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# Calculation of the cross section and the transverse-longitudinal asymmetry of the process $^3$He$(e,e'p)pn$ at medium energies within the unfactorized generalized Glauber approach

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The cross section and the transverse-longitudinal asymmetry $A_{\mathrm{TL}}$ of the three-body breakup process $^3$He$(e,e'p)pn$ have been calculated by a parameter-free approach based upon realistic few-body wave functions corresponding to the AV18 interaction, treating the rescattering of the struck nucleon within the unfactorized generalized eikonal approximation. The results of calculations exhibit a good agreement with recent JLab experimental data and show the dominant role played by the final state interaction, which, however, in the region of missing momentum, $300 \lesssim p_m \lesssim 600\,\mathrm{MeV}/c$, and removal energy corresponding to the two-body kinematic peak and higher, $E_m \gtrsim p_m^2/4m_N$, is dominated by single-nucleon rescattering, providing evidence that the final state interaction is mainly due to the one between the struck nucleon and a nearby correlated one.

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The exclusive reaction $A(e,e'p)X$, when an energetic electron knocks out quasielastically from the nucleus $A$ a proton $p$, which is detected in coincidence with the scattered electron $e'$, has provided in the past extremely rich information on shell-model single-particle (s.p.) properties, like the low momentum part of s.p. wave functions and the corresponding s.p. energies (see, e.g., Ref. [1]). Nowadays, the possibility to perform exclusive experiments at high-momentum transfer makes the $A(e,e'p)X$ process a powerful tool to investigate also those nuclear properties that are expected to exhibit deviations from independent particle motion, generated by short-range correlations (SRC) that should arise when two or more nucleons approach the range of the intermediate-distance tensor and short-distance repulsive nucleon-nucleon ($NN$) forces. Obtaining information on SRC in nuclei is a primary goal of modern nuclear physics [2], because not only would the detailed knowledge of SRC help to firmly establish the limits of validity of the standard model of nuclei, based on the assumption of pointlike nucleons interacting via two- and three-body interactions, but it would also have a strong impact on various fields of physics, like, to cite only a few examples, quark physics and QCD, the structure and properties of dense matter (e.g., neutron stars), relativistic heavy-ion collisions, and high-energy scattering processes [3]. Recently, a combined analysis of $A(e,e'p)X$ and related processes on the $^{12}$C nucleus at high-momentum transfer provided quantitative evidence on SRC (see Ref. [4] and references therein quoted) and has, at the same time, triggered intensive theoretical and experimental activities. Extracting information on SRC from $(e,e'p)X$ processes is, however, no easy task, because reconstructing the initial behavior in the nucleus of the detected proton is strictly possible only within the so-called plane wave impulse approximation (PWIA), that is, by assuming that when the proton is instantaneously removed from the nucleus it leaves it without any further interaction with the medium. In reality, final state interaction (FSI) effects of the proton with the medium (together with other effects, such as meson exchange currents (MEC), isobar production, etc.) may be very important (see, e.g., Ref. [5]). Nonetheless, by choosing proper kinematics these effects can be minimized and relevant information on bound nucleon dynamics can be obtained. Cross-section calculations based upon realistic correlated wave functions for initial and final states, coupled with sound models for the treatment of FSI, would help to choose the correct kinematics and will make the interpretation of the experimental data free from ambiguities. In this respect, few-nucleon systems represent an ideal test ground for any theoretical approach of the exclusive $A(e,e'p)X$ process, because various realistic microscopic approaches to generate initial and final nuclear wave functions are at our disposal. In our previous work [6] we calculated the cross section of the two-body breakup channel (2BBU), $^3$He$(e,e'p)^2$H, and the three-body breakup channel (3BBU), $^3$He$(e,e'p)np$, recently measured at JLab [7] in the range $0 \lesssim p_m \lesssim 1000\,\mathrm{MeV}/c$; here $p_m \equiv |\mathbf{p}_m|$ is the three-momentum of the undetected system $X$ ($^2$H or $pn$) that, within the PWIA, equals the momentum that the proton had before interaction with the virtual photon. In Ref. [6] a parameter-free theoretical approach has been used in which: (i) initial state correlations (ISC) in the target nucleus $^3$He have been taken care of by using state-of-the-art few-body wave functions obtained [8] by a variational solution of the Schrödinger equation containing realistic nucleon-nucleon interactions [9], and (ii) FSI have been treated by a generalized eikonal approximation (GEA) [10], which represents an extended Glauber approach (GA) based upon the evaluation of the relevant Feynman diagrams, which describe the rescattering of the struck nucleon in the final state, in analogy with the Feynman

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diagrammatic approach developed for the treatment of elastic hadron-nucleus scattering [11,12]. Calculations of Ref. [6] have been, however, performed within the factorization approximation (FA), that is, by writing the electron-nucleus cross section as a product of the off-shell electron-proton cross section and a distorted nuclear structure function, which is strictly valid only within the PWIA. The approach of Ref. [6] has been improved in Ref. [13] where the results of the calculation based upon an unfactorized generalized eikonal approximation (UGEA) of the 2BBU channel $^3$He$(e,e'p)^2$H have been presented, showing that: (i) the FSI of the struck proton in the 2BBU channel can satisfactorily be described within single and double rescattering of the struck proton (cf. Fig. 1 of Ref. [13]); (ii) the low-momentum part ($p_m \lesssim 300\,\mathrm{MeV}/c$) of the cross section is practically not affected by FSI and can even be described within the FA; on the contrary at $300 \lesssim p_m \lesssim 1000\,\mathrm{MeV}/c$, FSI becomes very important, particularly at "negative" values of the missing momentum (when $\phi$, the angle between the scattering and reaction planes, equals zero), and the FA breaks down (cf. Fig. 2 of Ref. [13]); (iii) in the region $300 \lesssim p_m \lesssim 700\,\mathrm{MeV}/c$ the effects of FSI are almost entirely exhausted by single rescattering (cf. Fig. 2 of Ref. [13]); and (iv) the transverse-longitudinal asymmetry $A_{\mathrm{TL}}$ is extremely sensitive to FSI and cannot be explained within the FA (cf. Fig. 3 of Ref. [13]). In the present article the results of UGEA for the 3BBU channel $^3$He$(e,e'p)pn$ are presented and compared with the experimental data of Ref. [14]. By considering the general case of a target nucleus $A$, the relevant quantities that characterize the process $A(e,e'p)(A-1)$ are the energy and momentum transfer $\nu$ and $Q^2$, respectively, the missing momentum $\mathbf{p}_m = \boldsymbol{q} - \boldsymbol{p}_1$ (i.e., the momentum of the recoiling system $A-1$), and the *missing energy* $E_m = \sqrt{P_{A-1}^2} + m_N - M_A = \nu - T_{A-1} - T_{\boldsymbol{p}_1} = |E_A| - |E_{A-1}| + E_{A-1}^* = E_{\min} + E_{A-1}^*$. Here, $\boldsymbol{p}_1$ is the momentum of the detected proton, $E_A(E_{A-1})$ the (negative) ground state energy of the target (recoiling) nucleus, and $E_{A-1}^*$ the intrinsic excitation energy of the latter ($E_{A-1}^* = E_m - |E_{\min}|$). The cross section of the process has the form

$$\frac{d^6\sigma}{d\Omega' dE' d\Omega_{p_1} dE_{p_1}} = \left| \frac{\mathbf{p}_1^2}{\frac{|\mathbf{p}_1|}{E_{p_1}} + \frac{|\mathbf{p}_1| - |\mathbf{q}|\cos\theta}{E_{A-1}^*}} \right| \frac{dE_m}{dE_{p_1}} \sigma_{\mathrm{Mott}} \times \sum_i V_i W_i^A(\nu, Q^2, \mathbf{p}_m, E_m), \quad (1)$$

where $i \equiv \{\mathrm{L, T, TL, TT}\}$, $V_i$ are kinematical factors, and the nuclear structure functions $W_i^A$ result from proper combinations of the polarization vector of the virtual photon, $\varepsilon_\lambda^\mu$, and the hadronic tensor, $W_{\mu\nu}^A$, the latter depending upon the nuclear current operators $\hat{J}_\mu^A(0)$. We consider the interaction of the incoming virtual photon $\gamma^*$ with a bound nucleon (the active nucleon) of low virtuality ($p^2 \sim m_N^2$) in the quasielastic kinematics, that is, corresponding to $x \equiv Q^2/2m_N\nu \sim 1$. In quasielastic kinematics, the virtuality of the struck nucleon after $\gamma^*$ absorption is also rather low and, provided $\mathbf{p}_1$ is sufficiently high, nucleon rescattering with the "spectator" $A-1$ can be described to a large extent in terms of multiple elastic scattering processes in the eikonal approximation. Let us now consider the process $^3$He$(e,e'p)pn$. In coordinate representation the initial three-body wave function is $\Phi_{1/2,\mathcal{M}_3}(\boldsymbol{r}_1, \boldsymbol{r}_2, \boldsymbol{r}_3)$, whereas the wave function of the final state is

$$\Psi_f^* = \hat{\mathcal{A}}\big[S_\Delta^{\mathrm{FSI}}(\boldsymbol{r}_1, \boldsymbol{r}_2, \boldsymbol{r}_3) e^{-i\mathbf{p}_1\mathbf{r}_1} e^{-i\mathbf{p}_m(\boldsymbol{r}_2+\boldsymbol{r}_3)/2} \times \Phi_{S_f\sigma_f}^{\mathbf{k}_{\mathrm{rel}}*}(\boldsymbol{r}_2 - \boldsymbol{r}_3)\chi^+_{\frac{1}{2},\lambda_f}\big], \quad (2)$$

where $\hat{\mathcal{A}}$ is a proper antisymmetrization operator, $\Phi_{S_f\sigma_f}^{\mathbf{k}_{\mathrm{rel}}*}(\boldsymbol{r}_2 - \boldsymbol{r}_3)$ describes the relative motion of two interacting particles in the continuum, and, eventually, $\chi_{\frac{1}{2},\lambda_f}$ is the spin wave function of the struck nucleon and

$$S_\Delta^{\mathrm{FSI}}(\boldsymbol{r}_1, \boldsymbol{r}_2, \boldsymbol{r}_3) = 1 + T_{(1)}^{\mathrm{FSI}}(\boldsymbol{r}_1, \boldsymbol{r}_2, \boldsymbol{r}_3) + T_{(2)}^{\mathrm{FSI}}(\boldsymbol{r}_1, \boldsymbol{r}_2, \boldsymbol{r}_3) \quad (3)$$

is the GEA $S$ matrix, which describes both the case of no FSI and single and double rescattering of the struck nucleon off the spectator nucleons [6,10]. By introducing the corresponding Jacobi coordinates $\mathbf{R}$, $\boldsymbol{\rho}$, and $\boldsymbol{r}$, the relevant matrix elements can be computed solely in terms of relative coordinates $\boldsymbol{\rho}$ and $\boldsymbol{r}$. For ease of presentation we give explicit expressions of the single rescattering contribution only, which has the form $T_{(1)}^{\mathrm{FSI}}(\boldsymbol{\rho}, \boldsymbol{r}) = -\sum_{i=2}^3 \theta(z_i - z_1) e^{i\Delta_z(z_i - z_1)}\Gamma(\mathbf{b}_1 - \mathbf{b}_i)$, where $\boldsymbol{r}_i \equiv (z_i, \mathbf{b}_i)$. It can be seen that, unlike the usual GA, the GEA $S$ matrix gets also a contribution from a parallel momentum $\Delta_z$ of pure nuclear origin depending upon the external kinematics and the removal energy of the struck proton (within the "frozen approximation" $\Delta_z = 0$, and the usual Glauber profile is recovered). By assuming that the nuclear current operator is a sum of nucleonic currents $j_\mu(i)$, its matrix elements resulting from Feynman diagrams can be written in momentum space as follows:

$$J_\mu^3 = \sum_\lambda \int \frac{d\mathbf{p}}{(2\pi)^3} \frac{d\boldsymbol{\kappa}}{(2\pi)^3} S_\Delta^{\mathrm{FSI}}(\boldsymbol{p}, \boldsymbol{\kappa}) \langle \lambda_f | j_\mu(\boldsymbol{\kappa} - \boldsymbol{p}_m; \boldsymbol{q}) | \lambda \rangle \times \mathcal{O}(\boldsymbol{p}_m - \boldsymbol{\kappa}, \boldsymbol{p}, \mathbf{k}_{\mathrm{rel}}; \mathcal{M}_3, S_f, \sigma_f, \lambda) = J_\mu^{3(\mathrm{PWIA})} + J_\mu^{3(1)} + J_\mu^{3(2)}, \quad (4)$$

where $S_\Delta^{\mathrm{FSI}}(\boldsymbol{p}, \boldsymbol{\kappa}) = \int d\boldsymbol{r} d\boldsymbol{\rho} e^{-i\mathbf{p}\mathbf{r}} e^{i\boldsymbol{\kappa}\boldsymbol{\rho}} S_\Delta^{\mathrm{FSI}}(\boldsymbol{\rho}, \boldsymbol{r})$ is the Fourier transform of the GEA $S$ matrix [Eq. (3)], $\mathcal{O}$ is the nuclear overlap in momentum space,

$$\mathcal{O}(\boldsymbol{p}_m - \boldsymbol{\kappa}, \boldsymbol{p}, \mathbf{k}_{\mathrm{rel}}; \mathcal{M}_3, S_f, \sigma_f, \lambda) = \int d\boldsymbol{\rho} d\mathbf{r} e^{i(\mathbf{p}_m - \boldsymbol{\kappa})\boldsymbol{\rho}} e^{i\mathbf{p}\mathbf{r}} \Phi_{\frac{1}{2},\mathcal{M}_3}(\boldsymbol{\rho}, \boldsymbol{r}) \Phi_{S_f,\sigma_f}^{\mathbf{k}_{\mathrm{rel}}*}(\boldsymbol{r}) \chi^\dagger_{\frac{1}{2}\lambda}, \quad (5)$$

and, eventually, $\langle \lambda_f | j_\mu(\boldsymbol{\kappa} - \boldsymbol{p}_m; \boldsymbol{q}) | \lambda \rangle$ is the nucleon current operator, with $\lambda$ and $\lambda_f$ being the spin projections of the struck proton before and after $\gamma^*$ absorption. Note that in Eq. (5) the quantity $\boldsymbol{\kappa} - \boldsymbol{p}_m \equiv \mathbf{k}_1$ represents exactly the momentum of the struck proton before the electromagnetic interaction, with $\boldsymbol{\kappa}$ being the momentum transfer in single rescattering. It is worth emphasizing that, because our calculations are performed in momentum space, nonrelativistic reduction of the current operator $j_\mu$ is not needed. In the absence of any FSI the momentum space $S$ matrix is $S_\Delta^{\mathrm{FSI}}(\boldsymbol{\kappa}, \mathbf{p}) = (2\pi)^6\delta(\boldsymbol{\kappa})\delta(\mathbf{p})$ and only the PWIA contribution $J_\mu^{3(\mathrm{PWIA})}$ survives in Eq. (4). When FSI is active, the contributions from single and double rescattering must be taken into account. Let us consider the former: it results from the single-scattering momentum space term of $S_{\Delta}^{\rm FSI}(\boldsymbol{p},\boldsymbol{\kappa})$, which has the following form,

$$T_{(1)}^{\rm FSI}(\boldsymbol{p},\boldsymbol{\kappa}) = \frac{(2\pi)^4}{k^*}\frac{f_{NN}(\boldsymbol{\kappa}_\perp)}{\boldsymbol{\kappa}_{||}+\Delta_z - i\varepsilon} \times \left[\delta\left(\boldsymbol{p}-\frac{\boldsymbol{\kappa}}{2}\right)+\delta\left(\boldsymbol{p}+\frac{\boldsymbol{\kappa}}{2}\right)\right], \qquad (6)$$

where $f_{NN}(\boldsymbol{\kappa}_\perp)$ [the Fourier transform of the profile function $\Gamma(\mathbf{b})$] represents the elastic scattering amplitude of two nucleons with center-of-mass momentum $k^*$. Placing Eq. (6) in Eq. (4) the single-scattering contribution $J_\mu^{3(1)}$ to the nuclear current is obtained as follows,

$$J_\mu^{3(1)} = \sum_\lambda \int \frac{d\boldsymbol{\kappa}}{(2\pi)^2 k^*}\langle\lambda_f|j_\mu(\boldsymbol{k}_1;\boldsymbol{q})|\lambda\rangle \frac{f_{NN}(\boldsymbol{\kappa}_\perp)}{\boldsymbol{\kappa}_{||}+\Delta_{||}-i\varepsilon} \times \left[\mathcal{O}\left(-\boldsymbol{k}_1,\frac{\boldsymbol{\kappa}}{2};\mathcal{M}_3,S_f,\sigma_f,\lambda\right) + \mathcal{O}\left(-\boldsymbol{k}_1,-\frac{\boldsymbol{\kappa}}{2};\mathcal{M}_3,S_f,\sigma_f,\lambda\right)\right], \qquad (7)$$

where $\boldsymbol{k}_1$, as already mentioned, is the momentum of the proton before $\gamma^*$ absorption and $\boldsymbol{\kappa} = \boldsymbol{k}_1 + \boldsymbol{q} - \boldsymbol{p}_1 = \boldsymbol{k}_1 + \boldsymbol{p}_m$ is the momentum transfer in the $NN$ rescattering. The longitudinal part of the nucleon propagators has been computed using the relation $[\boldsymbol{\kappa}_{||}+\Delta_z \pm i\varepsilon]^{-1} = \mp i\pi\delta(\boldsymbol{\kappa}_{||}+\Delta_z) + PV[\boldsymbol{\kappa}_{||}+\Delta_z]^{-1}$. Note that in the eikonal approximation the trajectory of the fast nucleon is a straight line so that all "longitudinal" and "perpendicular" components are defined with respect to this trajectory, which means that the $z$ axis must be directed along the momentum of the detected fast proton. Looking at the structure of Eq. (7), it can be seen that, due to the coupling of the nucleonic current operator $\langle\lambda_f|j_\mu(\boldsymbol{k}_1;\boldsymbol{q})|\lambda\rangle$ with the nuclear overlap integral, a factorized form for Eq. (7) cannot be obtained; the same holds for the double-scattering contribution and, consequently, for the cross section. However, as shown in Ref. [6], if the longitudinal part of the nucleonic current can be disregarded, the factorization form can approximately be recovered. Using the aforementioned formalism, and including the contribution from double rescattering, $J_\mu^{3(2)}$, the cross section [Eq. (1)] and the left-right asymmetry defined by

$$A_{\rm TL} = \frac{d\sigma(\phi=0^\circ) - d\sigma(\phi=180^\circ)}{d\sigma(\phi=0^\circ)+d\sigma(\phi=180^\circ)} \qquad (8)$$

have been calculated. Following de Forest's prescription [15], the "CC1" form of the nucleon current operator has been adopted; the elastic amplitude $f_{NN}$ has been chosen in the usual form $f_{NN}(\boldsymbol{\kappa}_\perp) = k^*\frac{\sigma^{\rm tot}(i+\alpha)}{4\pi}e^{-b^2\boldsymbol{\kappa}_\perp^2/2}$, where the slope parameter $b$, the total nucleon-nucleon cross section $\sigma^{\rm tot}$, and the ratio $\alpha$ of the real to the imaginary parts of the forward scattering amplitude were taken from the world's experimental data. The results of our calculations are shown in Figs. 1, 2, and 3. In Fig. 1 the full unfactorized cross section at two values of $p_m$ is compared with the PWIA result and with the experimental data; in this figure the role played by single and double rescattering is clearly illustrated. In Fig. 2 the factorized and unfactorized results are compared at $p_m = 440\,\mathrm{MeV}/c$

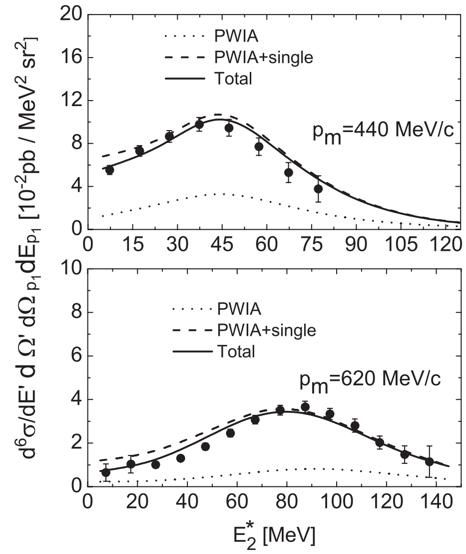


FIG. 1. The differential cross section [Eq. (1)] for the process $^3$He$(e,e'p)pn$ calculated at two values of the missing momentum versus the excitation energy of the two-body final state $E_2^* = E_m - E_{\rm min}$ ($E_{\rm min} = 2m_p + m_n - M_3$). Dotted lines, PWIA approximation; dashed and solid lines, unfactorized calculations with single and double rescattering in the final state, respectively. Experimental data are from Ref. [14].

(similar results are obtained at higher values of $p_m$). Finally in Fig. 3 the asymmetry $A_{\rm TL}$ is exhibited. As in Ref. [13], no approximations have been made in the evaluation of the single- and double-rescattering contributions: intrinsic coordinates have been used and the energy dependence of the profile function has been taken into account in the properly chosen CM system of the interacting pair. Note that the experimental kinematics has been intentionally chosen so as to cover the theoretically predicted two-nucleon correlation region characterized by a peak position given by $E_2^* = E_m - E_{\rm min} \sim 2m_N\left\{\sqrt{(1/2)[1+\sqrt{1+(p_m/m_N)^2}]}-1\right\} \to p_m^2/4m_N$ in the nonrelativistic limit. It can be seen that the location of the experimental peaks agrees with such a prediction. The amplitude of the peak is largely affected by FSI, whose inclusion brings the theoretical calculation into very good agreement with the experimental data. The relevant effects of FSI can be explained by the fact that the kinematics of the experiment is such

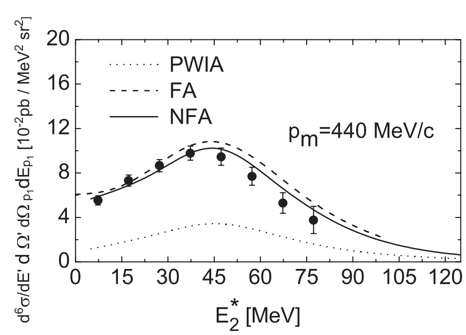


FIG. 2. Comparison at $p_m = 440\,\mathrm{MeV}/c$ of the results of calculations performed within the unfactorized approach (NFA) and the factorization approximation (FA).

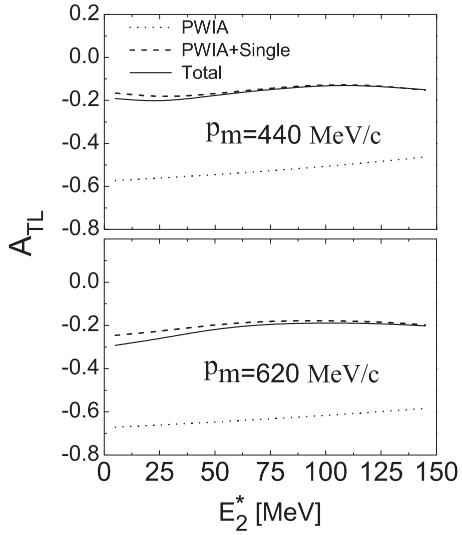


FIG. 3. The $A_{\rm TL}$ asymmetry [Eq. (8)] calculated in correspondence of two values of $p_m$ versus the excitation energy of the two-body final state. The solid line includes both single and double scattering.

that the missing momentum is always almost perpendicular to the momentum of the final proton and FSI is maximized. Concerning the effects of the latter, an interesting observation is in order: it can be seen that in the region of the correlation peak, and at higher values of $E_2^*$, the effects from double rescattering are very tiny; the same appears to be the case in the 2BBU channel at $300 \lesssim p_m \lesssim 700\,{\rm MeV}/c$. We consider this as more evidence that in these regions of missing momentum and energy FSI mainly occurs within the correlated pair. In Fig. 2 the factorized and unfactorized results are compared at $p_m = 440\,{\rm MeV}/c$; it can be seen that the unfactorized calculation differs by 5–7% from the factorized one and better agrees with the experimental data. Eventually, in Fig. 3 the transverse-longitudinal asymmetry $A_{\rm TL}$ [Eq. (8)], calculated within the PWIA and taking into account FSI, is presented. It is well known that, when the explicit expressions of $V_i$ and $W_i^A$ are placed in Eq. (8), the numerator is proportional to the transverse-longitudinal response $W_{\rm TL}$, whereas the denominator does not contain $W_{\rm TL}$ at all, which means that $A_{\rm TL}$ is a measure of the relevance of the transverse-longitudinal response relative to the other responses. In the $eN$ cross section the behavior of the asymmetry $A_{\rm TL}$ is known to be a negative and decreasing function of the missing momentum [15] and the same behavior should be expected in $eA$ scattering within the PWIA. The asymmetry presented in Fig. 3 clearly shows that a measurement of the 3BBU channel at $\phi = 0^\circ$ would provide results not very different from the ones obtained at $\phi = 180^\circ$. To sum up, we have calculated in momentum space the cross section of the processes $^3{\rm He}(e,e'p)pn$, using realistic ground-state two- and three-body wave functions, treating the FSI of the struck nucleon with the spectator nucleon pair within the generalized eikonal approximation. As in the case of the 2BBU calculation [13], our approach is a parameter-free one, for it only requires the knowledge of the nuclear wave functions, the FSI factor being fixed directly by $NN$ scattering data. At the same time, calculations are very involved mainly because of the complex structure of the wave function of Ref. [8], which has to be first transformed to momentum space and then used in the calculations of multidimensional integrals, including also the computation of principal values. The main results of our calculations can be summarized as follows: (i) a good agreement between theoretical calculations and experimental data on the 3BBU process $^3{\rm He}(e,e'p)pn$ has been achieved; (ii) FSI effects are very relevant, but in the region of the "correlation peak" ($E_2^* \simeq \frac{p_m^2}{4m}$) and at higher values of removal energies ($E_2^* \geqslant \frac{p_m^2}{4m}$), they can be accounted for by single rescattering, the double-rescattering contribution being negligible, which means that, within the Glauber approach, the struck, or "active," particle is within the short range of nuclear force, whereas the third particle is outside it; this represents an implicit demonstration that at high-momentum transfer, the final state interaction is mainly concentrated in the correlated pair; (iii) factorized and unfactorized results may differ at $\phi = 180^\circ$ by at most 10%; and (iv) the $A_{\rm TL}$ asymmetry is in qualitative agreement with unpublished experimental results [16]. In closing the present article, we would like once again to point out (see Ref. [13]) that unfactorized calculations for complex nuclei within particular models of the ground-state wave functions are nowadays a common practice (see, e.g., Ref. [5] and references therein quoted), whereas in the case of few-body systems described by realistically correlated wave functions they are very involved and not so common. As a matter of fact, in the case of the $^3{\rm He}(e,e'p)X$ process, besides our calculations, to our knowledge only two other ones have been published: the one of Ref. [17], similar in spirit to our approach, but where only the 2BBU channel has been considered, and the one of Ref. [18] where both the 2BBU and 3BBU have been calculated within the diagrammatic approach of Ref. [19], obtaining a good agreement with the experimental data by an alternative three-body mechanism that differs from the two-body Glauber rescattering of our approach.

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