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Photoinduced two-proton knockout and ground-state correlations in nuclei

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Abstract

A factorized and analytical form for the $A(\gamma, pp)$ and $A(e, e'pp)$ cross section is proposed. In the suggested scheme the two-proton knockout cross sections can be directly analyzed in terms of the ground-state correlation functions. Central, spin-spin and tensor correlations are considered. In the longitudinal channel, the $(e, e'pp)$ cross section is predicted to exhibit a peculiar sensitivity to ground-state correlation effects.

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The nuclear shell-model is a well-established theory for understanding the structure of atomic nuclei. Despite its obvious success, for a long time it has been realized that the strong repulsive nature of the nuclear force at short internucleon distances is likely to be at the origin of nuclear effects that are incompatible with the independent-particle nature of the shell-model. Obvious signs for short-range effects have recently been obtained from quasi-elastic $(e, e'p)$ reactions. Extensive programs at several high-duty electron facilities predict values for the spectroscopic strength in the low-energy part of the residual-nucleus spectrum that are consistently lower than what was expected in the independent-particle model (IPM). This deviation has been the subject of several investigations, most of them pointing to short- and long-range correlations to be at the origin of the “missing” spectroscopic strength [1,2]. In analyzing and interpreting the quasi-elastic $(e, e'p)$ data it turned out to be extremely useful to make use of the following factorized form of the differential cross section [3]:

$$\frac{d^4\sigma}{d\epsilon' d\Omega_{e'} d\Omega_p dT_p} = E_p \sigma_{ep} \rho(|\mathbf{p}_m|, E_x), \quad (1)$$

where $E(p)$ is the energy (momentum) of the detected proton and σ_{ep} the elementary cross section for electron scattering on an off-shell proton. The spectral function $\rho(|\mathbf{p}_m|, E_x)$ is related to the probability of removing a nucleon with momentum \mathbf{p}_m from the target nucleus and finding the residual nucleus at an excitation energy E_x . Strictly speaking, the above factorized form for the $(e, e'p)$ cross section is only valid in the plane wave

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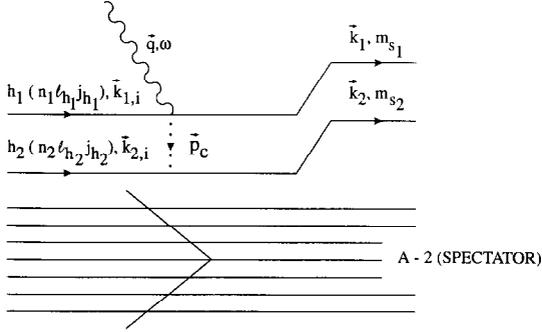


Fig. 1. The emission of two-protons from a nucleus via (virtual) photon absorption.

impulse approximation (PWIA) [3], which puts aside effects like final state interactions (FSI) and photon absorption on two-body currents.

A better understanding of the short-range correlations in nuclei is believed to come from a profound study of (γ, pp) and $(e, e'pp)$ reactions. The underlying idea is that reactions with an electromagnetic probe which induce two particles to escape are extremely sensitive to the two-body dynamics inside the target nucleus. When the main purpose of the two-nucleon knockout studies is the short-range part of the two-nucleon dynamics, proton-proton knockout is to be preferred above proton-neutron knockout. Indeed, the latter are also sensitive to correlations of the charge-exchange type (the most important one being one-pion exchange) that can heavily mask the effects of short-range nature [4]. With two protons in the final state one has an enormous freedom when it comes to determining the position of the hadron detectors. Therefore, one could benefit from a simplified form for the cross section along the lines of Eq. (1). For this simplified form to be useful, one expects it to have some predictive power so that it can be used to map the main sensitivities of the cross sections and optimize the kinematical conditions. In this paper we aim at deriving such a factorized form for the $(e, e'pp)$ and (γ, pp) cross section. Factorized expressions for the (γ, pn) reaction have been derived in Refs. [5,6]. The effect of factorization was studied in Ref. [7] and found to be a reasonable approximation at higher photon energies ($\omega > 150$ MeV). Therefore, at moderate and larger values of the momentum transfer, it is to be expected that a factorized form for the $(e, e'pp)$ cross section provides reasonable estimates.

In order to arrive at a factorized form of a coincidence differential cross section it is a common procedure to work in a plane-wave model for the outgoing nucleons. Furthermore, we adopt the spectator approximation which means that the residual $A - 2$ nucleons are assumed not to participate actively in the reaction process. Similar assumptions are at the basis of the factorized form (1) for the $(e, e'p)$ cross section. In the two-nucleon knockout case, however, the above assumptions are not sufficient for the cross section to factorize. This can best be explained by considering the two-nucleon knockout process as schematically depicted in Fig. 1. There we assume that the photon couples to a correlated pair of protons. The center-of-mass (c.o.m.) momentum \mathbf{P} of the pair can be straightforwardly derived and reads $\mathbf{P} = \mathbf{k}_{1,i} + \mathbf{k}_{2,i} = \mathbf{k}_1 + \mathbf{k}_2 - \mathbf{q}$. The relative momentum $\mathbf{p}_{\text{rel}} = (\mathbf{k}_{1,i} - \mathbf{k}_{2,i})/2$ on the other hand, will generally depend on the momentum \mathbf{p}_C that is exchanged between the two correlated nucleons. Thus, in its most general form the cross section involves an integral over the \mathbf{p}_C which prevents it from being written in a factorized form. Now, the key approximation for the two-nucleon knockout cross section to factorize is that $\mathbf{p}_C \approx \mathbf{0}$. This assumption is related to the “quasi-deuteron approximation” which adopts the view that dinucleons in finite nuclei are mainly residing in relative S states [7,8]. Let $\mathbf{p}_C = \mathbf{0}$, then the relative momentum of the pair is determined by

$$\mathbf{p}_{\text{rel}} = \mathbf{k}_{\pm} = \frac{\mathbf{k}_1 - \mathbf{k}_2}{2} \pm \frac{\mathbf{q}}{2},$$

respectively corresponding with the situation that the photon is absorbed on proton “1” (\mathbf{k}_-) and “2” (\mathbf{k}_+). With all of the above assumptions the $(e, e'pp)$ cross section can be written in a factorized form:

$$\frac{d^5\sigma}{d\epsilon' d\Omega_{e'} d\Omega_1 d\Omega_2 dT_{p_2}} = E_1 p_1 E_2 p_2 \sigma_{\text{ep}}(k_+, k_-, q) F_{h_1, h_2}(P), \quad (2)$$

with $k_{\pm} = |\mathbf{k}_{\pm}|$ and $F_{h_1, h_2}(P)$ denoting the probability to find a nucleon pair in the single-particle states $h_1(n_1 l_1)$ and $h_2(n_2 l_2)$ with center-of-mass (c.o.m.) momentum P . In the IPM this function reads

$$F_{h_1, h_2}(P) = \sum_{m_1, m_2} \left| \int d\mathbf{R} e^{i\mathbf{P}\cdot\mathbf{R}} \phi_{n_1 l_1 m_1}(\mathbf{R}) \phi_{n_2 l_2 m_2}(\mathbf{R}) \right|^2, \quad (3)$$

with ϕ_{nlm} the single-particle wave function for the orbital (nlm). The elementary cross section σ_{ep} can be considered as the equivalent of the σ_{ep} in Eq. (1) and describes the physics of virtual photoabsorption on a diproton embedded in the target nucleus. Generally, σ_{ep} will depend on the photoabsorption mechanisms and the relative motion of the pair. It is precisely in the relative motion that the highest sensitivity to the correlation effects could be expected. We outline a method to derive an analytical expression for σ_{ep} . In doing this, we account for dinucleon correlations that go beyond the IPM.

Adopting the PWA for the outgoing particles one has to evaluate the following type of matrix elements when calculating two-nucleon knockout cross sections

$$M_{m_{s_1} m_{s_2}}^{h_1 h_2 J_R M_R} = \int d\mathbf{r}_1 \dots \int d\mathbf{r}_A \mathcal{A} \left(e^{-ik_1 \cdot \mathbf{r}_1} e^{-ik_2 \cdot \mathbf{r}_2} \Psi_{A-2}^{h_1 h_2 J_R M_R}(3, \dots, A) \right) \\ \times \left\langle 1/2m_{s_1}, 1/2m_{s_2} \left| \left(\sum_{i=1}^A J_{\mu}^{[11]}(i) + \sum_{i<j=1}^A J_{\mu}^{[21]}(i, j) \right) \right| \Psi_i(1, 2, \dots, A) \right\rangle, \quad (4)$$

where \mathcal{A} is the anti-symmetrization operator and m_{s_1} (m_{s_2}) the spin of escaping particle 1 (2). The wave function for the residual nucleus $\Psi_{A-2}^{h_1 h_2 J_R M_R}$ is the two-hole state that is created after knocking two protons out of the target nucleus. The operators $J_{\mu}^{[11]}$ and $J_{\mu}^{[21]}$ are the one- and two-body parts of the nuclear current. As proton-proton emission is not sensitive to charge-exchange, the major component of the two-body operator is the isobaric current. The $J_{\mu=0}^{[11]}$ is related to the charge density operator $\sum_i G_E(q^{\mu} q_{\mu}) e_i \delta(\mathbf{r} - \mathbf{r}_i)$. The transverse components of the one-body current $J_{\mu}^{[11]}(\mu = \pm 1)$ is determined by the convection and magnetization current. The target wave function in the above expression is written as

$$\Psi_i(1, 2, \dots, A) = F(1, 2, \dots, A) \psi_i(1, 2, \dots, A), \quad (5)$$

where ψ_i is the IPM wave function and the operator F induces the correlations. In general, the operator F has many components, reflecting the full complexity of the nucleon-nucleon interaction. Calculations, however, have shown that the major correlation effects can be incorporated by considering an operator of the form [9,10]

$$F(1, 2, \dots, A) = \mathcal{S} \prod_{i<j=1}^A [f_C(r_{ij}) + (f_{\sigma\sigma}(r_{ij}) \boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_j + f_{T\tau}(r_{ij}) S_{ij}) \boldsymbol{\tau}_i \cdot \boldsymbol{\tau}_j]. \quad (6)$$

The first term accounts for central short-range correlations (commonly referred to as Jastrow correlations), whereas the other two induce spin-spin ($\boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_j$) and tensor (S_{ij}) correlations. The operator \mathcal{S} is the symmetrization operator.

In determining σ_{ep} we sum over the spins of the escaping nucleons. Further, instead of considering the contribution to each individual state we evaluate the cross section for emission out of a particular shell-model combination ($n_1 l_{h_1}, n_2 l_{h_2}$). This means that we compute the averaged integrated cross section for a range of excitation energies in the $A - 2$ system. The range of excitation energies will be rather narrow when hole states close to the Fermi level are probed and grow wider as one or both nucleons are escaping from a deeper lying shell.

One can expand the matrix element (4) in terms of $g \equiv 1 - f_C$, $f_{\sigma T}$ and $f_{\tau T}$. In doing this we have adopted the so-called “single-pair approximation” (SPA) [11]. This procedure is equivalent with an expansion into first order in g , $f_{\sigma T}$ and $f_{\tau T}$, retaining only those terms that contain the coordinates of both active nucleons. The SPA was earlier applied in Ref. [5] in the context of (γ, pn) reactions. Physically, the SPA is equivalent with multiplying the IPM relative wave function of the active pair with the correlation operator of Eq. (6). With all these assumptions one obtains the following general expression for σ_{epn} :

$$\sigma_{\text{epn}} = \sigma_M f_{\text{rec}}^{-1} \left[\frac{q_\mu^4}{q^4} w_L + \left(\frac{-q_\mu^2}{2q^2} + \tan^2 \frac{\theta_e}{2} \right) w_T + \frac{q_\mu^2}{2q^2} w_{TT} + \frac{1}{\sqrt{2}} \frac{q_\mu^2}{q^3} (\epsilon + \epsilon') \tan \frac{\theta_e}{2} w_{LT} \right], \quad (7)$$

where f_{rec} is the recoil factor and σ_M the Mott cross section. The w 's depend on the current operator and the different terms in the correlation operator (6). After lengthy calculations one arrives at the following analytical expressions:

$$\begin{aligned} w_L &= 4e^2 (g(k_+) + g(k_-))^2 (G_E(q_\mu q^\mu))^2 + 40e^2 (f_{\sigma T}(k_+) + f_{\sigma T}(k_-))^2 (G_E(q_\mu q^\mu))^2 \\ &\quad + 24e^2 (g(k_+) + g(k_-)) (f_{\sigma T}(k_+) + f_{\sigma T}(k_-)) (G_E(q_\mu q^\mu))^2 \\ &\quad + \frac{16}{3} \sqrt{\frac{\pi}{5}} e^2 (g(k_+) + g(k_-)) (f_{\tau T}^0(-\mathbf{k}_+) + f_{\tau T}^0(-\mathbf{k}_-)) (G_E(q_\mu q^\mu))^2, \\ w_T &= \frac{\mu_p^2 e^2 q^2}{M_p^2} (g(k_+) - g(k_-))^2 (G_E(q_\mu q^\mu))^2 \\ &\quad + \frac{e^2}{2M_p^2} [(k_{1,x}g(k_-) + k_{2,x}g(k_+))^2 + (k_{1,y}g(k_-) + k_{2,y}g(k_+))^2] (G_E(q_\mu q^\mu))^2 \\ &\quad + \frac{256}{81} \left(\frac{f_{\gamma N \Delta} f_{\pi N \Delta} f_{\pi NN}}{m_\pi^3} \right)^2 G_\Delta^2 (G_E(q_\mu q^\mu))^2 \left(\mathbf{q} \times \left(\frac{\mathbf{k}_1 - \mathbf{k}_2}{2} \right) \right)^2 \\ &\quad \times \left[k_+^2 \left(\frac{1}{k_+^2 + m_\pi^2} \right)^2 + k_-^2 \left(\frac{1}{k_-^2 + m_\pi^2} \right)^2 - 2\mathbf{k}_+ \cdot \mathbf{k}_- \frac{1}{k_+^2 + m_\pi^2} \frac{1}{k_-^2 + m_\pi^2} \right], \\ w_{LT} &= 0, \\ w_{TT} &= \frac{-2\mu_p^2 e^2 q^2}{M_p^2} (g(k_+) - g(k_-))^2 (G_E(q_\mu q^\mu))^2 \\ &\quad - \frac{e^2}{M_p^2} [(k_{1,x}g(k_-) + k_{2,x}g(k_+))^2 - (k_{1,y}g(k_-) + k_{2,y}g(k_+))^2] (G_E(q_\mu q^\mu))^2 \\ &\quad - \frac{256}{81} \left(\frac{f_{\gamma N \Delta} f_{\pi N \Delta} f_{\pi NN}}{m_\pi^3} \right)^2 G_\Delta^2 (G_E(q_\mu q^\mu))^2 \left[\left(\mathbf{q} \times \left(\frac{\mathbf{k}_1 - \mathbf{k}_2}{2} \right) \right)_x^2 - \left(\mathbf{q} \times \left(\frac{\mathbf{k}_1 - \mathbf{k}_2}{2} \right) \right)_y^2 \right] \\ &\quad \times \left[k_+^2 \left(\frac{1}{k_+^2 + m_\pi^2} \right)^2 + k_-^2 \left(\frac{1}{k_-^2 + m_\pi^2} \right)^2 - 2\mathbf{k}_+ \cdot \mathbf{k}_- \frac{1}{k_+^2 + m_\pi^2} \frac{1}{k_-^2 + m_\pi^2} \right]. \end{aligned} \quad (8)$$

It should be noted that the tensor and spin-spin correlations have only been considered for the longitudinal channel w_L . We have defined the xz plane as the electron scattering plane with \mathbf{q} along the z axis. In deriving the isobaric current contribution to the above expression we have used the operator specified in Ref. [12] from which also all the coupling constants are taken. The G_Δ is the Δ propagator in which an energy-dependent Δ_{33} has been introduced [12]. The functions $g(k)$, $f_{\sigma T}(k)$ and $f_{\tau T}(k)$ occurring in the above expression are the Fourier transforms of the central functions occurring in the correlation operator (6):

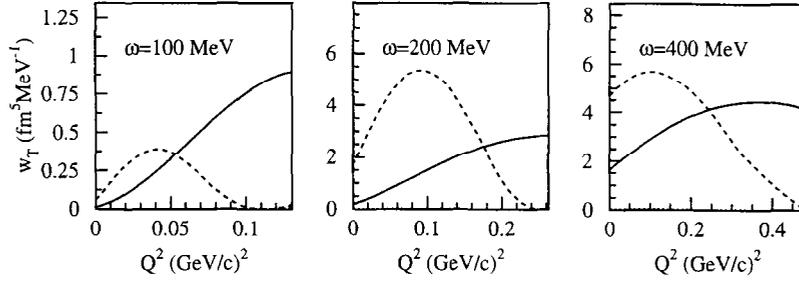


Fig. 2. The Q^2 dependence of the transverse term w_T in σ_{ep} for three values of the energy transfer. The kinematical condition is fixed through: $P = 50$ (MeV) $\mathbf{1}_x$, $\theta_1 = 135^\circ$ and $\phi_1 = 0^\circ$. The dotted (solid) line shows the contribution from photoabsorption on the one-body convection (magnetization) current. The contribution from intermediate Δ_{33} production is shown with the dashed line. In these calculations a Gaussian central correlation function was used. See text for further details.

$$\begin{aligned}
 g(\mathbf{p}) &\equiv \int d\mathbf{r} e^{i\mathbf{p}\cdot\mathbf{r}} (1 - f_C(r)), \\
 f_{\sigma T}(\mathbf{p}) &\equiv \int d\mathbf{r} e^{i\mathbf{p}\cdot\mathbf{r}} f_{\sigma T}(r), \\
 f_{tr}^0(\mathbf{p}) &\equiv \int d\mathbf{r} e^{i\mathbf{p}\cdot\mathbf{r}} Y_{20}(\Omega) f_{tr}(r).
 \end{aligned} \tag{9}$$

In the absence of ground-state correlations only the transverse Δ current would give a non-vanishing contribution to σ_{ep} . The longitudinal channel w_L is totally determined by the ground-state correlations.

Note that the σ_{ep} as written in Eq. (7) has formally the same form as the σ_{ep} in the $(e, e'p)$ case. The physics situation is, however, very different. In the $(e, e'p)$ case, the σ_{ep} contains information on electron scattering on a bound, off-shell nucleon. In quasi-elastic kinematics, this scattering process is predominantly sensitive to the one-body aspects of the target nucleus which are relatively well understood. For the $(e, e'pp)$ case the situation is completely different. Here, σ_{ep} contains information on the two-body aspects of the nuclear system. The $F_{h_1, h_2}(P)$ occurring in the factorized $(e, e'pp)$ cross section (2) is a less-challenging quantity. As it deals with the c.o.m. motion of dinucleons it is rather insensitive to nucleon-nucleon correlations at small and moderate values of P . Recent calculations [11] predict some sensitivity at large momenta P . However, by then the function $F_{h_1, h_2}(P)$ becomes so small in absolute magnitude that extremely small two-nucleon knockout cross sections can be expected. More favourable kinematical situations for probing correlations are created when considering relatively small values of P : in those cases larger cross sections are faced and the IPM predictions for $F_{h_1, h_2}(P)$ can be considered as realistic choices.

The (γ, pp) cross section is uniquely determined by the w_T term and reads

$$\frac{d^3\sigma}{d\Omega_1 d\Omega_2 dT_{p_2}} = E_1 p_1 E_2 p_2 \sigma_{\gamma pp}(k_+, k_-, q_\gamma) F_{h_1, h_2}(P), \tag{10}$$

with

$$\sigma_{\gamma pp}(k_+, k_-, q_\gamma) = \frac{1}{2E_\gamma (2\pi)^5} f_{\text{rec}}^{-1} w_T. \tag{11}$$

In the transverse response function w_T the contribution from the isobaric current and the terms related to ground-state correlations are competing. The $Q^2 = -q_\mu q^\mu$ dependence of the different terms in w_T is investigated in Fig. 2. We have considered in-plane kinematics (which means that both protons are escaping in the electron scattering plane) with a fixed \mathbf{P} (which is chosen to point along the x axis), θ_1 and ϕ_1 (polar and azimuthal angle of escaping proton 1). For the results of Fig. 2 the central correlation function of the Gaussian type

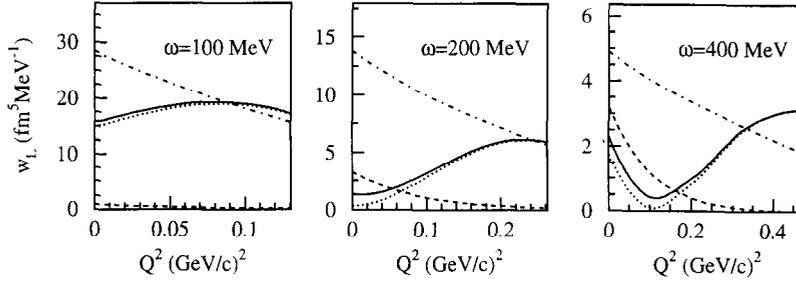


Fig. 3. As in Fig. 2 but now for the longitudinal term w_L . The dashed curves only include Jastrow correlations, the dotted curves include Jastrow and spin-spin correlations, and finally the solid curves include Jastrow, spin-spin and tensor correlations. All these results have been obtained with the variational correlation functions of Ref. [10]. The dot-dashed line includes only Jastrow correlations, but now calculated with the Gaussian correlation function of Ref. [13].

$g(r) = 0.51e^{1.52(\text{fm}^{-2})r^2}$ as suggested in Ref. [13] was used. For all three energy transfers considered, the $\Delta(1232)$ current produces the largest contribution at the real photon point ($Q^2 = 0$). Accordingly, (γ, pp) reactions are predicted to exhibit a rather small sensitivity to Jastrow correlations. The second term in w_T , which corresponds with photoabsorption on the one-body convection current, has a marginal effect on the cross section. The term related to the magnetization current (first term in the expression for w_T) has a q^2 dependence which makes it a dominant contribution at higher Q^2 .

In Fig. 3 the Q^2 dependence of the longitudinal term (w_L) in σ_{ep} is displayed for the same kinematical conditions as for Fig. 2. The first and second term in w_L are originating from the central and spin-spin correlations respectively. The interference between the spin-spin and central correlations gives rise to the third term of w_L . We remark that within the adopted approximations the contribution from the tensor correlations to w_L is restricted to an interference term with the central correlations. For the w_L we have investigated the role of the different terms in the correlation operator (6). Whereas central Jastrow correlations are frequently addressed in the literature, studies that include also spin-spin and tensor correlations are far less abundant. Using variational techniques the functions f_C , $f_{\sigma\tau}$ and f_{τ} have been determined starting from a realistic nucleon-nucleon interaction [9,10]. We have used the correlation functions from the ^{16}O calculations of Ref. [9]. It is striking that the Jastrow correlation function f_C that comes out of these variational methods is very soft. For the sake of reference, we show in Fig. 3 also the result when including only the Gaussian Jastrow correlation function that was introduced earlier. With this choice of the central correlation function the longitudinal $(e, e' \text{pp})$ strength is predicted to be considerably larger than the strength produced by the variational central correlation function f_C . The variational method, however, is a more complete model in the sense that it accounts for the full complexity of the ground-state correlations. It is clear from Fig. 3 that spin-spin correlations are predicted to contribute substantially to the longitudinal $(e, e' \text{pp})$ strength. The effect of tensor correlations is rather marginal. This is not too surprising as they are generally considered to be a predominant proton-neutron correlation effect [14]. Note that the predictions for w_L depend dramatically on the model assumptions with respect to the ground-state correlations.

In Fig. 4 we compare the predictions of the suggested factorized cross section with recent $^{12}\text{C}(e, e' \text{pp})$ data [15] and with unfactorized calculations. Two different regions in the excitation spectrum of the $A - 2$ system were considered: $E_x \leq 13 \text{ MeV}$ and $23 \text{ MeV} \leq E_x \leq 48 \text{ MeV}$. The first region can be attributed to $(1p)^2$ knockout and the second to both $(1p)(1s)$ and $(1s)^2$ knockout. For both of these regions we have calculated the angular cross sections with the variational and Gaussian correlation functions. With both choices one gets cross sections that are compatible with the data.

For the Gaussian results we have compared the predictions of the factorized approach to the results of an unfactorized calculation. These calculations are based on an extension of the model outlined in Refs. [16,17] and involve the same physics components as the unfactorized approach in the sense that both central correlations

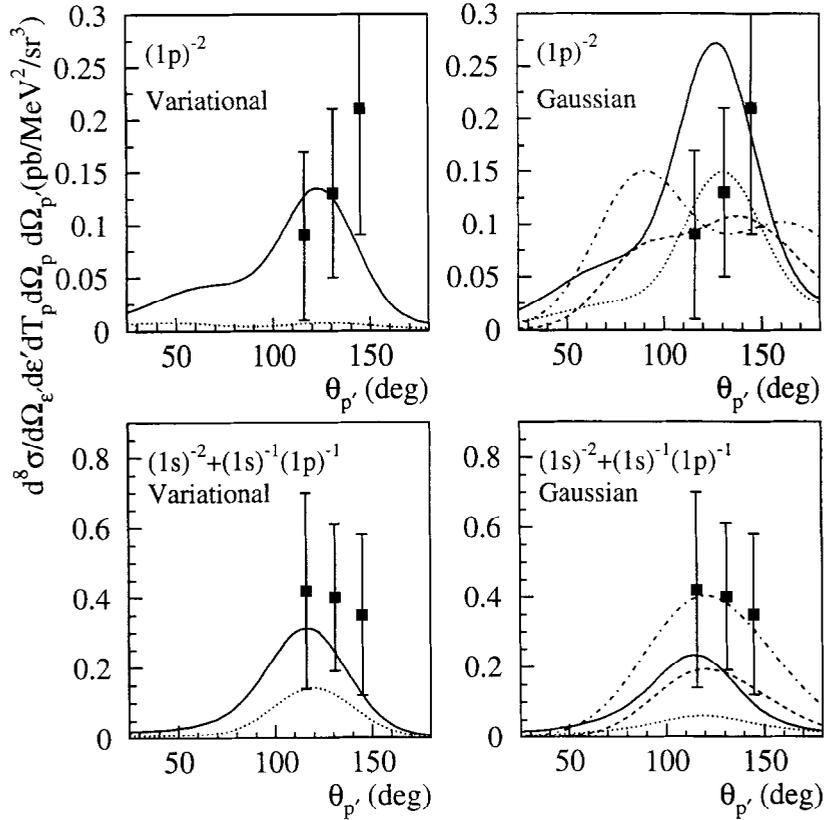


Fig. 4. The $^{12}\text{C}(e,e'pp)$ cross section for $\epsilon = 475$ MeV, $\omega = 212$ MeV and $q = 270$ MeV/c. One of the proton scattering angles was fixed at 27° . The dotted line is the calculated contribution from the ground-state correlations and the solid line is the prediction when including both ground-state correlations and intermediate Δ_{33} creation. Results for two types of correlation functions are shown. The data are from Ref. [15]. The dot-dashed (dashed) line shows the result of an unfactorized calculation with plane (distorted) outgoing proton waves.

and Δ_{33} effects are considered. Unlike in a factorized model, one is no longer bound to a plane wave description for the outgoing particles in an unfactorized approach [16,18,19]. The right panels of Fig. 4 show the results of the unfactorized calculations with both plane and distorted outgoing proton waves. Consequently the sole difference between the plane wave unfactorized calculations and the analytical approach discussed above is the treatment of the relative motion of the initial pair. It is noted that the unfactorized calculations produce angular cross sections that are wider than the analytical predictions. This is not too surprising given the fact that the relative motion of the initial diproton is now handled in its full complexity. After all, the factorized predictions could be considered as reasonable given that they can be performed in just a fraction of the computing time that the unfactorized calculations consume. Even though the analytical expression σ_{epp} should not be considered as a fully-fledged alternative for the cumbersome unfactorized calculations, it could help in optimizing the kinematical conditions and determining the major trends and sensitivities of the cross section.

Summarizing, we have derived a factorized form for the $(e,e'pp)$ and (γ,pp) cross section. This implies rather simple analytical expressions that should give a more transparent handle on the different physics components of photoinduced two-proton knockout reactions. Within the method developed here we have illustrated the sensitivity of the $(e,e'pp)$ cross sections to ground-state correlation effects.

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