Single Transverse Spin Asymmetries

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Using generalized factorization theorems, we calculate the leading single transverse spin asymmetry for high-transverse-momentum direct-photon production in pp collisions, in terms of partonic matrix elements. The leading contribution comes from a "twist-3" parton distribution, involving the correlation between quark fields and the gluonic field strength. With simple assumptions for this matrix element, the asymmetry increases with x_F, naturally giving effects of 10% or more at large x_F.

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Nucleon cross sections show important spin dependence. The measurement of longitudinal polarization dependence in deeply inelastic scattering (DIS) [1] has given rise to renewed interest in both longitudinal and transverse spin [2]. In fact, the dependence of hadronic cross sections on transverse polarizations has a long history [3,4], and important effects are seen in "single-spin" experiments, in which only one particle in the initial state is polarized, or in which only the polarization of a single final-state particle is observed. In this Letter we describe how recently developed methods in perturbative factorization at higher twist [5,6] may be used to describe single transverse spin asymmetries, and we discuss as an example high-x_F direct-photon cross sections at large transverse momentum l_T,

$$N(p,s_T)+N'(p')\to\gamma(l)+X. \tag{1}$$

$N(p,s_T)$ and $N'(p')$ represent a transversely polarized nucleon of momentum p and spin s_T, and an unpolarized nucleon of momentum p'. In the following, we shall take $p^\mu\sim\bar{n}^\mu\sqrt{s/2}$ with $\bar{n}^\mu=\delta_{\mu+}$ and $p'^\mu\sim n^\mu\sqrt{s/2}$ with $n^\mu=\delta_{\mu-}$. The same methods can be used to evaluate single transverse spin asymmetries in pion or jet production in the collision of these hadrons, and in Drell-Yan cross sections, where the azimuthal angular distribution of a lepton is observed. The direct-photon process is attractive for study because of its relative simplicity at lowest order in the strong coupling, and because our calculations provide predictions that we believe are experimentally accessible.

Christ and Lee [7] pointed out many years ago that time-reversal invariance forbids single transverse spin asymmetry in DIS to lowest order in $\alpha_{\rm EM}$. In hadron-hadron scattering, however, the presence of initial-state interactions allows single transverse spin asymmetries for final-state photons as well as hadrons.

Our results show a potentially healthy asymmetry in the direct-photon process at moderate transverse momenta, whose observation would provide new information on nucleon structure through the expectation value of a nonlocal operator that combines the quark fields with a single gluonic field strength,

$$T(x,s_T)=\int\frac{dy_1^-}{4\pi}e^{ixp^+y_1^-}\langle p,s_T|\bar{\psi}(0)\gamma^+\int dy_2^-\,\epsilon_{\sigma\rho\alpha\beta}s_T^\sigma n^\alpha\bar{n}^\beta F^{\rho+}(y_2^-)\psi(y_1^-)|p,s_T\rangle. \tag{2}$$

Here and below we take $s_T^2=-1$. This matrix element describes the color singlet coupling of the quark density at fractional momentum x to an averaged color field strength and the spin. [We have suppressed ordered exponentials of $n\cdot A$, which connect the fields along the light cone, and make $T(x,s_T)$ gauge invariant.] The operators in $T(x,s_T)$ may be compared to the spin-orbit coupling of an electron in a central field, although here, of course, the spin is not necessarily associated directly with the quark. We note that Ryskin [8] has described a model in which single-spin asymmetries are related to color spin-field strength couplings.

That perturbative QCD can be used to study the effects of transverse spin was realized some time ago by Kane, Pumplin, and Repko [9], and Efremov and Teryaev [10]. The relatively large size of experimentally observed effects remained, however, a difficulty [4,10]. By employing factorization methods, we show that the matrix element of Eq. (2) contributes at first order in the strong coupling, $g(p_T)$. This is to be compared with order $\alpha_s(p_T)$ for the operator [10],

$$T_m(x,s_T)=(p^2)^{1/2}\int\frac{dy^-}{4\pi}e^{ixp^+y^-}\langle p,s_T|\bar{\psi}(0)\gamma^+\gamma_5\gamma s_T\psi(y_1^-)|p,s_T\rangle, \tag{3}$$

with which $T(x,s_T)$ mixes. In fact, a possible role for matrix elements involving the field strength was suggested by Efremov and Teryaev in Ref. [10]. Subsequently [11], these authors discussed single-spin asymmetries in inclusive Compton scattering, $\gamma+N\to\gamma+X$, from a point of view very much like ours, using higher-twist techniques developed for DIS [12] in light-cone gauge. They identified the role of operators with a single covariant derivative $D_\rho=i\partial_\rho-gA_\rho$, such as

$$T_q(x,s_T)=p^+\int dy_1^-dy_2^-e^{ixp^+(y_1^--y_2^-)}\langle p,s_T|\bar{\psi}(0)\gamma^+\epsilon_{\sigma\rho\alpha\beta}s_T^\sigma n^\alpha\bar{n}^\beta D^\rho(y_2^-)\psi(y_1^-)|p,s_T\rangle. \tag{4}$$

There is a close relation between $T(x,s_T)$ [Eq. (2)] and $T_q(x,s_T)$ [Eq. (4)]. Both are off-diagonal contributions to a forward matrix element. $T(x,s_T)$ measures the overlap between states that differ by a single gluon with zero momentum fraction. In $T_q(x,s_T)$, on the other hand, the overlap is between states that differ by the transfer of the entire longitudinal momentum of a quark to a gluon, with the emission of a single soft quark. In perturbation theory, however, soft gluons are emitted much more readily than soft quarks. It is therefore natural to expect the new matrix element, $T(x,s_T)$, to be larger than $T_q(x,s_T)$, and we shall concentrate on the former in what follows.

It is important to emphasize that in any cross section described by a factorization formula, the overall normalization remains unknown until the relevant parton distribution (leading or higher twist) is measured. Given our comments above, the observation of an asymmetry in our direct-photon cross section would be essentially a measurement of the matrix element of Eq. (2). In future work, we hope to present an analysis of asymmetries in production cross sections for hadrons as well. Depending on the factorization formulas found in that case, it may be possible to relate the normalizations of the various cross sections. As we shall see, a modest size for the matrix element in Eq. (2) is sufficient to produce a sizable asymmetry in the transverse-momentum distribution of direct photons.

Physical observables that depend on the transverse polarization of a single hadron are typically power corrections, in comparison with the leading contributions encountered in spin-averaged or longitudinally polarized cross sections. (Leading contributions in the scattering of *two* transversely polarized hadrons have been discussed in Ref. [13].) For example, the transverse spin asymmetry in deeply inelastic scattering provides a measurement of the combination of structure functions $g_1(x)+g_2(x)$, which is related to matrix elements of twist-3 operators. The factorization of power corrections, which gives this result, follows directly from the operator-product expansion in DIS, but its extension to hadron-hadron scattering requires additional arguments. In Ref. [14], we extended the factorization program to $O(1/Q^2)$ corrections for unpolarized hadron-hadron cross sections, and in [6,15] to $O(1/Q)$ corrections in polarized cross sections, just far enough to include the corresponding dependence on transverse spin. The general factorization formula for cross sections of polarized hadronic scatterings is of the form

$$\sigma(Q)=H^0\otimes f_2\otimes f_2+\left(\frac{1}{Q}\right)H^1\otimes f_2\otimes f_3+O\left(\frac{1}{Q^2}\right). \quad (5)$$

H^0 and H^1 are perturbatively calculable coefficient functions, and the f_n's, with $n=2,3$, are twist-n matrix elements.

As mentioned above, for direct-photon production we shall consider for simplicity the range of large $x_F=(t-u)/s$, with $s=(p+p')^2$, $t=(p-l)^2$, and $u=(p'-l)^2$. In this kinematic region, we expect "Compton" processes to dominate, initiated in the hard-scattering functions of Eq. (5) by a valence quark from the polarized nucleon $N(p,s_T)$ in Eq. (1), and a gluon from the unpolarized nucleon $N'(p')$ [16]. In this case, the relevant twist-3 matrix element of Eq. (2) arises from gluonic corrections in the polarized nucleon.

To define the asymmetry in the cross section, we denote by $\sigma(s_T,l)$ the direct-photon-production cross sections for the scattering of Eq. (1). The transverse asymmetry in the cross section for the polarized process is then

$$\Delta\sigma_\perp \equiv \tfrac{1}{2}[\sigma(s_T,l)-\sigma(-s_T,l)]. \quad (6)$$

A typical Feynman diagram that gives a leading contribution in the strong coupling to $\Delta\sigma_\perp$ is shown in Fig. 1. We work in Feynman gauge, where diagrams with more gluons also contribute at the same power of Q, but only through unphysical polarizations. These higher-order corrections produce gauge-invariant matrix elements, but do not affect the short-distance hard parts [6]. It is our purpose to exhibit the leading short-distance hard parts in the following, and we therefore need only the diagrams with one gluon from the polarized nucleon. The single gluon here is essential for obtaining a nonvanishing asymmetry $\Delta\sigma_\perp$ at tree level [11].

The factorization procedure described in Ref. [6] includes: (i) a "collinear expansion" in terms of the parton momenta entering the hard scattering and of the polarizations of gluon fields that link the hard scattering to the incoming hadron, (ii) the separation of spinor and Lorentz indices between the hard part and the resulting long-distance matrix elements, and (iii) the calculation of the partonic hard part. In Feynman gauge it is relatively easy to identify contributions to matrix elements involving the field strength $F^{\rho+}$ from the expansion of the hard scattering in the momentum k^ρ of the gauge-field component $n\cdot A$. $T(x,s_T)$ [Eq. (2)] comes from this expansion, after a contour integral in the component $n\cdot k$, which flows from the polarized hadron through the hard scattering.

By applying the above procedure, and using invariance under parity and time reversal, we find that the leading nonvanishing contribution to the asymmetry defined in

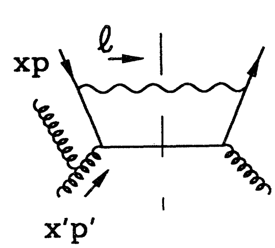

FIG. 1. Sample diagram contributing to single transverse spin asymmetry in direct-photon production.

Eq. (6) is

$$E_l \frac{d\Delta\sigma_\perp}{d^3 l} = \alpha_{\rm EM}\alpha_s \int \frac{dx'}{x'} G(x') \int \frac{dx}{x} \left(\frac{1}{s}\right) \delta(\hat{s}+\hat{t}+\hat{u}) \epsilon_{\rho\sigma\alpha\beta} s_T^\rho l^\sigma n^\alpha \bar{n}^\beta \times \left\{ T(x,s_T) H_1^1(x,x',l) + \left[T(x,s_T) - x\frac{\partial}{\partial x} T(x,s_T) \right] H_2^1(x,x',l) \right\}. \quad (7)$$

Here $G(x')$ is the usual gluon distribution, and T is the twist-3 matrix element given in Eq. (2), where a sum over quark flavor $\sum_f e_f^2$ is implicitly included, with e_f the fractional electric charge of quarks of flavor f. The "hard-scattering" functions H_i^1 are given by

$$H_1^1(x,x',l) = \frac{1}{2N} \frac{N^2}{N^2-1} 2g \left[\frac{1}{-\hat{t}} \right], \quad (8a)$$

$$H_2^1(x,x',l) = \frac{1}{2N} \frac{N^2}{N^2-1} 2g \left[\left(\frac{\hat{s}}{-\hat{t}} \right) + \left(\frac{-\hat{t}}{\hat{s}} \right) \right] \left(\frac{1}{-\hat{u}} \right), \quad (8b)$$

where g is the strong-coupling constant, and the carets refer to partonic invariants, $\hat{s}=(xp+x'p')^2=xx's$, $\hat{t}=(xp-l)^2=xt$, and $\hat{u}=(x'p'-l)^2=x'u$. We have exhibited explicitly the dependence on the number of colors, N. In this form, we see that the asymmetry is proportional to the transverse momentum of the observed photon, relative to the beam direction. Note that the H_i^1 are proportional to g, rather than α_s. Because it is part of H^1, the scale of the strong coupling is set by l_T^2.

The significance of the asymmetry in Eq. (7) is measured by its ratio to the unpolarized cross section, given at leading power by

$$E_l \frac{d\sigma}{d^3 l} = \alpha_{\rm EM}\alpha_s \int \frac{dx'}{x'} G(x') \int \frac{dx}{x} \left(\frac{1}{s}\right) \delta(\hat{s}+\hat{t}+\hat{u}) \left(\frac{F_2(x)}{x} \right) \hat{\sigma}(x,x',l) + p \leftrightarrow p', \quad (9)$$

where again, we keep only the Compton subprocess. In Eq. (9), $F_2(x) = \sum_f e_f^2 x q_f(x)$ with quark distributions $q_f(x)$, and the hard-scattering function $\hat{\sigma}$ at lowest order is

$$\hat{\sigma}(x,x',l) = 2\left(\frac{1}{2N}\right) \left[\frac{\hat{s}}{-\hat{t}} + \frac{-\hat{t}}{\hat{s}} \right]. \quad (10)$$

From the ratio of Eqs. (7) and (9), we may compute the fractional asymmetry as a function of l_T and x_F,

$$A(s_T,x_F,l_T) = E_l \frac{d\Delta\sigma_\perp}{d^3 l} \Big/ E_l \frac{d\sigma}{d^3 l}. \quad (11)$$

To estimate the asymmetry, we model the "twist-3" matrix element, $T(x,s_T)$, as a mass scale times a dimensionless function of x. We choose the mass scale to be 0.2 GeV ($\approx \Lambda_{\rm QCD}$), while for the function we compare two guesses. The first is

$$T(x,s_T) \approx 0.2 F_2(x)/x \ \ ({\rm GeV}), \quad (12)$$

with $F_2(x)$ the usual deeply inelastic scattering structure function, calculated according to parton-model quark distributions, which we will give in a moment. In the second guess, we simply multiply Eq. (12) by x, to produce a function that is suppressed at small x, in a manner reminiscent of the spin distribution $g_1(x)$,

$$T(x,s_T) \approx 0.2 F_2(x) \ \ ({\rm GeV}). \quad (13)$$

For simplicity, we parametrize conventionally normalized parton distributions, without scaling violation, as

$$xu_v(x) = \frac{2}{B(0.5,4)} x^{0.5}(1-x)^3, \quad (14a)$$

$$xd_v(x) = \frac{1}{B(0.5,4.5)} x^{0.5}(1-x)^{3.5}, \quad (14b)$$

$$xS(x) = 8\left[\frac{1}{2} - 2\frac{B(1.5,4)}{B(0.5,4)} - \frac{B(1.5,4.5)}{B(0.5,4.5)} \right] (1-x)^7, \quad (14c)$$

$$xG(x) = 3(1-x)^5, \quad (14d)$$

where B is the beta function. The resulting ratios A [Eq. (11)], for s_T perpendicular to l_T, are plotted in Fig. 2 as a function of x_F at $\sqrt{s}=30$ GeV and $l_T=4$ GeV. The lower (solid) line results from using Eq. (13) for the distribution $T(x,s_T)$, and the upper (dashed) line from using Eq. (12). In both cases, A rises to over 20% as x_F approaches 0.8. The origin of this effect is easy to understand. In the large-x_F region, the typical value of x in Eq. (7) is large, and the term with $x(\partial/\partial x)T(x,s_T)$ dominates and increases with x_F, relative to the unpolarized cross section. This is because the derivative of any function which behaves as $(1-x)^a$, $a>0$, vanishes less rapidly than the function itself as $x \to 1$. Such derivative terms have been found before at twist-4 [5,12], but in the asymmetry they occur at twist-3 only times the matrix element $T(x,s_T)$, Eq. (2). This further justifies our con-

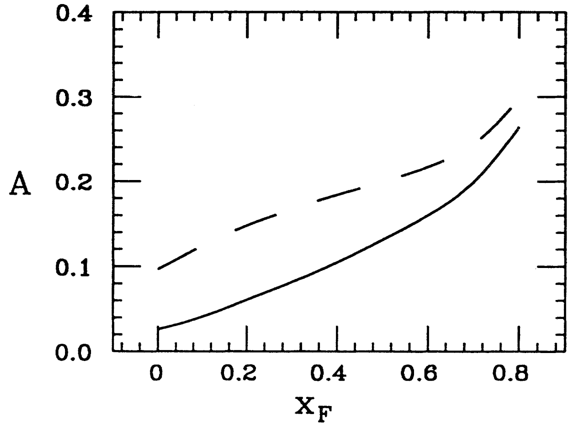

FIG. 2. Single transverse spin asymmetry A in pp direct-photon production as a function of x_F, at $\sqrt{s}=30$ GeV, $l_T=4$ GeV. The solid line is computed using the matrix element in Eq. (13), the dashed line using the one in Eq. (12).

centration on $T(x,s_T)$ in the large-x_F region.

We emphasize that the asymmetry shown in Fig. 2 is dependent on our model of the twist-3 matrix element. For instance, although the asymmetry is positive for the specific choices for $T(x,s_T)$ in Eqs. (12) and (13), its actual sign is not fixed by our considerations. In addition, we anticipate substantial higher-order corrections, as have already been observed at leading power [17]. Such corrections typically factor from the underlying hard scattering [18], and we expect them to cancel, at least approximately, in the ratio A, Eq. (11). This question, of course, bears more study. Nevertheless, it is clear that a modest size for the matrix element is sufficient to produce a very significant asymmetry in the large-x_F region for photons, comparable to those found for pion production [4]. As mentioned above, an experimental measurement of nonzero asymmetry over a range in x_F could be used to determine $T(x,s_T)$. This would supply qualitatively new information about partonic correlations in the nucleon.

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