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FINAL-STATE INTERACTIONS AND SINGLE-SPIN ASYMMETRIES IN SEMI-INCLUSIVE DEEP INELASTIC SCATTERING*

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Recent measurements from the HERMES and SMC collaborations show a remarkably large azimuthal single-spin asymmetries A_{UL} and A_{UT} of the proton in semi-inclusive pion leptoproduction $\gamma^*(q)p \to \pi X$. We show that final-state interactions from gluon exchange between the outgoing quark and the target spectator system leads to single-spin asymmetries in deep inelastic lepton-proton scattering at leading twist in perturbative QCD; i.e., the rescattering corrections are not power-law suppressed at large photon virtuality q^2 at fixed x_{bj}. The existence of such single-spin asymmetries requires a phase difference between two amplitudes coupling the proton target with $J^z_p = \pm\frac{1}{2}$ to the same final-state, the same amplitudes which are necessary to produce a nonzero proton anomalous magnetic moment. We show that the exchange of gauge particles between the outgoing quark and the proton spectators produces a Coulomb-like complex phase which depends on the angular momentum L^z of the proton's constituents and thus is distinct for different proton spin amplitudes. The single-spin asymmetry which arises from such final-state interactions does not factorize into a product of structure function and fragmentation function, and it is not related to the transversity distribution $\delta q(x,Q)$ which correlates transversely polarized quarks with the spin of the transversely polarized target nucleon.

Single-spin asymmetries in hadronic reactions have been among the most difficult phenomena to understand from basic principles in QCD. The problem has become more acute because of the observation by the HERMES[1] and SMC[2] collaborations of a strong correlation between the target proton spin $\vec{S}_p$ and the plane

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of a produced pion in semi-inclusive deep inelastic lepton scattering $\ell p^\uparrow \to \ell' \pi X$ at photon virtuality as large as $Q^2 = 6$ GeV^2. Large azimuthal single-spin asymmetries have also been seen in hadronic reactions such as $pp^\uparrow \to \pi X$,[3] where the target antiproton is polarized normal to the pion production plane, and in $pp \to \Lambda^\uparrow X$,[4] where the hyperon is polarized normal to the Λ production plane.

In the target rest frame, single-spin correlations correspond to the $T-odd$ triple product $i\vec{S}_p \cdot \vec{p}_\pi \times \vec{q}$, where the phase i is required by time-reversal invariance. The differential cross section thus has an azimuthal asymmetry proportional to $|\vec{p}_\pi||\vec{q}|\sin\theta_{q\pi}\sin\phi$ where ϕ is the angle between the plane containing the photon and pion and the plane containing the photon and proton polarization vector $\vec{S}_p$. In a general frame, the azimuthal asymmetry has the invariant form $\frac{i}{M}\epsilon_{\mu\nu\sigma\tau}P^\mu S_p^\nu p_\pi^\sigma q^\tau$ where the polarization four-vector of the proton satisfies $S^2 = -1$ and $S \cdot P = 0$.

In order to produce a correlation involving a transversely polarized proton, there are two necessary conditions: (1) There must be two proton spin amplitudes $M[\gamma^* p(J^z_p) \to F]$ with $J^z_p = \pm\frac{1}{2}$ which couple to the same final-state $|F>$; and (2) The two amplitudes must have different, complex phases. The analysis of single-spin asymmetries thus requires an understanding of QCD at the amplitude level, well beyond the standard treatment of hard inclusive reactions based on the factorization of structure functions and fragmentation functions. Since we need the interference of two amplitudes which have different proton spin $J^z_p = \pm\frac{1}{2}$ but couple to the same final-state, the orbital angular momentum of the two proton wavefunctions must differ by $\Delta L^z = 1$. The anomalous magnetic moment for the proton is also proportional to the interference of amplitudes $M[\gamma^* p(J^z_p) \to F]$ with $J^z_p = \pm\frac{1}{2}$ which couple to the same final-state $|F>$.

Final-state interactions (FSI) in gauge theory can affect deep inelastic scattering reactions in a profound way, as has been demonstrated recently.[5] The rescattering of the outgoing quark leads to a leading twist contribution to the deep inelastic cross section from diffractive channels $\gamma^* p \to q\bar{q}p'$, and the interference effects induced by these diffractive channels cause nuclear shadowing. Here we shall show that FSI also provide the required phases needed to produce single-spin asymmetries in deep inelastic scattering.

The dynamics of the constituents in the target can be described by its light-front wavefunctions, $\psi_{n/p}(x_i, \vec{k}_{\perp i}, \lambda_i)$, the projections of the hadronic eigenstate on the free color-singlet Fock state $|n>$ at a given light-cone time $\tau = t + z/c$. The wavefunctions are Lorentz-invariant functions of the relative coordinates $x_i = k^+_i/P^+_\pi = (k^0 + k^z_i)/(P^0 + P^z)$ and $\vec{k}_{\perp i}$ [with $\sum_{i=1}^n x_i = 1$ and $\sum_{i=1}^n \vec{k}_{\perp i} = \vec{0}_\perp$], and they are independent of the bound state's physical momentum P^+ and $\vec{P}_\perp$.[6] The physical transverse momenta are $\vec{p}_{\perp i} = x_i \vec{P}_\perp + \vec{k}_{\perp i}$. The λ_i label the light-front spin S^z projections of the quarks and gluons along the quantization z direction. If a target is stable, its light-front wavefunction must be real. Thus the only source of a nonzero complex phase in leptoproduction in the light-front frame are final-state interactions. The rescattering corrections from final-state exchange of gauge particles produce Coulomb-like complex phases which, however, depend on the proton spin. Thus $M[\gamma^* p(J_p^z = \pm\frac{1}{2}) \to F] = |M[\gamma^* p(J_p^z = \pm\frac{1}{2}) \to F]|\, e^{i\chi_\pm}$. Each of the phases is infrared divergent; however the difference $\Delta\chi = \chi_+ - \chi_-$ is infrared finite and nonzero. The resulting single-spin asymmetry is then proportional to $\sin\Delta\chi$.

We shall calculate the single-spin asymmetry in semi-inclusive electroproduction $\gamma^* p \to HX$ induced by final-state interactions in a model of a spin-$\frac{1}{2}$ proton of mass M composed of charged spin-$\frac{1}{2}$- spin-0 constituents of mass m and λ, respectively, as in the QCD-motivated quark-diquark model of a nucleon. The basic electroproduction reaction is then $\gamma^* p \to q(qq)_0$, as illustrated in Figs. 1 and 2. We shall take the case where the detected particle H is identical to the quark. One can compute the asymmetry for a detected hadron by convoluting the jet asymmetry result with a realistic fragmentation function; e.g. $D_{q\to\pi X}(z, Q^2)$.

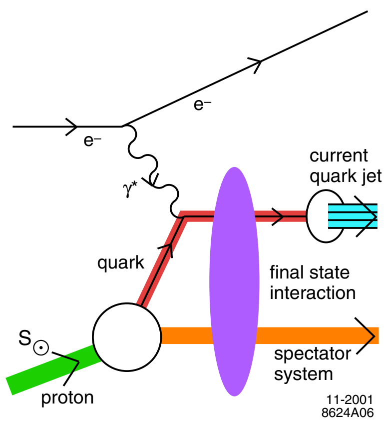

Fig. 1. The final-state interaction in the semi-inclusive deep inelastic lepton scattering $\ell p^\uparrow \to \ell' \pi X$.

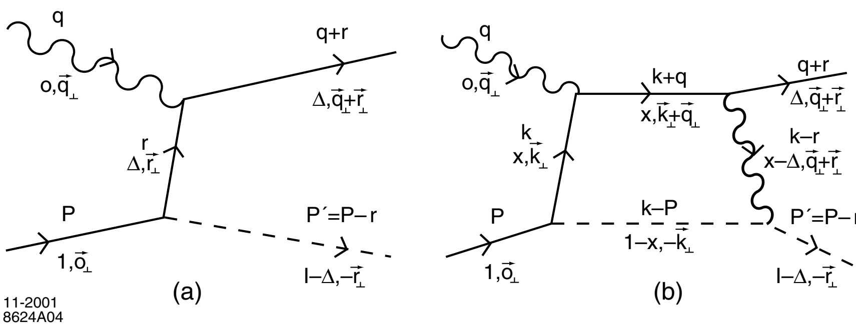

Fig. 2. The tree (a) and one-loop (b) graphs for $\gamma^* p \to q(qq)_0$. The interference of the two amplitudes with $J_p^z = \pm 1/2$ provides the proton's single-spin asymmetry.

The amplitude for the $\gamma^* p \to q(qq)_0$ can be computed from the tree and one-loop graphs illustrated in Fig. 2. A single-spin asymmetry will arise from the final-state interactions of the outgoing charged lines. The $J^z = +\frac{1}{2}$ two-particle Fock state is given by[7,8]

$$
\begin{aligned}
&\left|\Psi^{\uparrow}_{\rm two\ particle}(P^+, \vec{P}_\perp = \vec{0}_\perp)\right\rangle \\
&= \int \frac{{\rm d}^2\vec{k}_\perp {\rm d}x}{\sqrt{x(1-x)}16\pi^3} \Big[\ \psi^{\uparrow}_{+\frac{1}{2}}(x,\vec{k}_\perp) \left|+\frac{1}{2};\ xP^+\ ,\ \vec{k}_\perp\right\rangle \\
&\quad + \psi^{\uparrow}_{-\frac{1}{2}}(x,\vec{k}_\perp) \left|-\frac{1}{2};\ xP^+\ ,\ \vec{k}_\perp\right\rangle\ \Big]\ ,
\end{aligned}
\tag{1}
$$

where

$$
\left\{
\begin{array}{l}
\psi^{\uparrow}_{+\frac{1}{2}}(x,\vec{k}_\perp) = (M + \frac{m}{x})\,\varphi\ , \\
\psi^{\uparrow}_{-\frac{1}{2}}(x,\vec{k}_\perp) = -\frac{(+k^1 + {\rm i}k^2)}{x}\,\varphi\ .
\end{array}
\right.
\tag{2}
$$

The scalar part of the wavefunction φ depends on the dynamics. In the perturbative theory it is simply

$$
\varphi = \varphi(x,\vec{k}_\perp) = \frac{\frac{g}{\sqrt{1-x}}}{M^2 - \frac{\vec{k}_\perp^2 + m^2}{x} - \frac{\vec{k}_\perp^2 + \lambda^2}{1-x}}\ . \tag{3}
$$

In general one normalizes the Fock state to unit probability.

Similarly, the $J^z = -\frac{1}{2}$ two-particle Fock state has components

$$
\left\{
\begin{array}{l}
\psi^{\downarrow}_{+\frac{1}{2}}(x,\vec{k}_\perp) = \frac{(+k^1 - {\rm i}k^2)}{x}\,\varphi\ , \\
\psi^{\downarrow}_{-\frac{1}{2}}(x,\vec{k}_\perp) = (M + \frac{m}{x})\,\varphi\ .
\end{array}
\right.
\tag{4}
$$

The spin-flip amplitudes in (2) and (4) have corresponding orbital angular momentum projections $L^z = +1$ and -1. The numerator structure of the wavefunctions is characteristic of the orbital angular momentum and is the same for both perturbative and non-perturbative couplings.

We require the interference between the tree amplitude of Fig. 2a and the one-loop amplitude of Fig. 2b. The contributing amplitudes for $\gamma^* p \to q(qq)_0$ have the following structure through one loop order:

$$
\mathcal{A}(\Uparrow \to \uparrow) = (M + \frac{m}{\Delta})\ C\ (h + i\frac{e_1 e_2}{8\pi} g_1) \tag{5}
$$

$$
\mathcal{A}(\Downarrow \to \uparrow) = \ (\frac{+r^1 - ir^2}{\Delta})\ \ C\ (h + i\frac{e_1 e_2}{8\pi} g_2) \tag{6}
$$

$$
\mathcal{A}(\Uparrow \to \downarrow) = \ (\frac{-r^1 - ir^2}{\Delta})\ \ C\ (h + i\frac{e_1 e_2}{8\pi} g_2) \tag{7}
$$

$$
\mathcal{A}(\Downarrow \to \downarrow) = (M + \frac{m}{\Delta})\ C\ (h + i\frac{e_1 e_2}{8\pi} g_1)\ , \tag{8}
$$

where

$$
C = -\ g\ e_1\ P^+\ \sqrt{\Delta}\ 2\ \Delta\ (1-\Delta) \tag{9}
$$

$$
h = \frac{1}{\vec{r}_\perp^{\,2} + \Delta(1-\Delta)(-M^2 + \frac{m^2}{\Delta} + \frac{\lambda^2}{1-\Delta})}\ . \tag{10}
$$

The quark light-cone fraction $\Delta = \frac{k^+}{P^+}$ is equal to the Bjorken variable x_{bj} up to corrections of order $1/Q$. The label $\Uparrow / \Downarrow$ corresponds to $J^z_p = \pm\frac{1}{2}$. The second label $\uparrow / \downarrow$ gives the spin projection $J^z_q = \pm\frac{1}{2}$ of the spin-$\frac{1}{2}$ constituent. Here e_1 and e_2 are the electric charges of q and $(qq)_0$, respectively, and g is the coupling constant of the proton-q-$(qq)_0$ vertex. The first term in (5) to (8) is the Born contribution of the tree graph. The crucial result will be the fact that the contributions g_1 and g_2 from the one-loop diagram Fig. 2b are different, and that their difference is infrared finite. A gauge particle mass λ_g will be used as an infrared regulator.

The covariant expression for the four one-loop amplitudes of diagram Fig. 2b is:

$$\begin{aligned}
&\mathcal{A}^{\rm one-loop}(I) \qquad\qquad (11)\\
&= ig\ e_1^2\ e_2 \int \frac{d^4k}{(2\pi)^4} \\
&\times \frac{\mathcal{N}(I)}{(k^2-m^2+i\epsilon)\ ((k+q)^2-m^2+i\epsilon)((k-r)^2-\lambda_g^2+i\epsilon)((k-P)^2-\lambda^2+i\epsilon)} \\
&= -ig\ e_1^2\ e_2 \int \frac{d^2\vec{k}_\perp}{2(2\pi)^4} \int P^+ dx\ \frac{\mathcal{N}(I)}{P^{+4}\ x\ x\ (x-\Delta)\ (1-x)} \\
&\times \int dk^- \frac{1}{\left(k^- - \frac{(m^2+\vec{k}_\perp^2)-i\epsilon}{xP^+}\right)\left((k^-+q^-) - \frac{(m^2+(\vec{k}_\perp+\vec{q}_\perp)^2)-i\epsilon}{xP^+}\right)} \\
&\times \frac{1}{\left((k^- - r^-) - \frac{(\lambda_g^2+(\vec{k}_\perp-\vec{r}_\perp)^2)-i\epsilon}{(x-\Delta)P^+}\right)\left((k^- - P^-) + \frac{(\lambda^2+\vec{k}_\perp^2)-i\epsilon}{(1-x)P^+}\right)},
\end{aligned}$$

where we used $k^+ = xP^+$. The numerators $\mathcal{N}(I)$ are given by

$$\mathcal{N}(\Uparrow\to\uparrow) = 2P^+\sqrt{\Delta}\ x\ (M + \frac{m}{x})\ q^- \qquad (12)$$

$$\mathcal{N}(\Downarrow\to\uparrow) = 2P^+\sqrt{\Delta}\ x\ (+k^1 - ik^2)\ q^- \qquad (13)$$

$$\mathcal{N}(\Uparrow\to\downarrow) = 2P^+\sqrt{\Delta}\ x\ (-k^1 - ik^2)\ q^- \qquad (14)$$

$$\mathcal{N}(\Downarrow\to\downarrow) = 2P^+\sqrt{\Delta}\ x\ (M + \frac{m}{x})\ q^-\ , \qquad (15)$$

where $q^- = \frac{Q^2}{\Delta P^+} = \frac{2M\nu}{P^+}$. The integration over k^- in (11) does not give zero only if $0 \le x \le 1$. We first consider the region $\Delta < x \le 1$.

$$\begin{aligned}
&\mathcal{A}^{\rm one-loop}(I) \qquad\qquad (16)\\
&= -ig\ e_1^2\ e_2\ \times\ (2\pi i) \int \frac{d^2\vec{k}_\perp}{2(2\pi)^4} \int P^+ dx\ \frac{\mathcal{N}(I)}{P^{+4}\ x\ x\ (x-\Delta)\ (1-x)} \\
&\times \frac{1}{\left(P^- - \frac{(\lambda^2+\vec{k}_\perp^2)-i\epsilon}{(1-x)P^+} - \frac{(m^2+\vec{k}_\perp^2)-i\epsilon}{xP^+}\right)\left(P^- - \frac{(\lambda^2+\vec{k}_\perp^2)-i\epsilon}{(1-x)P^+} + q^- - \frac{(m^2+(\vec{k}_\perp+\vec{q}_\perp)^2)-i\epsilon}{xP^+}\right)} \\
&\times \frac{1}{\left(P^- - \frac{(\lambda^2+\vec{k}_\perp^2)-i\epsilon}{(1-x)P^+} - r^- - \frac{(\lambda_g^2+(\vec{k}_\perp-\vec{r}_\perp)^2)-i\epsilon}{(x-\Delta)P^+}\right)},
\end{aligned}$$

The result is identical to that obtained from light-cone time-ordered perturbation theory.

The phases χ_i needed for single-spin asymmetries come from the imaginary part of (16), which arises from the potentially real intermediate state allowed before the rescattering. The imaginary part of the second propagator (light-cone energy denominator) in (16) gives

$$-i\pi\ \delta\ \left(P^- - \frac{(\lambda^2+\vec{k}_\perp^2)}{(1-x)P^+} + q^- - \frac{(m^2+(\vec{k}_\perp+\vec{q}_\perp)^2)}{xP^+}\right)$$
$$= \ -i\pi\ \frac{1}{P^+}\ \frac{\Delta^2}{\vec{q}_\perp^{\,2}}\ \delta(x\ -\ \Delta\ -\ \bar{\delta})\ , \tag{17}$$

where

$$\bar{\delta}\ =\ 2\ \Delta\ \frac{\vec{q}_\perp\cdot(\vec{k}_\perp-\vec{r}_\perp)}{\vec{q}_\perp^{\,2}}\ . \tag{18}$$

For the current-gauge propagator-current factor, in Feynman gauge only the $-g^{+-}$ term of the gauge propagator $-g^{\mu\nu}$ contributes in the Bjorken limit, and it provides a factor proportional to q^- in the numerator which cancels the q^- in the denominator of the gauge propagator. Therefore the result scales in the Bjorken limit. We have verified that the result is the same in the light cone gauge. The small numerator coupling of the light-cone gauge particle is compensated by the small value for the exchanged $l^+ = \bar{\delta}P^+$ momentum.

Since the exchanged momentum $\bar{\delta}P^+$ is small, the light-cone energy denominator corresponding to the gauge propagator is dominated by the $\frac{(k_\perp-r_\perp)^2+\lambda_p^2}{(x-\Delta)}$ term. This gets multiplied by $(x-\Delta)$, so only $(k_\perp-r_\perp)^2+\lambda_p^2$ appears in the propagator, independent of whether the photon is absorbed or emitted. The contribution from the region $0 \le x < \Delta$ is thus the same as that from the region $\Delta < x \le 1$.

We can integrate (16) over the transverse momentum using a Feynman parametrization to obtain the one-loop terms in (5) to (8):

$$g_1 = \int_0^1 d\alpha\ \frac{1}{\alpha(1-\alpha)\vec{r}_\perp^{\,2}+\alpha\lambda_g^2+(1-\alpha)\Delta(1-\Delta)(-M^2+\frac{m^2}{\Delta}+\frac{\lambda^2}{1-\Delta})} \tag{19}$$

$$g_2 = \int_0^1 d\alpha\ \frac{\alpha}{\alpha(1-\alpha)\vec{r}_\perp^{\,2}+\alpha\lambda_g^2+(1-\alpha)\Delta(1-\Delta)(-M^2+\frac{m^2}{\Delta}+\frac{\lambda^2}{1-\Delta})}\ . \tag{20}$$

The final-state interactions generate a phase when exponentiated. The rescattering phases $e^{i\chi_i}$ $(i=1,2)$ with $\chi_i = \tan^{-1}(\frac{e_1e_2}{8\pi}\frac{g_i}{h})$ are thus distinct for the spin-parallel and spin-antiparallel amplitudes. The difference in phase arises from the orbital angular momentum $k_\perp$ factor in the spin-flip amplitude, which after integration gives the extra factor of the Feynman parameter α in the numerator of g_2. Notice that the phases χ_i are each infrared divergent for zero gauge boson mass $\lambda_g \to 0$, as is characteristic of Coulomb phases. However, the difference $\chi_1-\chi_2$ which contributes to the single-spin asymmetry is infrared finite.

The virtual photon and produced hadron define the production plane which we will take as the $\hat{z}-\hat{x}$ plane. The azimuthal single-spin asymmetry transverse to the production plane is given by

$$\mathcal{P}_y = \frac{e_1 e_2}{8\pi} \, \frac{2\left(\Delta M + m \right) r^1}{\left[\left(\Delta M + m \right)^2 + \vec{r}_\perp^2 \right]} \left[\vec{r}_\perp^2 + \Delta(1-\Delta)(-M^2 + \frac{m^2}{\Delta} + \frac{\lambda^2}{1-\Delta}) \right]$$
$$\times \; \frac{1}{\vec{r}_\perp^2} \ln \frac{\vec{r}_\perp^2 + \Delta(1-\Delta)(-M^2 + \frac{m^2}{\Delta} + \frac{\lambda^2}{1-\Delta})}{\Delta(1-\Delta)(-M^2 + \frac{m^2}{\Delta} + \frac{\lambda^2}{1-\Delta})} \; . \qquad (21)$$

The linear factor of $r^1 = r^x$ reflects the fact that the single-spin asymmetry is proportional to $\vec{S} \cdot \vec{q} \times \vec{r}$ where $\vec{q} \sim -\nu\hat{z}$ and $\vec{S}_p = \pm\hat{y}$. Here $\Delta = x_{bj}$.

Our analysis can be easily generalized to QCD. The coupling strength $\frac{e_1 e_2}{4\pi}$ of the final-state interaction in the Abelian model becomes $C_F \alpha_s(\mu^2)$ in QCD, where the scale of α_s in the $\overline{MS}$ scheme can be identified with the momentum transfer carried by the gluon $\mu^2 = e^{-5/3}(k-r)^2_\perp$.[9] The numerator structure of the Abelian and QCD wavefunctions are the same since the matrix elements coupling the proton to its constituents are determined by the orbital angular momentum. The strength of the matrix elements can be normalized by the anomalous magnetic moment and the total charge.

In QCD, $r_\perp$ is the magnitude of the momentum of the current quark jet relative to the q direction. Notice that for large r^1, $\mathcal{P}_y$ decreases as $\alpha_s(r^2_\perp) x_{bj} M r_\perp [\ln r^2_\perp]/\vec{r}^2_\perp$. The mass M of the physical proton mass appears here since it determines the ratio of the $L_z = 1$ and $L_z = 0$ matrix elements.

We thus predict that the single-spin asymmetry in electroproduction is independent of Q^2 at fixed $\Delta = x_{bj}$. It satisfies Bjorken scaling and is therefore leading twist. The existence of the single-spin asymmetries requires a phase difference between two amplitudes coupling the proton target with $J^z_p = \pm\frac{1}{2}$ to the same final-state, the same amplitudes which are necessary to produce a nonzero proton anomalous magnetic moment. We have shown that the exchange of gauge particles between the outgoing quark and the proton spectators produces a Coulomb-like complex phase which depends on the angular momentum L^z of the proton's constituents and is thus distinct for different proton spin amplitudes. Unlike the Coulomb phase, the phase difference which appears in the single-spin asymmetry is infrared finite.

In conclusion, we have shown that the final-state interactions from gluon exchange between the outgoing quark and the target spectator system leads to single-spin asymmetries in deep inelastic lepton-proton scattering at leading twist in perturbative QCD; i.e., the rescattering corrections are not power-law suppressed at large photon virtuality q^2 at fixed x_{bj}. The nominal size of the spin asymmetry is thus $C_F \alpha_s(r^2_\perp) a_p$ where a_p is the proton anomalous magnetic moment. [We have estimated the scale of α_s as $\mathcal{O}(r^2_\perp)$.] Our result is directly applicable to the azimuthal correlation of the proton spin with the virtual photon to current quark jet plane, which can be deduced from jet measures such as the thrust distribution.

The $\sin\phi$ correlation of the proton spin with the photon-to-pion production plane as measured in the HERMES and SMC experiments can then be obtained using the usual fragmentation function. Detailed comparisons with experiment will be presented elsewhere.

It is usually assumed that the cross section for semi-inclusive deep inelastic scattering at large Q^2 factorizes as the product of quark distributions times quark fragmentation functions.[10,11] Our analysis shows that the single-spin asymmetry which arises from final-state interactions does not factorize in this way since the result depends on the $< p|\bar{\psi}_q A\psi|p >$ proton correlator, not the usual quark distribution derived from $< p|\bar{\psi}_q(\xi)\psi_q(0)|p >$ evaluated at equal light-cone time $\xi^+ = 0$. In particular, the spin asymmetry is not related to the transversity distribution $\delta q(x, Q)$ which correlates transversely polarized quarks with the spin of the transversely polarized target nucleon.

References

1. HERMES Collaboration, A. Airapetian *et al.*, *Phys. Rev. Lett.* **84** 4047 (2000); *Phys. Rev.* **D64**, 097101 (2001).
2. A. Bravar, for the SMC Collaboration, *Nucl. Phys. B* (Proc. Suppl.) **79**, 520 (1999).
3. E704 Collaboration, A. Bravar *et al.*, *Phys. Rev. Lett.* **77**, 2626 (1996).
4. K. Heller, in Proceedings of Spin 96, C. W. de Jager, T. J. Ketel and P. Mulders, Eds., World Scientific (1997).
5. S. J. Brodsky, P. Hoyer, N. Marchal, S. Peigne and F. Sannino, hep-ph/0104291.
6. G. P. Lepage and S. J. Brodsky, *Phys. Rev.* **D22**, 2157 (1980); S. J. Brodsky and G. P. Lepage, in: *Perturbative Quantum Chromodynamics*, edited by A. H. Mueller (World Scientific, Singapore 1989).
7. S. J. Brodsky and S. D. Drell, *Phys. Rev.* **D22**, 2236 (1980).
8. S. J. Brodsky, D. S. Hwang, B. Q. Ma and I. Schmidt, *Nucl. Phys. B* **593**, 311 (2001).
9. S. J. Brodsky, A. H. Hoang, J. H. Kuhn and T. Teubner, *Phys. Lett. B* **359**, 355 (1995).
10. J. C. Collins, *Nucl. Phys. B* **396**, 161 (1993).
11. D. Boer and P. J. Mulders, *Phys. Rev.* **D57**, 5780 (1998).