The longitudinal asymmetry of the $(e, e'p)$ missing momentum distribution as a signal of color transparency

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Abstract

We use multiple scattering theory to evaluate the $Q^2$ dependence of the forward-backward asymmetry, $A_z$, of the missing momentum distribution in quasielastic $A(e, e'p)$ scattering, which is expected to be affected by color transparency. The novel features of our analysis are a consistent treatment of the structure of the ejectile state formed after absorption of the virtual photon by the struck proton and a careful evaluation of the background asymmetry induced by the nonvanishing real part of the proton-nucleon scattering amplitude. We find that the absolute magnitude of $A_z$ is dominated by this background at the $Q^2$'s attainable at the Continuous Electron Beam Accelerator Facility. However, the $Q^2$ dependence of the asymmetry is sensitive to the onset of color transparency, whose observation in high statistics experiments appears to be feasible, particularly using light targets such as $^{12}C$ or $^{16}O$. 
The recent $A(e,e'p)$ experiment carried out at SLAC by the NE18 collaboration did not observe any substantial $Q^2$ dependence of the integrated nuclear transparency up to $Q^2 \lesssim 7 \text{ GeV}^2$. On the other hand, theory predicts that final state interactions (FSI) should vanish at asymptotically high $Q^2$ on account of color transparency (CT). From the point of view of multiple scattering theory (MST), CT corresponds to a cancellation between the rescattering amplitudes with elastic (diagonal) and inelastic (off-diagonal) intermediate states. The NE18 finding was anticipated in ref.4, where it was argued that in the integrated nuclear transparency the effect of the inelastic rescatterings of the struck proton is still weak in the region $Q^2 \lesssim 7 \text{ GeV}^2$, covered by NE18 and achievable by the forthcoming experiments at CEBAF. The cancellation of the diagonal and off-diagonal rescatterings depends upon the missing momentum [5, 6, 7], which can be used to enhance the contribution from inelastic rescatterings and make the onset of CT observable at CEBAF. The enhancement and/or suppression of the contribution of inelastic rescatterings result in a non-vanishing forward-backward asymmetry of the missing momentum distribution $w(\vec{p})$, defined as

$$A_z(x, y) = \frac{N_+ - N_-}{N_+ + N_-},$$

where $N_+$ ($N_-$) is the number of events in the kinematical region $x < p_z < y (-y < p_z < -x)$ and the z-axis is parallel to the ($e, e'$) momentum transfer $\vec{q}$.

Unfortunately, besides CT, there is another competing source of asymmetry associated with FSI: the nonvanishing ratio between the real and the imaginary part of the elastic proton-nucleon ($pN$) scattering amplitude, $\alpha_{pN}$. The forward-backward asymmetry, associated with $\alpha_{pN}$, disregarded in refs.[3, 4, 5], was recently found to be sizeable [8]. The role played by the real part of the rescattering amplitude in electron nucleus scattering at high $Q^2$ has been also pointed out in ref.[9], within the context of an analysis of the inclusive data from SLAC.

In this paper we give a quantitative evaluation of how much the CT signal can be obscured by the effect of a nonzero $\alpha_{pN}$. From the technical point of view, one must calculate $A_z$ taking into account the coupled-channel treatment of the off-diagonal rescatter-
terings and the nonzero $\alpha_{pN}$ on the same footing. Besides including the effect of $\alpha_{pN}$, we improve upon the previous works treating the coherency properties of off-diagonal rescatterings in a more accurate manner.

The missing momentum distribution equals

$$w(p) = \frac{1}{(2\pi)^3} \int d^3\vec{r}_1 d^3\vec{r}_1' d^3\vec{r}_2 \cdots d^3\vec{r}_A \rho_p(\vec{r}_1, \vec{r}_2, \ldots, \vec{r}_A; \vec{r}_1', \vec{r}_2', \ldots, \vec{r}_A') \times S^*(\vec{r}_1', \vec{r}_2', \ldots, \vec{r}_A) S(\vec{r}_1, \vec{r}_2, \ldots, \vec{r}_A) \exp[i\vec{p}(\vec{r}_1' - \vec{r}_1)], \quad (2)$$

where $\rho_p(\vec{r}_1, \vec{r}_2, \ldots, \vec{r}_A; \vec{r}_1', \vec{r}_2', \ldots, \vec{r}_A')$ is the proton $A$-body semidiagonal density matrix \cite{10} and

$$S(\vec{r}_1, \vec{r}_2, \ldots, \vec{r}_A) = \frac{\langle p| \hat{S}_{3q}(\vec{r}_1, \vec{r}_2, \ldots, \vec{r}_A) |E \rangle}{\langle p|E \rangle}. \quad (3)$$

In eq. (3), $\hat{S}_{3q}(\vec{r}_1, \vec{r}_2, \ldots, \vec{r}_A)$ is the operator describing the evolution of the ejectile state during its propagation through the nuclear medium and the ejectile state $|E\rangle$ is a three-quark state formed from the struck proton after absorption of the virtual photon at position $\vec{r}_1$. In terms of the electromagnetic current operator $\hat{J}_{em}$, $|E\rangle = \hat{J}_{em}(Q)|p\rangle = \sum |i\rangle \langle i|J_{em}(Q)|p\rangle = \sum G_{ip}(Q)|i\rangle$, where the sum includes the complete set of three-quark states and $G_{ip}(Q) = \langle i|J_{em}(Q)|p\rangle = \langle i|E\rangle$ denotes both the electromagnetic form factor of the proton ($|i\rangle = |p\rangle$) and all the transition form factors associated with the electroexcitation processes $e^+ + p \rightarrow e'^+ + i$ \cite{11}.

We will consider the forward-backward asymmetry at moderate missing momentum $p \lesssim 200$ MeV, where the main effect of short range $NN$ correlations in the initial state is an overall renormalization of the momentum distribution \cite{11}, which does not affect the calculation of $A_z$. Although $NN$ correlations also contribute to the integrated transparency \cite{12, 13}, their effect is expected to be less important in $A_z$. Neglecting exchange and dynamical correlations between the struck nucleon and the spectator nucleons, the proton $\nu$-body semidiagonal density matrix, $\rho(\vec{r}_1, \vec{r}_2, \ldots, \vec{r}_\nu; \vec{r}_1', \vec{r}_2', \ldots, \vec{r}_\nu)$ takes the factorized form $\rho_p(\vec{r}_1, \vec{r}_1')\rho(\vec{r}_2, \ldots, \vec{r}_\nu)$, where $\rho(\vec{r}_2, \ldots, \vec{r}_\nu)$ is the $(\nu - 1)$-body diagonal density matrix which can be recast in the form \cite{14} $\Pi_{i=2}^{\nu-1} \rho(\vec{r}_i) g(\vec{r}_2, \ldots, \vec{r}_\nu)$ where $\rho(\vec{r}_i) = n_A(\vec{r}_i)/A$ is the one-body density and the $(\nu - 1)$-body distribution function \cite{10}, $g(\vec{r}_2, \ldots, \vec{r}_\nu)$, brings
in the effect of NN correlations. We also neglect NN correlations among the spectator nucleons in this first study of longitudinal asymmetry. This implies that $g(\vec{r}_2, \cdots, \vec{r}_\nu) = 1$ and one ends up with the results of ref. [8] that at moderate missing momenta $p \lesssim 200$ MeV/c, the integrand of (2) - the FSI-modified one-body density matrix - can be approximated as

$$
\rho_p(\vec{r}_1, \vec{r}_1') \int d^3 \vec{r}_2 \cdots d^3 \vec{r}_A \prod_{i=1}^A \rho(\vec{r}_i) S^*(\vec{r}_1', \vec{r}_2, \cdots, \vec{r}_A) S(\vec{r}_1, \vec{r}_2, \cdots, \vec{r}_A) = \rho_p(\vec{r}_1, \vec{r}_1') \Phi(\vec{r}_1, \vec{r}_1')
$$

(4)

where, $\rho_p(\vec{r}_1, \vec{r}_1') = \frac{1}{Z} \sum_n \phi_n^*(\vec{r}_1') \phi_n(\vec{r}_1)$ is the proton shell model one-body density matrix and $\phi_n$ are the shell model wavefunctions.

In the MST, the FSI operator $\hat{S}^*_3 \hat{S}_3$ can be very schematically represented as $S^*_3 S_3 = \prod_{j=2}^A (1 - \hat{\Gamma}^*_1 \hat{\Gamma}_j)(1 - \hat{\Gamma}_1 \hat{\Gamma}_j)$, where the profile function $\hat{\Gamma}$ is an operator acting on the state describing the internal structure of the three-quark system [8]. Then, the calculation of $\Phi(\vec{r}_1, \vec{r}_1')$ involves a sum of diagrams, as shown in Fig. 1. Every dotted line attached to the straight-line trajectory originating from the point $\vec{r}_1$ ($\vec{r}_1'$) denotes a profile function $\hat{\Gamma}_j$ ($\hat{\Gamma}^*_1 \hat{\Gamma}_j$), where the subscripts $1j$($1'j$) stands for the position of the $j$-th spectator and its projection on the trajectory. The interaction between the two trajectories, due to terms $\hat{\Gamma}_1 \hat{\Gamma}_j$ like the one shown in Fig. 1b, makes the FSI factor $\Phi(\vec{r}_1, \vec{r}_1')$ a non factorizable function of $\vec{r}_1$ and $\vec{r}_1'$. It can be shown that the non factorizable part of $\Phi(\vec{r}_1, \vec{r}_1')$ is only connected with the incoherent elastic rescatterings, which come into play only at $p \gtrsim 200$ MeV [8]. For this reason, at $p \lesssim 200$ MeV, we can neglect $\hat{\Gamma}_1 \hat{\Gamma}^*_1$ terms in the FSI operator $\hat{S}^*_3 \hat{S}_3$ and obtain the factorizable $\Phi(\vec{r}_1, \vec{r}_1') = S^*_3(\vec{r}_1) S_3(\vec{r}_1')$ and the following expression for $w(\vec{p})$,

$$
w(\vec{p}) = \frac{1}{Z} \sum_n \left| \int d^3 \vec{r} \phi_n(\vec{r}) \exp[-i \vec{p} \cdot \vec{r}] S_{coh}(\vec{r}) \right|^2,
$$

(5)

where

$$
S_{coh}(\vec{r}_1) = \int \prod_{j=2}^A \rho(\vec{r}_j) d^3 \vec{r}_j \frac{\langle p| \hat{S}_3(\vec{r}_1, \vec{r}_2, \cdots, \vec{r}_A)|E \rangle}{\langle p|E \rangle}.
$$

(6)

Finally, using the Glauber form of $\hat{S}_3$, the coupled channel MST series for $S_{coh}(\vec{r})$ in terms of the $\nu$-fold off-diagonal rescatterings can be cast in the form $S_{coh}(\vec{r}) = \sum_{\nu=0}^\infty S_{coh}^\nu(\vec{r})$, where
the first few terms read
\[
S_{coh}(\vec{b}, z) = 1 - \exp[ik_{i_1p}z] \sum_{i_1} \frac{\sigma'_{pi_1}}{2} \frac{\langle i_1 | E \rangle}{\langle p | E \rangle} \int_z^\infty dz_1 n_A(\vec{b}, z_1) \exp[ik_{pi_1}z_1 - \frac{1}{2}t(\vec{b}, z_1, z)\sigma_{i_1i_1}]
\]
\[
+ \exp[ik_{i_2p}z] \sum_{i_1i_2} \frac{\sigma'_{pi_2}}{2} \frac{\sigma'_{i_2i_1}}{2} \frac{\langle i_1 | E \rangle}{\langle p | E \rangle} \int_z^\infty dz_1 n_A(\vec{b}, z_1) \exp[ik_{i_2i_1}z_1 - \frac{1}{2}t(\vec{b}, z_1, z)\sigma_{i_1i_1}]
\]
\[
\times \int_{z_1}^\infty dz_2 n_A(\vec{b}, z_2) \exp[ik_{pi_2}(z_2 - z_1) - \frac{1}{2}t(\vec{b}, z_2, z_1)\sigma_{i_2i_2}]
\]
\[
+ \cdots,
\]
Eq. (7)

where \( \vec{r} = (\vec{b}, z) \), \( \sigma'_{ik} = \sigma_{ik} - \delta_{ik}\sigma_{ii} \), the matrix \( \hat{\sigma} = 2\int d^2\vec{b}\hat{\Gamma}(\vec{b}) \) is connected with the forward diffraction scattering matrix \( \hat{f} = i\hat{\sigma} \) (notice that \( \text{Re} \sigma_{ii} = \sigma_{tot}(iN) \)) and \( t(\vec{b}, z_2, z_1) = \int_{z_1}^{z_2} dz n_A(\vec{b}, z) \) is the partial optical thickness. Eq. (7) shows contributions from up to two off-diagonal rescatterings and any number of elastic rescatterings in between, higher order off-diagonal rescatterings are implied by the dots. The onset of CT is controlled by the oscillating factors \( \exp[ik_{i_ni_{n-1}}z_n] \), taking into account the longitudinal momentum transfer associated with each off-diagonal transition \( i_{n-1}N \rightarrow i_nN \) \[13\] according to
\[
k_{i_{n-1}i_n} = \frac{m_{i_n}^2 - m_{i_{n-1}}^2}{2\varepsilon},
\]
where \( m_i \) denotes the mass of the proton excitation \( |i\rangle \) and \( \varepsilon \) is the energy of the struck proton in the laboratory frame.

From the point of view of CT, the most important ingredient in the coupled-channel formalism is the Pomeron contribution to the matrix \( \hat{\sigma} \). As in ref. \[4\], we will construct \( \hat{\sigma} \) using the oscillator quark-diquark model of the proton. As a result, the Pomeron contribution to the matrix element \( \sigma_{ik} \) can be written as
\[
\text{Im} f(kN \rightarrow iN) = \text{Re} \sigma_{ik}^P = \int dz d^2\rho \Psi_i^*(\vec{r}, z)\sigma(\rho)\Psi_k(\vec{r}, z),
\]
where \( \Psi_{i,k} \) are the oscillator wave functions describing the quark-diquark states and \( \sigma(\rho) \) is the dipole cross section describing the interaction of the quark-diquark system with a nucleon, which was taken in the form \( \sigma(\rho) = \sigma_0 \left[ 1 - \exp\left(-\frac{\rho^2}{R_0^2}\right) \right] \). Following ref. \[4\], we set \( \sigma_0 = 2\sigma_{tot}(pN) \) and adjust \( R_0 \) to reproduce \( \sigma_{tot}^{pp}(pN) \). For the diagonal \( iN \rightarrow iN \) scattering we take \( \text{Re} f(iN \rightarrow iN) = -\text{Im} \sigma_{ii} = \frac{1}{2} (\alpha_{pp}\sigma_{tot}(pp) + \alpha_{pn}\sigma_{tot}(pn)) \). Even though this
choice can be justified within the framework of the dual parton model \[16\], we are fully aware that it should only be regarded as an estimate. The real parts \(\text{Re} f(iN \rightarrow kN)\) of the off-diagonal amplitudes are not known experimentally. We shall give an estimate of the sensitivity of our results to these quantities and the associated uncertainty.

For the oscillator frequency of the quark-diquark system we use the value \(\omega_{qd} = 0.35\) GeV, leading to a realistic mass spectrum of the proton excitations. Since we start with a \(\sigma(\rho)\) which vanishes as \(\rho \rightarrow 0\), our diffraction matrix satisfies the CT sum rules \[6, 17\] by construction, and leads to vanishing FSI in \((e,e'p)\) at asymptotically large \(Q^2\). Moreover, with the above form of \(\sigma(\rho)\), we obtain a diffraction matrix yielding a realistic description of the mass spectrum observed in diffractive \(pN\) scattering \[17\]. For instance, the present model gives a value of the ratio between the diffractive and the elastic \(pN\) cross sections in agreement with the experimental data: \(\sigma_{\text{diff}}(pp)/\sigma_{\text{el}}(pp) \approx 0.25\) \[18\]. The effect of the truncation of the sum over the intermediate states has been evaluated by direct \(z\)-integration in eq. \((7)\). This procedure is more accurate than the effective diffraction matrix approximation used in earlier works \[4, 6\]. The main difference is that the direct evaluation of the \(z\)-integrations in eq. \((7)\) produces a somewhat weaker suppression of the contribution of heavy intermediate states and a somewhat faster onset of CT effects. Since our diffraction matrix gives a reasonable description of diffractive \(pN\) scattering, we believe that the multiple scattering approach employed in this paper is a good tool for a quantitative study of the onset of CT in \((e,e'p)\) reactions.

Besides the matrix \(\hat{\sigma}\), the evaluation of \(w(\vec{p})\) requires the initial wave function \(|E\rangle\). Assuming the dominance of the hard mechanism in the electromagnetic form factors \[19\], it can be shown that the projection of the ejectile onto the subspace of hadronic mass eigenstates which satisfy the coherency constraint \(m^2 - m^2_\rho \lesssim Q^2/R_A m_p\), coincides with the similar projection of a compact state, with transverse size \(\rho \sim 1/Q\) \[20, 17\]. For this reason, the latter can be taken for \(|E\rangle\), and in our analysis we use the initial wave function of the form \(\langle \rho | E \rangle \propto \exp(-C \rho^2 Q^2)\), with \(C = 1\). It is worth noting that the missing momentum distribution is not sensitive to the specific choice of \(C\) as long as
\[ C \gtrsim 1/Q^2 \rho_o^2, \] where \( \rho_o \) denotes the position of the first node in the wave functions of the excited states satisfying the coherency requirement. In the region of \( 2 \lesssim Q^2 \lesssim 20 \text{ GeV}^2 \), that we discuss in the present paper, the numerical results are practically independent of the parameter \( C \) for \( C \gtrsim 0.05 \).

Before discussing our numerical results, a comment on the treatment of \( A_z \) of refs.\{3, 7, 21\} is in order. First, these authors have not accounted for the large contribution to \( A_z \) coming from the nonzero \( \alpha_{pN} \). Second, we disagree with the treatment of the initial state \( |E\rangle \) suggested in ref.\{3\}. The formalism of the coupled channel eikonal wave equations used in \{3, 4, 21\} is equivalent to our coupled channel multiple scattering series of eq.(7). However, it has to be noticed that, within the Plane Wave Impulse Approximation (PWIA), eq. (7) combined with our definition of the ejectile wave function \( |E\rangle \) gives
\[
d\sigma(A(e, e'p))/d\sigma(A(e, e'p)) = B^2_i(\vec{p})/B^2_p(\vec{p}),
\]
and \( n(\vec{p}) \) is the nucleon momentum distribution. Eq. (10) shows the standard kinematical correlation between the mass of the excited baryon state \( |i\rangle \) and the missing momentum, which can produce either enhancement or suppression of the resonance abundance in the final state. The procedure used in refs.\{3, 4, 21\} amounts to a redefinition of the ejectile state
\[
|E(\vec{p})\rangle = \sum_i G_{ip}(Q) \frac{B_i(\vec{p})}{B_p(\vec{p})} |i\rangle,
\]
leading to a \( \vec{p} \) dependence in ejectile wave function \( |E\rangle \) that is not born out by the quantum-mechanical treatment of the Fermi motion. It corresponds to a double counting and evidently does not give the correct PWIA limit of the production cross sections. Consequently, the validity of the numerical results of refs. \{3, 4, 21\} appears to be questionable even within the crude models employed to describe the diffraction matrix.

The main results of our work are presented in Fig. 2. Following \{4\} we have calculated \( A_z(x, y) \) assuming parallel kinematics, i.e. taking in eq.(11) \( N_+ = \int_x^y dp_z w(\vec{p}_\perp = 0, p_z) \) and \( N_- = \int_{-\infty}^{-x} dp_z w(\vec{p}_\perp = 0, p_z) \). Fig. 2 shows \( A_z \) calculated for two kinematical windows,
\[(x, y) = (0, 200) \text{ and } (x, y) = (50, 200) \text{ MeV, and two different targets, } ^{16}O \text{ and } ^{40}Ca. \] The parameters of the \textit{pp}- and \textit{pn}-amplitudes were taken from the recent review of ref. \cite{22}. The number of the included excited states and the off-diagonal rescatterings were equal to 4 and 3 respectively. Contributions from higher excitations and rescatterings with \( \nu > 3 \) have been found to be negligible at \( Q^2 \ll 20 \text{ GeV}^2 \). In the region of \( Q^2 \ll 5 \text{ GeV}^2 \) the effect of CT is almost exhausted by the first excitation. To illustrate the contribution of the elastic intermediate state to \( A_z \), the results obtained using the Glauber model are also shown in Fig. 2. In this case the only source of the \( z \)-asymmetry of the missing momentum distribution is the nonzero \( \alpha_{pN} \). As one can see from Fig. 2, in the region of \( Q^2 \sim 2 - 5 \text{ GeV}^2 \), relevant to the experimental study of CT effects at CEBAF, the inelastic intermediate states are only responsible for 25-30 \% (for \( ^{16}O \)) and 10-15 \% (for \( ^{40}Ca \)) of \( A_z \) for both the kinematical windows in \( p_z \). The large contribution of the elastic rescatterings to \( A_z \) can make it difficult to extract definite conclusions on the onset of CT from measurements of the magnitude of \( A_z \) at CEBAF. According to our results, the situation is particularly unfavourable in the case of heavy targets. However, in light nuclei the \( Q^2 \)-dependence of \( A_z \) is dominated by CT.

Since CT effects on \( A_z \) are small, assessing the possiblility of their observation at CEBAF energies requires a study of the sensitivity of the theoretical predictions to the unknown ratios between the real and the imaginary parts of the resonance-nucleon amplitudes. At \( Q^2 \ll 10 \text{ GeV}^2 \), our calculations show a weak sensitivity of \( A_z \) to variations of \( \alpha_{iN} \) for the diagonal \( iN \to iN \) rescatterings of the excited states \( i \). However, the effect of the uncertainty associated with the \( \alpha(iN \to kN) \)'s is not negligible. The hatched area In Fig. 2 shows the band of variation of \( A_z \) corresponding to variations of \( \alpha(iN \to kN) \) between -0.5 and 0.5. It appears that, at least for nuclear mass number \( A \sim 10 \), the uncertainty of the theoretical predictions does not rule out the possibility of using the \( Q^2 \) dependence of \( A_z \) in the \( Q^2 \) range attainable at CEBAF to explore the onset of CT.

In conclusion, we have performed a MST calculation of the longitudinal asymmetry \( A_z \) in \((e, e'p)\) scattering. Our analysis improves upon the previous works on this problem in
several aspects. We have for the first time taken into account the background asymmetry induced by the nonvanishing ratio between the real and the imaginary part of the proton- and resonance-nucleon amplitudes, which turns out to be very important and dominates the asymmetry at moderate $Q^2$. Moreover, we have developed a consistent treatment of the ejectile state, which is also free of problems of double counting. The onset of CT, connected with the occurrence of inelastic rescatterings, has been studied using a realistic model of the diffraction scattering matrix, without making approximations in the treatment of coherency effects in the coupled channel multiple scattering series. Our results suggest that the $Q^2$ dependence of $A_z$, if measured in high statistics experiments off light targets, can be used as a signature of the onset of CT in the kinematical range accessible at CEBAF.

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Figure captions:

Fig. 1 - The typical diagrams contributing to the FSI factor $\Phi(\vec{r}_1, \vec{r}'_1)$ in the MST: (a) the diagram without the interaction between the two trajectories outgoing from $\vec{r}_1$ and $\vec{r}'_1$, (b) the diagram containing the interaction between the trajectories generated by the term $\hat{\Gamma}(b_j)\hat{\Gamma}^*(b'_j)$.

Fig. 2 - The $Q^2$-dependence of the longitudinal asymmetry $A_z(x, y)$ as given by Eq. (1) in the case of parallel kinematics $p_\perp = 0$ for $^{16}O(e, e'p)$ and $^{40}Ca(e, e'p)$ scattering. The solid lines show the results obtained with inclusion of the elastic and inelastic intermediate states within the coupled-channel MST, while the dashed lines were obtained in the one-channel Glauber model. The shaded areas show the limits of variations of $A^{MST}_z$ for variations of $\alpha(iN \rightarrow kN)$ between -0.5 and 0.5.
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