gathered evidence for electron scattering reactions proceeding in part via multinucleon components is relatively new, over the years overwhelming evidence for the occurrence of multinucleon components in real photon reactions has been produced. Rather than attempting to be complete we mention some illustrative examples.

(i) Using the tagged photon technique, a Glasgow-Edinburgh-Mainz collaboration succeeded in measuring the \(^{12}\text{C}\) photoproton cross section up to high excitation energies in the residual nucleus [10]. Just as for the \(^{12}\text{C}(e, e'p)\) results of refs. [2,3,4,5], the \(^{12}\text{C}(\gamma, p)\) results of Ref. [10] reported excess strength at high missing energies, which was shown to be unlikely due to one-body knockout from the deep-lying hole states. In the meantime, the findings of Ref. [10] have been confirmed by several independent measurements [11,12].

(ii) Strong indications for the occurrence of multinucleon mechanisms have also been obtained for the exclusive regime. For a long time it has been realized that exclusive \((\gamma, n)\) reactions at high photon energies are good candidates to reveal information about the role of multihadron currents in photoinduced reactions. Only one experiment of this type has been reported up to now. Exclusive \(^{16}\text{O}(\gamma, n_0)^{15}\text{O}(g.s.)\) measurements at MIT-BATES [13] confirmed the similarity of \((\gamma, p_0)\) and \((\gamma, n_0)\) angular cross sections for photon energies ranging from just above the particle emission threshold, which are typically probing the giant resonance region, to the \(\Delta_{33}\)-production region. This is rather surprising given the uncharged nature of the neutron and the fact that the squared ratio of the respective magnetic moments \((\mu_p/\mu_n)^2\) equals 2.13.
(iii) As a third illustrative example of obvious multinucleon components showing up in photonuclear reactions, we mention the high-resolution $^{12}$C($\gamma, p$) measurements of Refs. [12,14] and the high-resolution $^{40}$Ca($\gamma, p$) measurements of Ref. [15]. In these experiments a strong feeding of states with a 2hole-1particle ($2h - 1p$) character has been observed. In a recent publication [16] we have pointed out that the angular cross section for these transitions can be naturally explained by assuming direct proton emission following the absorption on the pion-exchange currents. A similar strong feeding of the $2h - 1p$ states has been observed in the recent $^{12}$C($e, e'p$) measurements at high missing momenta at NIKHEF-Amsterdam [17].

Clearly, photonuclear reactions offer a good testing ground for any model that aims at describing the multihadron mechanisms in electromagnetically induced nucleon emission reactions. In addition, the similarities in some of the qualitative features of ($\gamma, p$) and ($e, e'p$) are so obvious that a combined analysis is likely to result in a better insight in both types of reactions. This procedure might be particularly useful to arrive at a better quantitative understanding of the physics of the dip and the $\Delta_{33}$ region, for which the ($e, e'p$) spectra exhibit similar qualitative features as their ($\gamma, p$) counterparts [5,6].

The multinucleon mechanisms in photonuclear reactions have been customary interpreted in terms of the phenomenological quasideuteron (QD) model [18] in which the photon is assumed to be predominantly absorbed by $np$ pairs. This model gives a natural explanation of the measured excess strength in the ($\gamma, p$) and ($\gamma, n$) spectra at high missing energies. In the QD phenomenology the measured nucleon is not exclusive and is accompanied by an other nucleon which remains either undetected or gets reabsorbed. The quasideuteron phenomenology has been confirmed by double coincidence measurements of the type ($\gamma, pn$) [19]. In line with the predictions of the quasideuteron model, the ($\gamma, pn$) data were shown to scale with the missing momentum $P = p_p + p_n - q_\gamma$, the $P$ dependence being determined by the probability of finding in the target nucleus a $np$ pair with total momentum $|P|$ and zero separation.

In this paper we present a non-relativistic microscopic model for the calculation of cross sections for one and two-nucleon knockout processes following photoabsorption on finite nuclei. Our main focus will be on estimating the effect of two-nucleon emission processes to ($\gamma, N$) cross sections starting from principal grounds. This involves a microscopic model for the photoabsorption mechanism and a treatment of the final state interaction between the escaping nucleons and the residual nucleus. Concerning the final state interaction, we rely on a shell-model approach to deal with the distortions that the struck nucleons undergo in their way out of the nucleus. By doing this we do not have to worry about spurious contributions to the cross sections due to non-orthogonality problems. Within this shell-model framework for the treatment of the final state interaction we explore the relevance of pionic and $\Delta_{33}$ degrees of freedom in inclusive and exclusive photonucleon processes.

The plan of this paper is as follows. In Sect. II the formalism is sketched. This includes a model for the calculation of ($\gamma, NN$) cross sections and their contribution to the ($\gamma, N$) spectra. In Sect. III the numerical results of the ($\gamma, N$) and ($\gamma, NN$) cross sections are presented. In particular, Sect IIIA deals with the contributions stemming from pionic and $\Delta_{33}$ degrees of freedom to exclusive $^{16}$O($\gamma, n_0$)$^{15}$O(g.s.) processes at higher photon energies. In Sect. IIIB we summarize some results of calculations aiming at estimating the influence of two-nucleon knockout on the ($\gamma, N$) processes leaving the residual nucleus in a continuum
state above the two-particle emission threshold. In Sect. IIIC the results of the $^{16}$O$(\gamma,NN)$ calculations are compared with the data. We conclude with a summary and some outlooks in Sect. IV.

II. FORMALISM

In line with the above discussion, photonucleon spectra reflect multinucleon components in both the discrete (exclusive $(\gamma,N)$) and the continuous part of their spectrum. As explained before there are strong indications that the continuum strength can be attributed to $(\gamma, pn)$ processes. Consequently, a model for two-nucleon emission is essential to calculate $(\gamma, N)$ spectra above the two-particle emission threshold.

In this Section we will first sketch an unfactorized model for the calculation of $(\gamma, NN)$ cross sections. This model does account for the distortions which the outgoing nucleon pair undergoes in its way out of the target nucleus and has been described in more detail in Ref. [20]. In the process of calculating the two-nucleon knockout cross sections, the nuclear structure of the target and the residual nucleus reflects itself in the two-hole spectral function. A schematic model for these spectral functions will be presented. Subsequently, by integrating the derived $(\gamma, NN)$ cross sections over one of the nucleon coordinates, we will obtain an expression for the inclusive $(\gamma, N)$ cross section. Lastly, we will elaborate upon the two-body currents on which the initial photoabsorption is assumed to take place.

A. $(\gamma, NN)$ cross sections and two-hole spectral functions

In the laboratory frame, the coincidence angular cross section for a $(\gamma, N_aN_b)$ reaction (Figure 1) is given by ($h=c=1$):

$$\frac{d^4\sigma_{\text{LAB}}}{d\Omega_a d\Omega_b d^2k_a d^2k_b} = \frac{1}{(2\pi)^5} \sum_f \sum_{m_{s_a}m_{s_b}} \frac{k_a^2 k_b^2}{2E_{\gamma}} \frac{1}{2} \sum_\lambda |m_{\gamma}^{fi}|^2 \delta(E_{A-2} + E_a + E_b - E_A - E_{\gamma}),$$  

where the Feynman amplitude $m_{\gamma}^{fi}$ reads:

$$m_{\gamma}^{fi} = \langle \Psi_f^{(A-2)}(E_x); \mathbf{k}_a m_{s_a}; \mathbf{k}_b m_{s_b} | J_\lambda(q_\gamma) | \Psi_0 \rangle.$$  

Here, $f$ is a shorthand notation for all quantum numbers specifying the eigenstates of the $(A-2)$ system and $E_x$ denotes the positive excitation energy measured with respect to the ground state of the residual $(A-2)$ nucleus. In the cross section of Eq. (1) we have summed over the spin projections of the escaping nucleons $(m_{s_a}, m_{s_b})$ and averaged over the initial photon polarization $\lambda$. The sum over $f$ extends over the discrete and the continuum states of the residual $(A-2)$ nucleus. We assume that for the present purposes the target and residual nucleus can be well described with the aid of Slater determinants of the independent-particle model. Further, we discard all effects due the rescattering (multi-step mechanisms). Within these assumptions the main contribution to the $(\gamma, NN)$ cross sections is supposed to come from direct two-nucleon knockout following the photoabsorption on a two-body current. In such a reaction picture the residual nucleus will be created in a 2 hole $(2h)$ state relative to the ground state $|\Psi_0\rangle$ of the target nucleus:
\begin{align}
|\Psi_f^{(A-2)}(E_\pm)| &= \equiv |(hh')^{-1} J_R M_R >
\begin{align}
&= \sum_{m_h m_{h'}} \frac{1}{\sqrt{1 + \delta_{hh'}}} < j_h m_h j_{h'} m_{h'} | J_R M_R > \\
&\times (-1)^{j_h + m_h + j_{h'} + m_{h'}} c_{h - m_h} c_{h' - m_{h'}} | \Psi_0 > . \quad (3)
\end{align}
\end{align}

By analogy with the partial-wave expansion techniques which are commonly employed in a shell-model approach to one-nucleon emission processes and which are extensively described in Ref. [21], a double partial wave expansion has recently been suggested for the two-nucleon emission case [22]. Here, the proper asymptotic behaviour of the A-body wave function \(|\Psi_f >\equiv (hh')^{-1} J_R M_R; k_a m_s a; k_b m_{s b} >\) is determined by:

\begin{align}
\langle \mathbf{r}_1 \sigma_1, \mathbf{r}_2 \sigma_2 | \Psi_f \rangle = \frac{1}{\sqrt{A(A - 1)}} A_{2(A-2)} \chi_{\frac{1}{2}m_{s a}}(\sigma_1) \left( e^{i k_{a} \mathbf{r}_1} + f_{k_{a}}(\theta_a) \frac{e^{i k_{a} \mathbf{r}_1}}{r_1} \right)
\times \chi_{\frac{1}{2}m_{s b}}(\sigma_2) \left( e^{i k_{b} \mathbf{r}_2} + f_{k_{b}}(\theta_b) \frac{e^{i k_{b} \mathbf{r}_2}}{r_2} \right) |(hh')^{-1} J_R M_R > \right), \quad (4)
\end{align}

and is reached through an expansion in terms of the continuum states \(p(\ell j m)\) of the mean-field potential:

\begin{align}
|\Psi_f > &= \sum_{l m j m'} \sum_{J M_1 M_1} \sum_{J M_1 M_1} (4\pi)^2 \frac{\pi}{2M_N \sqrt{k_a k_b}} e^{i (\delta_l + \sigma_l + \delta_{l'} + \sigma_{l'})} Y_{lm}(\Omega_{k_a}) Y_{l'm'}(\Omega_{k_b}) \\
&\times < lm \frac{1}{2} m_{s a} | j m >= < l'm \frac{1}{2} m_{s b} | j' m' > < J R M_R J_1 M_1 | J M > \\
&\times < j m j' m' > < J M_1 > |(hh')^{-1} J_R; (p(\ell j m) p'(\ell' j' m')) J_1; J M > , \quad (5)
\end{align}

where \(\epsilon_a = k_a^2/2M_a, \epsilon_b = k_b^2/2M_b, M_N\) is the nucleon mass, \(\delta_i\) is the central phase shift and \(\sigma_i\) the Coulomb phase shift of the continuum single-particle state \(p\). In Eq. (4) the operator \(A_{2(A-2)}\) makes sure that the total A-body wave function is properly antisymmetrized and \(r_A\) is a measure for the radius of the target nucleus. Asymptotic wave functions of the type \(\Psi_f^{(A-2)}\) refer to a situation in which the detected nucleon pair has the momentum-spin characteristics \((k_a, k_b, m_{s a}, m_{s b})\) and in which the residual nucleus is created in the \(2h\) state \(|(hh')^{-1} J_R M_R >\). The sum over the partial waves \((l, j, m)\) runs over all continuum states of the single-particle mean-field potential at a particular excitation energy \(\epsilon\), the latter being set by the kinetic energy of the escaping nucleon under consideration. The wave function of Eq. (3) has been derived under the following normalization convention for the continuum single-particle states:

\begin{align}
\varphi_{lj}(r, \epsilon) \rightarrow_{r \gg r_d} \left[ \frac{2M_N}{\pi k} \sin(kr - \eta ln(2k r) - \frac{\pi l}{2} + \delta_l + \sigma_l) \right]. \quad (6)
\end{align}

In order to calculate the Feynman amplitude of Eq. (2) it is convenient to expand the nuclear current operator in terms of its multipole components. This is commonly done with the aid of the electric and magnetic transition operators [22, 23]:
\[ J_\lambda(q_\gamma) = -\sqrt{(2\pi)} \sum_{J} i^J \sqrt{2J + 1} \left( T_{J\lambda}^{el}(q_\gamma) + \lambda T_{J\lambda}^{mag}(q_\gamma) \right). \] (7)

Inserting the Eqs. (3) and (7) in the Feynman amplitude (4) and performing some basic manipulations we obtain the following expression:

\[
m_f^{fi} = -\sqrt{S_{hh'}(E_x)} \sqrt{2\pi} \sum_{J} i^J j \sum_{lmjm'lm'j'm'} \sum_{J_1M_1} (4\pi)^2 (-i)^l+l' \frac{\pi}{2M_N \sqrt{k_a k_b}} e^{-i(\delta_l+\delta_l+\delta_{l'}+\delta_{l'})} \\
\times Y_{lm'}(\Omega_{k_b}) Y_{lm}(\Omega_{k_a}) < lm \frac{1}{2} ms | jm < < l'm' \frac{1}{2} ms' | j'm'> \\
\times \langle jm j' m' | J_1 M_1 > \frac{(-1)^{J_R-M_R+1}}{J_1} < J_R - M_R \lambda | J_1 M_1 > \\
\times [ < p(\epsilon_a l j) p'(\epsilon_b l' j'); J_1 || T_{J}^{el}(q_\gamma) + \lambda T_{J}^{mag}(q_\gamma)|| h h'; J_R > \\
- (-1)^{J_R+J_R+J_R} < p(\epsilon_a l j) p'(\epsilon_b l' j'); J_1 || T_{J}^{el}(q_\gamma) + \lambda T_{J}^{mag}(q_\gamma)|| h' h; J_R > ], \] (8)

where the function \( S_{hh'}(E) \) is equal to the joint probability of removing two nucleons remaining in the states \( h \) and \( h' \) from the ground state of the target nucleus and of finding the resulting system (with \( (A-2) \) nucleons) with an excitation energy in the interval \((E,E'+dE)\). The function \( S_{hh'}(E) \) is commonly referred to as the two-hole spectral function \([24]\) and is defined according to:

\[ S_{hh'}(E_x) = \sum_f \left| \langle \Psi_f^{(A-2)}(E_x) | c_h c_{h'} | \Psi_0 \rangle \right|^2. \] (9)

In the actual calculations we have adopted a very schematic model for the two-hole spectral functions. Under the assumption that the removal of a nucleon does not affect the subsequent removal of a second nucleon, the function \( S_{hh'}(E) \) can be approximated by the product of two probabilities: the probability to remove a nucleon in the state \( h \) and create the \((A-1)\) system at an excitation energy \( E' \) and the probability to remove a nucleon in the state \( h' \) from the \((A-1)\) system and create the \((A-2)\) nucleus at an excitation energy \( E \), or formally:

\[ S_{hh'}(E) = \int_{0}^{E} S_h(E') S_{h'}(E-E') dE', \] (10)

where the hole spectral function \( S_h(E) \) is given by:

\[ S_h(E) = \sum_f \left| \langle \Psi_f^{(A-1)}(E) | c_h | \Psi_0 \rangle \right|^2. \] (11)

In the optical model the hole spectral function is distributed according to a Breit-Wigner law, centered on the quasi-particle energy \( | \epsilon_h - \epsilon_F | \) (here, \( \epsilon_F \) denotes the Fermi energy) and with a full width at half maximum given by \( 2W(E) \) \([25]\):

\[ S_h(E) = \frac{1}{\pi} \frac{W(E)}{(E - | \epsilon_h - \epsilon_F |)^2 + (W(E))^2}. \] (12)

Parametrizations for the imaginary part of the optical potential \( W(E) \) are obtained from compilations of experimental data and can e.g. be found in Ref. \([26]\). In this work we have adopted the parametrization by Jeukenne and Mahaux \([26]\):
\[ W(E) = \frac{9E^4}{E^4 + (13.27)^4} \text{ (MeV)}. \] (13)

In Fig. 2 some two-hole spectral functions for \(^{16}\text{O}\) obtained in the outlined model are shown. In line with the results of the quasi-elastic \(^{16}\text{O}(e, e'p)\) measurements regarding the spreading of the hole strength in \(^{15}\text{N}\) [27], the quasi-particle energies \(\epsilon_h - \epsilon_F\) were determined to be 6 MeV for the \(1p_{3/2}\) and 30 MeV for the \(1s_{1/2}\) hole state. For all results of this paper the two-hole spectral functions \(S_{hh'}\) have been renormalized to unity: \(\int dE S_{hh'}(E) = 1\).

**B. \((\gamma, N)\) cross sections**

For many years, the \((\gamma, N)\) process for transitions in which the residual nucleus is created at high excitation energies, has been interpreted as the QD region with an undetected nucleon of opposite nature. Here, we will work out a microscopic model which will put us in the position to calculate cross sections for these inclusive processes starting from principal grounds. In line with the basic assumption of the QD model, we can assume that an important part of the \((\gamma, N_a)\) cross section at excitation energies above the two-particle emission threshold can be attributed to \((\gamma, N_aN_b)\) processes. Another mechanism which could be expected to contribute to \((\gamma, N)\) transitions at high excitation energies in the residual nucleus, is the exclusive process with excitation of the deep lying hole strength. Accordingly, we write the \((\gamma, N)\) cross section above the two-particle emission threshold as the sum of a one-nucleon and a two-nucleon knockout piece:

\[
\frac{d^2\sigma^{LAB}}{d\Omega_a dk_a} = \frac{d^2\sigma^{LAB}}{d\Omega_a dk_a} \bigg|_{[1]} + \frac{d^2\sigma^{LAB}}{d\Omega_a dk_a} \bigg|_{[2]},
\]

where the two-nucleon piece is determined according to:

\[
\frac{d^2\sigma^{LAB}}{d\Omega_a dk_a} \bigg|_{[2]} = \int d\Omega_b \int_0^\infty dk_b \frac{d^3\sigma^{LAB}}{d\Omega_a d\Omega_b dk_a dk_b}(\gamma, N_aN_b).
\]

In the calculation of the two-nucleon knockout contribution, the \((\gamma, N_aN_b)\) cross section is determined within the model outlined in the previous subsection. Since we are working in coordinate space the integration over the solid angle of the undetected nucleon can be performed analytically. After some basic manipulations we find with the aid of the Eq. (8) that:

\[
\frac{d^2\sigma^{LAB}}{d\Omega_a dk_a}(\gamma, N_a) \bigg|_{[2]} = \sum_{hh'} \int dk_b \int dE_x S_{hh'}(E_x) \delta(E_{A-2} + E_a + E_b - E_A - E_\gamma)
\]

\[
\times \sum_{J_1} \sum_{J_2} \sum_{i_1' j_1' j_2'} \sum_{J_1 j_2} \sum_{J_1' j_1'} \sum_{J_2' j_2'} \frac{A - 2k_aE_b}{A} \pi \frac{\pi}{4M_N E_\gamma} (-i)^{J_1 + j_1' + j_2 + J_2'} \hat{J}_1 \hat{j}_1' \hat{J}_2 \hat{j}_2' \hat{J}_1' \hat{j}_1' \hat{J}_2' \hat{j}_2' \hat{J}_1 \hat{j}_1 \hat{J}_2 \hat{j}_2
\]

\[
\times P_{J_2}(\cos \theta_a)(-1)^{j_2 - 1/2 + J_1 + j_1'} \epsilon^{-i(\delta_{j_1' + j_2' - \delta_{j_1' + j_2'}} - \sigma_{j_1' - \sigma_{j_1'}})(1 + (-1)^{I_1 + I_2})}
\]

\[
\times < j' - 1 | J 1 | J_2 0 > < j' 1/2 | j_1' - 1/2 | J_2 0 > \left\{ \begin{array}{c}
J_1 J_2
\end{array} \right\} \left\{ \begin{array}{c}
J_1' J_2'
\end{array} \right\}.
\]
subsection. The single-nucleon knockout contribution to the Eq. (14) is then given by:

\[
\times \left\{ (1 + (-1)^{J+J'+J_2}) \left[ \mathcal{M}_{pp',hh'}^{el}(J_1, J, J_R) \left( \mathcal{M}_{pp',hh'}^{el}(J_1', J', J_R) \right)^* + \mathcal{M}_{pp',hh'}^{mag}(J_1, J, J_R) \left( \mathcal{M}_{pp',hh'}^{mag}(J_1', J', J_R) \right)^* \right] + (1 + (-1)^{J+J'+J_2+1}) \left[ \mathcal{M}_{pp',hh'}^{el}(J_1, J, J_R) \left( \mathcal{M}_{pp',hh'}^{el}(J_1', J', J_R) \right)^* + \mathcal{M}_{pp',hh'}^{mag}(J_1, J, J_R) \left( \mathcal{M}_{pp',hh'}^{mag}(J_1', J', J_R) \right)^* \right] \right\},
\]

where \( \hat{J} \equiv \sqrt{2J+1} \), the \( P_J \) are the familiar Legendre Polynomials of degree \( J \), \( E_x \) is the excitation energy in the \((A-2)\) nucleus and the two-body matrix elements \( \mathcal{M} \) have been defined according to:

\[
\mathcal{M}_{pp',hh'}^{el}(J_1, J, J_R) = \langle p(\epsilon_b l')p'(\epsilon_a l')'; J_1 || T^{el, mag}_J(q_\gamma) || hh'; J_R > -(-1)^{j_h+j'_h+J_R} \langle p(\epsilon_b l) p'(\epsilon_a l')'; J_1 || T^{el, mag}_J(q_\gamma) || hh'; J_R > ,
\]

(17)

where \( \epsilon_a^2 = k_a^2/(2M_N) \).

The one-nucleon knockout contribution to the cross sections (14) is calculated with a coupled-channel continuum RPA technique. For an elaborate description of this model the interested reader is referred to Ref. [23]. In brief, the RPA model involves a coupled-channel calculation for all one-nucleon emission channels ((\( \gamma, p \)) and \(( \gamma, n \)) leaving the residual nucleus in a hole state relative to the ground state of the target nucleus.

An obvious shortcoming of the standard RPA is that it does not account for the spreading of the single-particle hole strength in the residual nucleus. Essentially, in the calculation of the cross sections for a particular reaction channel \( C \) (in the RPA formalism a channel \( C \) is characterized by the quantum numbers of the hole state excited in the residual nucleus and the momentum of the outgoing nucleon \( C(n_h l_h j_h; k'_a, m_m) \)), it is assumed that all hole strength is concentrated in the residual nucleus at an excitation energy \( | \epsilon_h - \epsilon_F | \). Here, \( \epsilon_h \) is the Hartree-Fock single-particle energy of the considered state. In order to account for the spreading of the deep-lying hole strength in the residual nucleus, we have folded the calculated \(( \gamma, N)\) angular cross sections with excitation of particular hole state \( h \) (denoted by \( \frac{d\sigma^{LAB}}{d\Omega_a} \big|_{RPA(h)} \), with the hole spectral function \( S_h \) as defined in the preceding subsection. The single-nucleon knockout contribution to the Eq. (14) is then given by:

\[
\left. \frac{d^2\sigma^{LAB}}{d\Omega_a dk_a} (\gamma, N_a) \right|_{[1]} = \sum_h S_h(E_x) \left. \frac{d\sigma^{LAB}}{d\Omega_a} \right|_{RPA(h)},
\]

(18)

where the sum extends over all occupied single-particle states in the residual nucleus and the excitation energy in the residual nucleus \( E_x \) is determined by \( E_x + S_p = E_M + k_a^2/2M + T_{A-1} \).

C. Absorption mechanisms

The next step is to provide a model for the dominant mechanisms in the photoabsorption process. Since our main focus will be on photon energies below 200 MeV, we assume the photon to couple predominantly with the pion dominated nucleon-nucleon correlations in the target nucleus. These correlations include terms with and without an intermediate \( \Delta_{33} \) excitation, as indicated in Fig. [3] For the terms with no \( \Delta_{33} \) lines we have considered the currents associated with the one-pion exchange potential (OPEP)
\[ V_\pi(k) = -\frac{f_{\pi NN}^2}{m_\pi^2} \frac{1}{m_\pi^2 + \frac{k^2}{2}} (\sigma_1 \cdot k) \cdot (\sigma_2 \cdot k) \cdot \tau_1 \cdot \tau_2, \]  

(19)

where \( m_\pi \) is the pion mass and \( f_{\pi NN} \) the pseudovector pion-nucleon coupling constant, \( f_{\pi NN}^2/(4\pi) = 0.079 \). The one-pion exchange current originating from the coupling of an external electromagnetic field with two nucleons interacting through the potential (19) can be found in many textbooks and is a sum of the seagull (diagram (a)) and the pion-in-flight term (diagram (b)) [28].

In the evaluation of the \( \Delta_{33} \) propagators in the diagrams (c) and (d) we have introduced an energy-dependent \( \Delta_{33} \) decay width \( \Gamma_\Delta \), such that the propagators read

\[ \frac{1}{M_\Delta - M_N - E_\gamma - \frac{i}{2} \Gamma_\Delta(E_\gamma)}, \]

(20)

with \( M_\Delta = 1232 \) MeV. The \( \Delta \)-decay width \( \Gamma_\Delta \) is considered to be exclusively the result of \( \Delta \to \pi + N \) decay and has been determined according to the expression given in Ref. [29] :

\[ \Gamma_\Delta(E_\gamma) \approx \frac{8 f_{\pi NN}^2 (E_\gamma^2 - m_\pi^2)^{3/2}}{12\pi} \frac{(M_\Delta - M_N)}{E_\gamma}. \]

(21)

In coordinate space, the \( \Delta_{33} \)-isobar current corresponding to the diagrams (c) and (d) is then:

\[ J^{(\pi\Delta)}(r, r_1, r_2) = \frac{2 f_{\gamma N \Delta} f_{\pi NN} f_{\pi NN}}{9\pi^2 (E_\Delta - E_\gamma - \frac{i}{2} \Gamma_\Delta(E_\gamma))} \left\{ \left[(\tau_1 \times \tau_2) \cdot \sigma_2 \cdot \nabla_2 (\sigma_1 \times \nabla_2) \times (\nabla_1 + \nabla_2) \right] + 4(\tau_2) \cdot \sigma_2 \cdot \nabla_2 (\nabla_1 \times \nabla_2) \delta(r - r_1) \right\} + \frac{e^{-m_\pi |r_2 - r_1|}}{4\pi |r_2 - r_1|}, \]

(22)

with \( E_\Delta \equiv M_\Delta - M_N \). In the absence of a convincing microscopic theory, we are forced to treat the \( \pi NN \) vertex in a phenomenological way. In this paper we shall resort to the widely used monopole form:

\[ F(\Lambda_\pi, p) = \frac{\Lambda_\pi^2 - m_\pi^2}{\Lambda_\pi^2 + p^2}, \]

(23)

with a cutoff parameter \( \Lambda_\pi = 1.2 \) GeV, a value obtained in investigations in which nucleon-nucleon scattering data are fitted in terms of the Bonn one-boson exchange potential.

### III. RESULTS

All results presented below have been obtained with single-particle wave functions obtained from a Hartree-Fock (HF) calculation with an effective interaction of the Skyrme type : SkE2 [30]. Apart from the bound state wave functions, the HF calculation determines the distorting potential in which the partial waves and phase shifts for the escaping nucleons are calculated. All results presented below were checked not to depend dramatically on this particular choice for generating the mean-field characteristics of the target nucleus. This observation can be understood by considering that the two-body matrix elements are not very sensitive to the high-momentum components in the mean-field wave functions.
A. Calculations for the exclusive $^{16}\text{O}(\gamma,n_0)^{15}\text{O(g.s.)}$ reaction

Figure 5 contains the calculated exclusive $^{16}\text{O}(\gamma,n_0)^{15}\text{N}$ cross sections at three different values of the photon energy, all lying above the pion threshold. The results are obtained in a direct knockout reaction formalism including the diagrams as listed in Fig. 4. Given the uncharged nature of the neutron, the one-body component is restricted to the magnetization current. The two-body components involve the currents of Sect. II C. Inspecting Fig. 5 it is clear that the calculations meet our expectations in the sense that the one-body mechanism represents but a fraction of the measured strength and that the angular cross sections are determined by neutron emission following the absorption on the two-body currents. At forward neutron angles a destructive interference effect between the one-body and the MEC contribution is noticed. Consequently, the cross section is dominated by the resonant $\Delta_{33}$ term at forward neutron angles. Remark further that the $\Delta_{33}$ contribution gains in importance with increasing photon energies. Whereas at $E_\gamma=150$ and 200 MeV a fair description of the data can still be obtained with the one-body and the nonresonant MEC contribution, the $\Delta_{33}$ produces the major contribution at $E_\gamma=250$ MeV. All curves drawn in Fig. 5 are obtained for a spectroscopic factor of 0.5. This corresponds with the ground state in $^{15}\text{O}$ exhausting 50% of the total $1p_{1/2}$ hole strength.

B. $(\gamma,p)$ and $(\gamma,n)$ reactions above the two-particle emission threshold

The results of our model calculations for the $^{12}\text{C}(\gamma,n)$ reaction above the $pn$ emission threshold are summarized in Fig. 6. In all figures of this subsection, the missing energy is defined according to $E_m = E_\gamma - T_N$, with $T_N$ the kinetic energy of the detected nucleon in the LAB system. In Fig. 6 we have plotted the calculated contributions from both one- and two-nucleon knockout. The contribution from exclusive one-neutron knockout (at the considered missing energies dominated by $1s$-shell removal) is calculated in the RPA formalism as outlined in Ref. [23]. In previous papers, it was illustrated that the RPA gives a fairly realistic account of the exclusive $(\gamma,N)$ data below 100 MeV photon energy [14,23]. For the spreading of the single-hole strength in the residual nucleus we used the hole spectral function of Eq. (12) in combination with the parametrization of Eq. (13). In line with the $^{12}\text{C}(p,2p)$ results [31] which find the s-shell knockout strength being distributed in the form of a wide peak between 10 and 30 MeV excitation energy, the quasiparticle energy $|\epsilon_h - \epsilon_F|$ was fixed at 20 MeV. The neutron separation energy $S_n$ being 18.7 MeV, the peak of the s-shell removal strength corresponds with a missing energy of about 39 MeV. Accordingly, the s-shell knockout strength is concentrated just above the $pn$ threshold ($S_{pn}=27.4$ MeV) in the $(\gamma,n)$ spectrum. It should be stressed that the missing-energy dependence of the $1s$ strength is mainly determined by the hole spectral function $S_h$. Whereas, at $E_\gamma=75$ MeV the $1s$ $(\gamma,N)$ strength is still substantial in comparison with the $pn$ strength, it is hardly visible at $E_\gamma=100$ MeV. Generally, for $E_\gamma \geq 100$ MeV the RPA predicts the one-body knockout contributions from the deep-lying hole states to be a negligible fraction of the measured $(\gamma,N)$ strength in the continuum. The calculated $(\gamma,pn)$ contributions are predicted to be substantially larger but are observed not to fully exhaust the measured strength. Nevertheless the calculations seem to account for the overall missing
energy behaviour of the data. This is particularly the case when the \((\gamma, pn)\) contribution is arbitrarily renormalized with a factor of two.

The results of the \(^{12}\text{C}(\gamma, p)\) calculations are shown in Figs. 7, 8 and 9 and compared with unpublished results of a Glasgow-Mainz collaboration [32]. The calculations have been performed at two photon energies, one at each side of the pion production threshold: \(E_\gamma = 123\) and 150 MeV. The missing energy dependencies at different values of the proton angle are drawn in Figs. 6 and 8, whereas in Fig. 9 the full angular cross sections are shown for different values of the proton energy. The proton kinetic energies were chosen such that they span the whole missing-energy spectrum. From Figs. 6 and 8 it emerges that the \((\gamma, pn)\) calculations give a reasonable description of the photoproton spectra, particularly for the backward proton angles. Nevertheless, in line with the \(^{12}\text{C}(\gamma, n)\) findings presented earlier, the \(pn\) calculations tend to underestimate the measured photonucleon cross sections. In particular, this seems to be the case at forward proton angles and large missing energies. The excess strength at higher missing energies, which corresponds with slow detected protons, is likely to be due to other mechanisms, like three-nucleon emission \((S_{ppn}=34.0 \text{ MeV}, S_{pnn}=35.8 \text{ MeV})\) which are totally discarded in our calculations. A striking feature of the data is the considerable amount of experimental strength which is observed for the forward proton angles in the region of the \(pn\) threshold. Right at the threshold this strength is unlikely to be due to two-nucleon knockout. Furthermore, exclusive proton removal from the 1s\(_{1/2}\) shell as calculated in the RPA, was found to represent but a very small fraction of the measured strength for the two considered photon energies.

Regarding the missing-energy behaviour, similar features as for the \(^{12}\text{C}(\gamma, p)\) are found in \(^{40}\text{Ca}\). The missing energy behaviour of the \(^{40}\text{Ca}(\gamma, p)\) cross sections for a fixed value of \(E_\gamma\) is presented in Fig. 10. Obviously, the predicted \((\gamma, pn)\) strength does not account for the experimental strength at forward proton angles, whereas a better description is reached at backward angles. Remark further the considerable amount of measured photoproton strength in the region of the \(pn\) threshold at \(\theta_p = 60^\circ\). In the process of calculating the contribution of exclusive one-nucleon knockout to the \(^{40}\text{Ca}(\gamma, p)\) spectrum, we considered removal from the 1s and 1p shell in addition to the 1d\(_{5/2}\) orbit.

The effect of the final state interaction of the struck nucleons with the residual nucleus has been estimated by doing the \(^{12}\text{C}(\gamma, p)\) calculations at \(E_\gamma = 150 \text{ MeV}\), with a plane wave description for the outgoing nucleon pair and comparing the results with the full distorted-wave cross sections. In the formalism outlined in Sect. II, a plane wave description can be simply achieved by replacing the partial waves \(p(\ell jm)\) by properly normalized spherical Bessel functions. In passing it is worth mentioning that the plane wave (PW) description does no longer guarantee the orthogonality between the initial and final states, such that spurious contributions could enter the cross sections. From Fig. 9 it becomes obvious that the effect of the distortions on the angular cross sections is not too dramatic, but for the region just above the threshold \((T_p=100 \text{ MeV})\). This particular behaviour can be easily explained by considering that from energy-conservation arguments a fast moving proton will be necessarily accompanied by a slow neutron, and therefore one of the outgoing nucleons will be subject to strong interactions with the residual system.

In Fig. 9 the effect of different absorption mechanisms is also studied. From the curves drawn in Fig. 9 it is clear that even at photon energies as low as 150 MeV a considerable fraction of the \(pn\) strength at backward proton angles is related to the resonant terms in the
nuclear current. In passing it is worth mentioning that strictly speaking also the \(pp\) channel could be expected to contribute to the \((\gamma, p)\) spectra. Within our model assumptions, the \(pp\) channel can only be fed through the \(\Delta_{33}\) diagrams of Fig. 3. The other diagrams are closed for two-proton emission since they involve a charge-exchange mechanism. In line with the experimental observations [35], however, the calculated \(pp\) strength represents but a small fraction of the photoabsorption strength emerging in the \(pn\) channel [20]. As will become clear in the forthcoming subsection, this finding even holds in the region of the \(\Delta_{33}\) resonance.

C. Results of the \((\gamma, pp)\) and \((\gamma, pn)\) calculations

From the results presented in previous subsection, it emerged that there are strong indications that at high excitation energies in the residual nucleus the measured \((\gamma, N)\) cross sections should not be interpreted as the result of an exclusive process but reflect substantial two-nucleon knockout contributions. In this sense, the \((\gamma, N)\) spectra above the two-nucleon emission threshold, could be expected to be largely set by the physics of \((\gamma, N_aN_b)\) processes. The latter type of reactions, however, offer some supplementary degrees of freedom which might be worth exploiting in order to reach a better understanding of two-nucleon mechanisms in finite nuclei. At present, a full determination of the fivefold differential cross sections \(d^4\sigma/d\Omega_a dk_a d\Omega_b dk_b\) is clearly at the edge of experimental feasibility. Recently, however, several labs have produced double coincidence data at fixed kinematical conditions for one of the outgoing nucleons [35,36]. To calculate the measured cross sections, we can employ Eq. (16). This expression was derived by integrating the full coincidence \((\gamma, N_aN_b)\) cross section over one of the outgoing nucleon coordinates and produced the predictions for the two-nucleon components in the \((\gamma, N)\) spectra. In order to get some idea regarding the realistic character of the proposed \((\gamma, NN)\) model, it is worth checking its predictions against the data. Here, we present some calculations under the kinematical conditions of the measurements of Ref. [36].

The results of the \(^{16}\text{O}(\gamma, pn)\) and \(^{16}\text{O}(\gamma, pp)\) calculations at \(E_\gamma = 281\) MeV and different values of the proton energy are summarized in Figs. 11 and 12. At this photon energy the cross sections are dominated by the \(\Delta_{33}\) current. The \((\gamma, pn)\) and \((\gamma, pp)\) cross sections are found to exhibit similar characteristics. For the high kinetic energies, the angular cross sections are clearly forwardly peaked. With decreasing proton kinetic energy flatter distributions are obtained. The data compromise the proton energy dependence of the cross section at a fixed proton angle \((\theta_p = 52^\circ)\). The comparison with the data is shown in Fig. 13. At low proton energies, the calculated \((\gamma, pp)\) and \((\gamma, pn)\) clearly underestimate the data. This region is usually interpreted as being dominated by pion production. Our calculations seem to suggest that even at low proton kinetic energies there is a considerable background of direct two-nucleon emission. Apart from the pion knockout, also three and more nucleon knockout processes will preferentially feed the low proton energy domain of the spectrum. For lack of a microscopic theory for three and more nucleon ejection processes, explicit \((\gamma, p\pi)\) measurements will be needed to gain insight into the pionproduction channels. At higher kinetic energies, the calculations give a reasonable account of the \(pp\) and \(pn\) emission channel. The dashed curve for the \((\gamma, pp)\) channel is the cross section obtained with a plane-wave description for the escaping protons. It is clear that the distortion effects from the FSI reduce the peak of the cross section and are substantial in explaining the data.
Remark further how the background of $pp$ strength at low proton energies can be partly ascribed to FSI effects. The $pp$ cross sections are about one order of magnitude smaller than the $pn$ cross sections, a feature which is nicely reproduced by the calculations.

IV. CONCLUSION

Summarizing, we have presented a microscopic study of multinucleon mechanisms in photon-induced nucleon-knockout processes. Our study encompasses both inclusive and exclusive photonucleon processes in addition to two-nucleon knockout reactions.

In the exclusive regime, we have reported on $^{16}\text{O}(\gamma,n_{0})$ results above the pion-production threshold. Here, the predominant role of pionic degrees of freedom, including the $\Delta_{33}$ excitation, in photonucleon processes is striking.

Our main focus has been on the contribution from two-nucleon knockout to the inclusive photonucleon spectra. The description relies on an unfactorized approach to two-nucleon knockout reactions. The fair description of the $^{16}\text{O}(\gamma,pn)$ and $^{16}\text{O}(\gamma,pp)$ data, makes us feel rather confident about the realistic character of the employed two-nucleon knockout formalism. Regarding the $(\gamma,N)$ processes, we have shown that at higher missing energies the $pn$ emission strength largely exceeds the strength related to one-nucleon removal from the deep-lying hole states. We find our model, which accounts for photoabsorption on the resonant and non-resonant pion currents, to give a reasonable description of the general features of the $(\gamma,N)$ spectra at high excitation energies in the residual nucleus. Nevertheless, the calculations tend to selectively underestimate the available $^{12}\text{C}(\gamma,p)$ and $^{12}\text{C}(\gamma,n)$ spectra above the $pn$ threshold. This is particularly the case at forward nucleon angles and higher missing energies. This feature, together with the observation that quite some strength resides in the region of the $pn$ threshold, points towards other mechanisms, besides $pn$ emission, contributing to the photonucleon processes with excitation of the residual nucleus in a continuum state. It would be worth investigating this in more detail, particularly in view of the fact that recent calculations predict the short-range effects to occur mainly at high excitation energies \cite{37,38}.

Finally, we stress that the techniques adopted in this paper can be easily applied to $(e,e'N)$ and $(e,e'NN)$ processes. It is to be hoped that a combined analysis of the $(\gamma,p)$ and $(e,e'p)$ spectra, together with new data from $(\gamma,NN)$ and $(e,e'NN)$ measurements, will lead to a better insight into the nature of the multinucleon mechanisms in electromagnetically induced nucleon knockout. Given their particular sensitivity to multinucleon mechanisms, reactions with real photons will play a substantial role in this program.

ACKNOWLEDGMENTS

One of us (J.R.) is indebted to D. van Neck for useful discussions on two-hole spectral functions. The authors would like to thank P. Harty and C. Van den Abeele for giving us the permission to show their data prior to publication. We are also indebted to K. Heyde for valuable discussions and a careful reading of the manuscript. This work was supported by the National Fund for Scientific Research and in part by the NATO through the research grant NATO-CRG920171.
REFERENCES

[1] L.B. Weinstein and W. Bertozzi, in Proc. of the Fourth Workshop on Perspectives in Nuclear Physics at Intermediate Energies (Trieste, 1988), Eds. S. Boffi, C. Ciofi degli Atti and M. Giannini, (World Scientific, Singapore, 1989).
[2] P.E. Ulmer et al., Phys. Rev. Lett. 59, 2259 (1987).
[3] L.B. Weinstein et al., Phys. Rev. Lett. 64, 1646 (1990).
[4] R.W. Lourie et al., Phys. Rev. C56, 2364 (1996).
[5] H. Baghaei et al., Phys. Rev. C39, 177 (1989).
[6] T. Takaki, Phys. Rev. C39, 359 (1989).
[7] A.E.L. Dieperink and P.K.A. de Witt Huberts, Annu. Rev. Nucl. Part. Sci. 40, 239 (1990).
[8] C.M. Spaltro, H.P. Blok, E. Jans, L. Lapikãs, M. van der Schaar, G. van der Steenhoven and P.K.A. de Witt Huberts, NIKHEF-K preprint.
[9] V. Van der Sluys, J. Ryckebusch and M. Waroquier, submitted for publication.
[10] J.C. McGeorge, G.I. Crawford, R.O. Owens, M.R. Sené, D. Branford, A.C. Shotter, B. Schoch, R. Beck, P. Jennewein, F. Klein, J. Vogt and F. Zettl, Phys. Lett. B179, 212 (1986).
[11] P.D. Harty, M.N. Thompson, G.J. O’Keefe, R.P. Rassool, K. Mori, Y. Fujii, T. Suda, I. Nomura, O. Konno, T. Terasawa and Y. Torizuka, Phys. Rev. C37, 13 (1988).
[12] L. Van Hoorebeke et al., Phys. Rev. C42, R1179 (1990).
[13] E.J. Beise, G. Dodson, M. Garçon, S. Hoibraten, C. Maher, L.D. Pham, R.P. Redwine, W. Sapp, K.E. Wilson and S.A. Wood, Phys. Rev. Lett. 62, 2593 (1989).
[14] S.V. Springham, D. Branford, T. Davinson, A.C. Shotter, J.C. McGeorge, J.D. Kellie, S.J. Kellie, S.J. Hall, R. Beck, P. Jennewein and B. Schoch, Nucl. Phys. A517, 93 (1990).
[15] C. Van den Abeele et al., Phys. Lett. B296, 302 (1992).
[16] J. Ryckebusch, K. Heyde, L. Machenil, D. Ryckbosch, M. Vanderhaeghen and M. Waroquier, Phys. Rev. C46, R829 (1992).
[17] L. Kester, Ph.D. thesis, University of Amsterdam (1993).
[18] J.S. Levinger, Phys. Rev. 84, 43 (1951).
[19] I.J.D. McGregor et al., Nucl. Phys. A533, 269 (1991).
[20] J. Ryckebusch, M. Vanderhaeghen, L. Machenil and M. Waroquier, Nucl. Phys. A
[21] C. Mahaux and H. Weidenmüller, in A shell model approach to nuclear reactions (North Holland, Amsterdam, 1969).
[22] T. De Forest and J.D. Walecka, Adv. Phys. 15, 1 (1966).
[23] J. Ryckebusch, M. Waroquier, K. Heyde, J. Moreau and D. Ryckbosch, Nucl. Phys. A476, 237 (1988).
[24] W. Kratschmer, Nucl. Phys. A298, 477 (1978).
[25] C. Mahaux, P.F. Bortignon, R.A. Broglia and C.H. Dasso, Phys. Rep. 120, 1 (1985).
[26] J.P. Jeukenne and C. Mahaux, Nucl. Phys A394, 445 (1983).
[27] S. Frullani and J. Mougey, Adv. Nucl. Phys. 14, 1 (1984).
[28] I.S. Towner, Phys. Rep. 155, 263 (1987).
[29] E. Oset, H. Toki and W. Weise, Phys. Rep. 83, 281 (1982).
[30] M. Waroquier, J. Ryckebusch, J. Moreau, K. Heyde, N. Blasi, S.Y. van der Werf and G. Wenes Phys. Rep. 148, 249 (1987).
[31] G. Jacob and Th.A.J. Marius, Rev. Mod. Phys. 45, 6 (1973).
[32] P.D. Harty, I.J.D. MacGregor, J.C. McGeorge, S.N. Dancer and R.O. Owens, Phys. Rev. C 47, 2185 (1993).
[33] P. Harty, private communication
[34] C. Van den Abeele, private communication.
[35] S.M. Doran, I.J.D. MacGregor, J.R.M. Anand, I. Anthony, S.N. Dancer, S.J. Hall, J.D. Kellie, J.C. McGeorge, G.J. Miller, R.O. Owens, P.A. Wallace, B. Schoch, H. Schmieden and S. Klein, Nucl. Phys. A559, 347 (1993).
[36] J. Arends, P. Detemple, N. Floss, S. Huthmacher, G. Kaul, B. Mecking, G. Noldike and R. Stenz, Nucl. Phys. A526, 479 (1991).
[37] C. Ciofi degli Atti, S. Simula, L.L. Frankfurt and M.I. Strikman, Phys. Rev. C 44, R7 (1991).
[38] H. Mütter and W.H. Dickhoff, preprint.
FIGURES

FIG. 1. Kinematics for the \((\gamma, N_aN_b)\) reaction. The figure sketches the situation in which the photon and the escaping nucleons remain in one plane (planar kinematics).

FIG. 2. Two-hole spectral functions for \(^{16}\text{O}\) as calculated with the schematic model outlined in the text. The two-hole spectral functions have been normalized to unity.

FIG. 3. Feynman diagrams included in the evaluation of the two-body matrix elements.

FIG. 4. Diagrammatic representation of an exclusive \((\gamma, N)\) process with one- and two-body absorption mechanisms in a direct knockout picture.

FIG. 5. Calculated \(^{16}\text{O}(\gamma,n_0)^{15}\text{O}\) \((\text{g.s., } (1p_{1/2})^{-1})\) angular cross sections in a direct knockout model at three values of the photon energy. Dotted line: photoabsorption on the magnetization current. Dashed line: photoabsorption on the magnetization and pion-exchange current. Solid line: photoabsorption on the magnetization, pion-exchange and \(\Delta_{33}\)-isobar current. The data are from Ref. [13].

FIG. 6. Missing energy dependence of the \(^{12}\text{C}(\gamma,n)\) cross section at \(\theta_n=66^\circ\). The dotted line shows the calculated cross sections for one-body knockout from the 1s\(_{1/2}\) shell. The dashed line represent the contribution from \((\gamma,pm)\). The solid line gives the sum of both contributions. For the dot-dashed line the \((\gamma,pn)\) contribution has been arbitrarily multiplied with a factor of two. The data are from Ref. [11].

FIG. 7. Missing energy dependence of the \(^{12}\text{C}(\gamma,p)\) cross section at \(E_{\gamma}=123\text{ MeV}\). The solid line shows the prediction of the \((\gamma,pm)\) calculations. The data are from Ref. [33].

FIG. 8. As in Fig. 7 but at \(E_{\gamma}=150\text{ MeV}\).

FIG. 9. Angular \(^{12}\text{C}(\gamma,p)\) cross sections at different values of the missing energy for \(E_{\gamma}=150\text{ MeV}\). The solid line is the prediction of a \((\gamma,pm)\) calculation with all MEC and \(\Delta_{33}\) diagrams of Fig. 3. For the dotted line only the MEC diagrams are accounted for. The dot-dashed line is the equivalent of the solid curve but is obtained with a plane wave description for the outgoing nucleons.

FIG. 10. Missing energy behaviour of the \(^{40}\text{Ca}(\gamma,p)\) cross section at \(E_{\gamma}=60\text{ MeV}\). The data are from Ref. [34]. The dot-dashed shows the calculated contribution from one-proton knockout. The dashed line represents the contribution from \((\gamma,pm)\). The solid line gives the sum of both contributions.
FIG. 11. $^{16}\text{O}(\gamma, pn)$ angular cross sections at $E_\gamma=281$ MeV and several values of the proton kinetic energy.

FIG. 12. As in Fig. 11 but for the $^{16}\text{O}(\gamma, pp)$ process.

FIG. 13. Proton energy dependence of the $^{16}\text{O}(\gamma, p)$, $^{16}\text{O}(\gamma, pn)$ and $^{16}\text{O}(\gamma, pp)$ reaction at $E_\gamma=281$ MeV and $\theta_p=52^\circ$. In the upper figure ($^{16}\text{O}(\gamma, p)$) both the calculated contribution from $^{16}\text{O}(\gamma, pp)$ (dotted line), $^{16}\text{O}(\gamma, pn)$ (dot-dashed line) and their sum (solid line) are shown. For the $(\gamma, pp)$ channel the dashed curve gives the plane-wave result. The data are from Ref. [36].
This figure "fig1-1.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig2-1.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig3-1.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig4-1.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig1-2.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig2-2.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig3-2.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig4-2.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig1-3.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig2-3.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig3-3.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig4-3.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1
This figure "fig3-4.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9310027v1