The Beam Single Spin Asymmetry in Semi-inclusive Deep
Inelastic Scattering

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Abstract

We study the beam single spin asymmetry in semi-inclusive hadron production in deep inelastic
scattering in at order $1/Q$. There are two competing contributions: the leading order transverse
momentum dependent parton distribution $h_{1}^{+}(x, k_{\perp})$ convoluted with chiral-odd fragmentation
function $\hat{e}(z)$, and the chiral-odd distribution function $e(x)$ convoluted with Collins fragmentation
function $H_{1}^{+}(z, k_{\perp}')$. We estimate this asymmetry and compare with the experimental measurements from CLAS and HERMES collaborations.

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The single spin asymmetry (SSA) is a novel phenomena in high energy spin physics, and has attracted much interest in recent years [1]. In particular, the measurements from the HERMES, SMC, and JLAB collaborations show a remarkably large SSA in the semi-inclusive deep inelastic scattering, such as pion production in $\gamma^* p \rightarrow \pi X$, when the proton is polarized transversely to the direction of the virtual photon [2, 3, 4]. On the theoretical side, there are many approaches to understanding SSA using Quantum Chromodynamics (QCD) phenomenology [1, 5]. Recent interest focuses on the transverse momentum dependent (TMD) parton distributions and their implications to the semi-inclusive processes in deep inelastic scattering [6, 7, 8, 9, 10, 11, 12, 13]. For example, the Sivers function is one of those TMD parton distributions representing the asymmetric distribution of quarks in a transversely polarized proton, which correlates the quark transverse momentum and the proton polarization vector $\vec{S}_\perp$ [7]. It contributes to the target SSA in semi-inclusive deep inelastic scattering. The existence of the Sivers function has been confirmed recently [14, 15, 16, 17], where the final state interactions from the gauge link in the gauge invariant definition of TMD parton distributions play an important role.

In this paper, we will study the SSA related to the beam polarization. Unlike the target SSA, the beam SSA is subleading in $1/Q$, which will eventually vanish as $Q^2 \rightarrow \infty$. However, at some intermediate $Q^2$, but still large enough to guarantee the factorization, this asymmetry might be important and measurable. Experimentally, HERMES collaboration found this asymmetry consistent with zero [2], but the CLAS collaboration at JLab found sizable asymmetry, in the order of 4% [18]. In the literature [10, 11, 19], this asymmetry has been associated with the twist-3 chiral-odd distribution function $e(x)$ [20] convoluted with the Collins fragmentation function $H_{1\perp}(z,k_{1\perp})$ [8]. However, this is not the complete picture at this order. There is an additional contribution: the leading order TMD parton distribution $h_{1\perp}(x,k_{1\perp})$ [12] convoluted with a chiral-odd fragmentation function $\hat{e}(z)$ [9, 21]. We will demonstrate the existence of this contribution below.

The TMD parton distribution $h_{1\perp}$ represents the correlation between the quark’s transverse momentum and polarization in an unpolarized proton state [12, 22]. It has the same features as the Sivers function: it is a leading-order distribution; it is nonvanishing due to the final state interactions; it depends on the quark orbital angular momentum of the nucleon [23]. Various model calculations have also shown that it has a similar size as the Sivers function [24, 25, 26]. On the other hand, since $h_{1\perp}$ is chiral-odd, it is very difficult to probe it in deep inelastic scattering, just like the transversity distribution. Our analysis shows that it contributes to the beam SSA, which can be used to extract the distribution itself. Another possible way to study $h_{1\perp}$ distribution is the asymmetry in the Drell-Yan process [22].

We first derive the beam SSA in deep inelastic scattering. The semi-inclusive hadron production cross section can be expressed as

$$
\frac{d^5\sigma}{dx_Bdydzd^2P_{\perp h}} = \frac{2\pi\alpha^2}{4xzBQ^2s}L_{\mu\nu}W^{\mu\nu},
$$

where $L_{\mu\nu}$ and $W^{\mu\nu}$ are leptonic and hadronic tensors, respectively. We work in a frame where the virtual photon’s momentum $q$ and the proton’s momentum $P$ are in the $z$ direction, and the incident and outgoing lepton’s momenta $\vec{l}$ and $\vec{l}'$ form a scattering plane. We can define the azimuthal angle of any momentum as an angle relative to the scattering plane. The variable $s$ is the lepton-hadron total energy square, $Q^2 = -q^2$ the virtuality of the photon, and the dimensionless variables $x_B$, $y$, and $z$ are defined as $x_B = Q^2/2P \cdot q$, $y = P \cdot q/P \cdot l$, and $z = 2P \cdot q/P \cdot l$. The azimuthal angle of any momentum is an angle relative to the scattering plane.
where $L$ we will focus on the contribution from the $e$ contribution from diagram (a). The antisymmetric part has two contributions. Since the $W$ fragmentation of a quark into a pion. The symmetric part of diagrams represent the parton distribution of the target, and the upper part involves the functions. The relevant Feynman diagrams are shown in Fig. 1. The lower part of these can be parameterized in Lorentz-invariant and gauge-invariant distribution and fragmentation can be calculated in perturbative QCD; and the soft parts are nonperturbative and can be parameterized in Lorentz-invariant and gauge-invariant distribution and fragmentation functions. The relevant Feynman diagrams are shown in Fig. 1. The lower part of these diagrams represent the parton distribution of the target, and the upper part involves the fragmentation of a quark into a pion. The symmetric part of $W^{\mu\nu}$ in the leading order has a contribution from diagram (a). The antisymmetric part has two contributions. Since the contribution from the $e(x)\otimes H^\perp_1(z,k^\perp_1)$ term has been calculated in Ref. [10], in the following we will focus on the contribution from the $h^\perp_1(x,k^\perp_1)\otimes \hat{e}(z)$ term. Because $\hat{e}(z)$ is a twist-3 fragmentation function, we need to include the diagrams (b) and (c) of Fig. 1 to get an electromagnetically gauge-invariant result.

The parton distributions can be defined from the following density matrix [9],
\[
\mathcal{M}_p(x,k_\perp) = p^z \int \frac{d\xi d^2\xi_\perp}{(2\pi)^3} e^{-i(\xi^z-k^z-\xi_\perp k_\perp)} \langle PS|\bar{\psi}(\xi^-) \mathcal{L}^\dagger \mathcal{L} \psi(0)|PS\rangle ,
\]
where $\mathcal{L}$ is the gauge link [13, 16]. The density matrix has the following expansion [12],
\[
\mathcal{M}_p = \frac{1}{2} f_1(x,k_\perp) \gamma^\mu p^\mu + \frac{1}{2M} h^\perp_1(x,k_\perp) \sigma^{\mu\nu} k_\mu p_\nu + \ldots ,
\]
where $f_1$ is the usual unpolarized unintegrated parton distribution, and both $f_1$ and $h^\perp_1$ are leading order in twist counting.

Similarly, for the pion fragmentation functions, we can define the density matrix as [9, 10, 11],
\[
\mathcal{M}_\pi(z,k'_\perp) = \frac{n^-}{2z} \int \frac{d\eta d^2\eta_\perp}{(2\pi)^3} e^{-i(\eta^z-k'^z-\eta_\perp k'_\perp)} \langle 0|\mathcal{L} \psi(\eta^+, \eta_\perp)|\pi X\rangle \langle \pi X|\bar{\psi}(0)\mathcal{L}^\dagger|0\rangle ,
\]
where $z = p_\pi \cdot p/k'_\perp \cdot p$, and $-z k'_\perp$ is the transverse momentum of pion relative to the quark’s momentum. This fragmentation matrix density has the expansion [9],
\[
\mathcal{M}_\pi = \frac{1}{2} \hat{f}_1(z,k'_\perp) \gamma^\mu p^\mu + \frac{1}{2 p_\pi \cdot p} \hat{e}(z,k'_\perp) + \ldots ,
\]
and the anti-symmetric parts, $\mathcal{M}_\pi = \frac{1}{2} M_p \hat{e}(z,k'_\perp) + \ldots$.
where $\hat{f}_1$ is the usual unpolarized unintegrated fragmentation function; and $\hat{e}$ is the twist-$3$ chiral-odd fragmentation function. There are two twist-$3$ and chiral-odd fragmentation functions for the pion [21], but we only keep the one which contributes to beam SSA. As argued in [4, 21], instead of the pion mass, we put the nucleon mass as the coefficient in front of $\hat{e}(z, k^\perp_1)$, because the pion mass vanishes in the chiral limit but the density matrix $M_\pi$ does not.

The contribution from Fig. 1(a) to the hadronic tensor $W^{\mu\nu}$ can be calculated as,

$$W^{\mu\nu(a)} = 2z \int d^2k_\perp d^2k'_\perp \delta^{(2)}(P_{\perp h}/z - k_\perp + k'_\perp) \{ f_1(x, k_\perp)\hat{f}_1(z, k'_\perp) [p^\mu n^\nu + p^\nu n^\mu - g^{\mu\nu}] \\
+ i\hbar \frac{1}{2} (x, k_\perp)\hat{e}(z, k'_\perp) \frac{1}{p_x \cdot p} [p^\mu k^\nu_\perp - p^\nu k^\mu_\perp] \},$$

where the electric charge of the quark and the sum over all quark flavor are implicitly assumed. The first term in the bracket is the symmetric part of the tensor, and the second one antisymmetric. The symmetric part itself is electromagnetic gauge invariant, while the antisymmetric part is not, i.e., $q_\mu W^{\mu\nu(a)} \neq 0$. However, after including the contributions from diagrams (b) and (c) in Fig. 1, we can recover the gauge-invariance [3]. This leads to the following result for the antisymmetric part: $2h_1^+(x, k_\perp)\hat{e}(z, k'_\perp)/z/Q^2 [T^\mu k^\nu_\perp - T^\nu k^\mu_\perp]$, where $T^\mu = 2xP^\mu + q^\mu$.

Including also the contributions from the convolution of $e(x)$ with the Collins function $H^+_1(z, k'_\perp)$ [10], we get the complete result for the antisymmetric part of the hadronic tensor at order $1/Q$,

$$W^a_{\mu\nu} = 2z \int d^2k_\perp d^2k'_\perp \delta^{(2)}(P_{\perp h}/z - k_\perp + k'_\perp) \{ i\hbar \frac{1}{2} (x, k_\perp)\hat{e}(z, k'_\perp) \frac{2}{z} Q^2 [T^\mu k^\nu_\perp - T^\nu k^\mu_\perp] \\
- ixe(x, k_\perp)H^+_1(z, k'_\perp) \frac{2}{Q^2} [T^\mu k^\nu_\perp - T^\nu k^\mu_\perp] \},$$

and it contributes to the beam SSA. We modified the definition of the Collins function $H^+_1$ in [10, 11] by a factor of $M_p/m_\pi$ by the same argument we used in Eq. (13) for the fragmentation function $e(z)$. The definition of $e(x)$ follows [20]. Since $T^\mu$ is on order of $Q$, we see that the above antisymmetric part will contribute to the cross section in the order of $1/Q$. That means the beam SSA will be $1/Q$ suppressed, which is different from the Sivers effect contribution to the target SSA being the leading order effect.

Substituting the hadronic tensor into the differential cross section formula Eq. (10), we will get:

$$\frac{d^5\sigma}{dx_Bdydzd^2P_{\perp h}} = \frac{4\pi\alpha^2}{x_By^2s} \int d^2k_\perp d^2k'_\perp \delta^{(2)}(\frac{P_{\perp h}}{z} - k_\perp + k'_\perp) \left\{ (1 - y + \frac{y^2}{2})f_1(x, k_\perp)\hat{f}_1(z, k'_\perp) \\
+ \lambda_e \frac{2y\sqrt{1-y}}{Q} x e(x, k_\perp)H^+_1(z, k'_\perp) |\vec{k}'_\perp| \sin \phi_k \\
- \lambda_e \frac{2y\sqrt{1-y}}{Q} x e(x, k_\perp)H^+_1(z, k'_\perp) |\vec{k}'_\perp| \sin \phi_k' \right\},$$

(9)
where $\phi_k$ and $\phi_{k'}$ are the azimuthal angles of the momenta $\vec{k}_\perp$ and $\vec{k'}_\perp$ relative to the scattering plane, respectively. For example, we define $\sin \phi_k = \vec{l} \times \vec{P}_{\perp h} / |\vec{l} \times \vec{P}_{\perp h}|$. Since only one of the two transverse momenta $k_\perp$ and $k'_\perp$ is relevant for the $\sin \phi$ asymmetry, we can integrate out the other without assuming the transverse momentum dependence of the distribution and fragmentation functions. After that, the $\phi_k$ or $\phi_{k'}$ dependence leads to $\phi_h$ dependence, where $\phi_h$ is the azimuthal angle of the observed hadron relative to the scattering plane: $\sin \phi_h = \vec{l} \times \vec{P}_{\perp h} / |\vec{l} \times \vec{P}_{\perp h}|$. Finally,

$$
\frac{d^5\sigma}{dx_Bdydzd^2P_{\perp h}} = \frac{4\pi\alpha^2}{x_By^2s} \left\{ (1 - y + \frac{y^2}{2}) \int d^2k_\perp \delta(2)(P_{\perp h} - zk_\perp)f_1(x, k_\perp)\hat{f}_1(z) \\
+ \lambda_a 2y\sqrt{1 - y} Q \int d^2k_\perp \delta(2)(P_{\perp h} - zk_\perp)h_1^{(1/2)}(x, k_\perp)|\vec{k}_\perp| \frac{\hat{e}(z)}{z}\sin \phi_h \\
+ \lambda_e 2y\sqrt{1 - y} \int d^2k'_\perp \delta(2)(P_{\perp h} + zk'_\perp)xe(x)z^2H_1^{(1/2)}(z, k'_\perp)|\vec{k'}_\perp| \sin \phi_h \right\},
$$

where a sign has changed in the last term because $\vec{k'}_\perp$ and $\vec{P}_{\perp h}$ have opposite directions. The integrated fragmentation functions are defined as $f_1(z) = z^2 \int d^2k'_\perp f_1(z, k'_\perp)$, the same for $\hat{e}(z)$, and the distribution function $e(x) = \int d^2k'_\perp e(x, k'_\perp)$.

We can further simplify the differential cross section by integrating out the transverse momentum $P_{\perp h}$ but keeping the dependence on $\phi$,

$$
\frac{d^5\sigma}{dx_Bdydzd\phi} = \frac{2\alpha^2}{x_By^2s} \left\{ (1 - y + \frac{y^2}{2})f_1(z)\hat{f}_1(z) \\
+ 2\lambda_a y\sqrt{1 - y} \frac{M_p}{Q} h_1^{(1/2)}(x)\frac{\hat{e}(z)}{z}\sin \phi \\
+ 2\lambda_e y\sqrt{1 - y} \frac{M_p}{Q} xe(x)H_1^{(1/2)}(z)\sin \phi \right\},
$$

where the integrated parton distribution $f_1(x) = \int d^2k_\perp f_1(x, k_\perp)$, $h_1^{(1/2)}(x) = \int d^2k_\perp |k_\perp|/M_ph_1^{(1/2)}(x, k_\perp)$, and fragmentation $H_1^{(1/2)}(z) = \int d^2k'_\perp |k'_\perp|/M_pH_1^{(1/2)}(z, k'_\perp)$. If we write the differential cross section as $d\sigma \propto 1 + A_y \sin \phi$,

$$
A_y = \frac{\lambda_a \int dydzdx_B \frac{2y\sqrt{1 - y} M_p}{x_By^2s} \left( h_1^{(1/2)}(x)\frac{\hat{e}(z)}{z} + xe(x)H_1^{(1/2)}(z) \right) f_1(z)}{\int dydzdx_B \frac{1 - y + y^2/2}{x_By^2s} f_1(x)\hat{f}_1(z)}.
$$

The $x_B$ and $z$ dependence of $A_y$ can also be similarly calculated. We note that the two contributions have exactly the same dependence on $y$, which makes it difficult to distinguish them experimentally.

With known distribution and fragmentation functions, we can predict the beam SSA. However, up to now, except for the unpolarized quark distribution $f_1$ and fragmentation $\hat{f}_1$, these functions can only be estimated in models. In addition, the model calculations are not consistent at the present stage. For example, controversial predictions exist for the Collins fragmentation function $H_1^{(1/2)}$ [27, 28], and we have a wide range of predictions for the leading-order TMD parton distribution $h_1^{(1/2)}(x, k_\perp)$ from models [24, 25, 26]. So, a reliable model prediction is not possible at present. However, we can still gain some insight for these functions by comparison with the experimental data. For example, in [19], the beam SSA
FIG. 2: The beam SSA prediction from $h_1^\perp(x) \otimes \hat{e}(z)$, compared with the experimental data from CLAS [18]. The $z$ dependence solely comes from the ratio of the fragmentation functions $\hat{e}(z)/f_1(z)$.

has been interpreted as the result of the Collins effect, and the experimental data were used to extract the distribution function $e(x)$.

In this paper, we take an alternative extreme. We interpret the beam SSA as a result of the first term in Eq. (12). To compare with the experimental data, we assume that the factorization works at the energy range covered by the experiment. This contribution depends on the chiral-odd fragmentation function $\hat{e}(z)$, which has been calculated in a chiral quark model in [29]. To a good approximation, we have

$$\hat{e}(z) = \frac{z}{1 - z} \frac{m_q}{M_n} \hat{f}_1(z) \approx \frac{1}{3} \frac{z}{1 - z} \hat{f}_1(z),$$

(13)

where $m_q$ is the constituent quark mass, and $M_n$ the nucleon mass. The above relation is only true at the scale of $\Lambda_\chi$, and at higher scale their relation might breakdown because the evolution of these two functions is different. However, as a rough estimate, we will adopt such approximations. The chiral quark model prediction for the usual unpolarized fragmentation function $\hat{f}_1(z)$ is consistent with the experimental data after considering the evolution effects [29]. We should also note that the chiral quark model is not suitable for the calculation of the fragmentation function at $z \rightarrow 1$ region, where the invariant mass of the fragmenting quark exceeds the cutoff of the model, $\Lambda_\chi$.

The $z$ dependence of the asymmetry $A_y$ only comes from the ratio of the two fragmentation functions $\hat{e}(z)$ and $\hat{f}_1(z)$ in Eq. (12), and the simple relation Eq. (13) can be used to predict the $z$ dependence of $A_y$. In Fig. 2, we show the normalized asymmetry prediction from this term compared with CLAS measurements. The most striking observation is that this simple relation Eq. (13) agrees with the experiment very well. The normalization of $A_y$ also depends on the TMD parton distribution $h_1^\perp$. This distribution involves more complicated dynamics [24, 25, 26], and hence is less reliable compared to the fragmentation function $\hat{e}(z)$ in Eq. (13). Nevertheless, from what we have now for $h_1^\perp$, we can make an order-of-magnitude estimate and compare with experiment. For example, a bag model calculation shows that the ratio of $h_1^{(1/2)}(x)/f_1(x)$ at the kinematic region of the CLAS measurement $0.15 < x < 0.4$ is about 0.04 for $u$ quark [26]. After taking into account other kinematic factors in Eq. (12), the asymmetry $A_y$ is predicted to be about 0.05, in rough agreement with the CLAS result of 0.038 [18], although the bag model prediction is very crude and the sign is inconsistent.
Extending the above estimate to the HERMES kinematics, the beam SSA is at least a factor of 5 less than what CLAS has found, consistent with the HERMES measurement. This is just the consequence of the beam SSA being 1/Q effect. So, there is no contradiction between these two experiments. This also agrees with the observation of [19].

We note that another interpretation of the beam SSA has been made in [31], where the photon “Sivers” effect was considered. In Ref. [30], an O(α_s^2) QCD effect to the beam SSA have also been investigated. We did not include these effects in our formalism.

In conclusion, we have calculated the beam single spin asymmetry in semi-inclusive hadron production in deep inelastic scattering. Up to 1/Q, there are two contributions: the distribution h_1^+(x, k_{⊥}) convoluted with fragmentation ˆe(z), and e(x) convoluted with H_1^+(z, k_{⊥}'). A simple chiral quark model prediction of ˆe(z) ≈ z/3/(1 − z) ˆf_1(z) agrees well with the experimental data on the z dependence of the asymmetry. Further experimental data can provide more information on the extraction of the leading order TMD parton distribution h_1^+.

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