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Double spin asymmetry A_{LT} in direct photon production

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A R T I C L E   I N F O

A B S T R A C T

We study the longitudinal-transverse double spin asymmetry A_{LT} for direct photon production in nucleon–nucleon scattering by using the collinear twist-3 approach. This asymmetry, which, for instance, could be measured at RHIC, contains a complete set of collinear twist-3 correlation functions in a transversely polarized nucleon.

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1. Introduction

High energy experiments with polarized beams and targets have opened a new window for revealing QCD dynamics and hadron structure. Based on QCD factorization theorems, polarization-dependent cross sections generally can be factorized into the convolution of perturbatively calculable hard parts and universal nonperturbative (soft) parts, which are expressed through various spin-dependent parton correlation functions. Among these functions, higher-twist spin-dependent correlation functions are poorly known in comparison to the three leading-twist ones: the spin-averaged parton distribution f_1, the helicity distribution g_1, and the quark transversity h_1 [1–3]. Along with leading-twist distributions, the higher-twist correlation functions provide us with important information on the structure of hadrons, even though they do not have a probability interpretation. The best way of extracting them is to investigate spin observables which have no leading-twist contribution. A classic example is the twist-3 double spin asymmetry A_{LT} (longitudinally polarized lepton beam, transversely polarized target) in inclusive deep-inelastic lepton–nucleon scattering (DIS), which allows one to study the parton correlator g_T. In the case of nucleon–nucleon scattering, the double spin asymmetry A_{LT} in the Drell–Yan process, involving two polarized incident hadrons, has been extensively investigated for the purpose of studying higher-twist correlators [3–6]. More recently, A_{LT} for inclusive lepton production from the decay of W-bosons in proton–proton scattering and for jet production in lepton–proton scattering has been derived [7,8]. In this Letter we focus on the double spin asymmetry A_{LT} for direct photon production in nucleon–

nucleon scattering, which also allows one to study twist-3 spin-dependent parton distributions as leading effects. In contrast to the aforementioned processes, direct photon production contains a complete set of collinear twist-3 correlation functions in a transversely polarized nucleon.

The crucial tool required for the extraction of twist-3 correlation functions is collinear higher-twist factorization, which has been established in studying transverse single spin asymmetries (SSAs) [9–14]. The exploration of SSAs in hadronic reactions has a long history, starting from the mid 1970's. In particular, the large size of the observed SSAs for single inclusive hadron production [15–20] came as a big surprise and, *a priori*, posed a challenge for QCD, because the collinear parton model predicts the asymmetries are proportional to $\alpha_s m_q / P_{h\perp}$ [21,22], where m_q is the quark mass and $P_{h\perp}$ is the transverse momentum of the final state hadron. However, significant SSAs in hadronic collisions may be generated by going beyond the naive parton model and including collinear twist-3 parton correlators, as was first pointed out in [9], and later on studied in more detail [10–14,23]. (We also note that alternative mechanisms underlying the large SSAs have been proposed [24–26].) In the process $p^\uparrow p \to hX$, the collinear twist-3 formulation has some relation to a description in terms of transverse momentum dependent parton correlators (TMDs) [27–30], provided that initial/final state interactions in the TMD approach are taken into account [31]. In particular, for semi-inclusive DIS and related processes, TMDs are of crucial importance (see, for instance [32–41] and references therein). For these reactions, also intriguing nonzero spin/azimuthal asymmetries were observed [42–45]. In the recent past, important progress was made in understanding various spin observables in terms of the collinear twist-3 approach and/or the TMD approach. To mention just one example, it was found that for certain structure functions in semi-inclusive DIS the two formalisms provide the same

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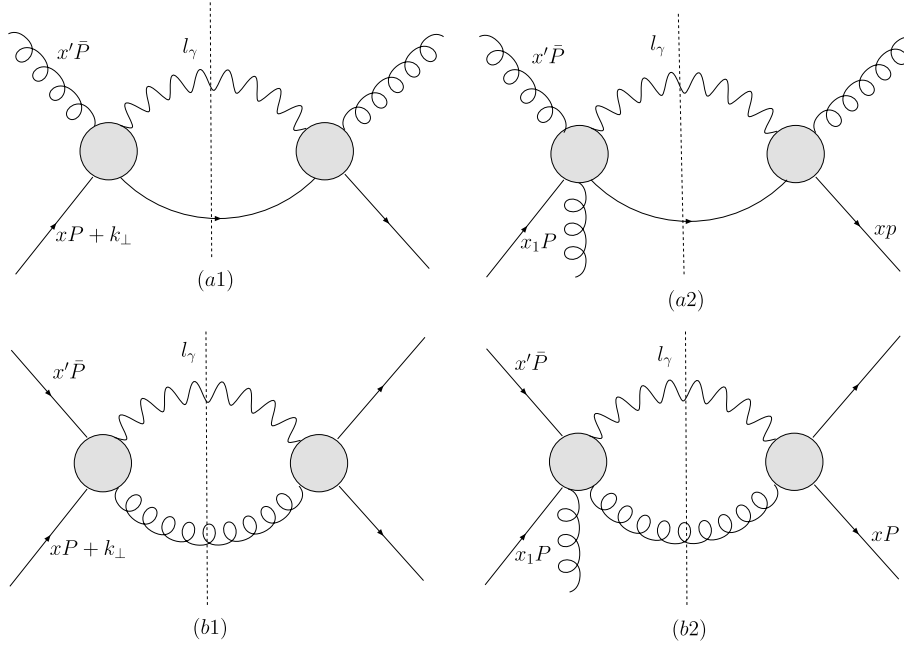*E-mail address:* zzzhoujian@gmail.com (J. Zhou).

Fig. 1. Generic Feynman diagrams for the partonic channels $qg \to q\gamma$ and $q\bar{q} \to g\gamma$. Diagrams (a1) and (b1) contribute to the hard part associated with the soft part $\langle \bar{\psi} \partial_\perp \psi \rangle$; diagrams (a2) and (b2) contribute to the hard part associated with the soft part $\langle \bar{\psi} A_\perp \psi \rangle$.

result at intermediate transverse hadron momenta [46–49], which can be viewed as a nontrivial consistency check.

By using the collinear higher-twist approach, we are able to compute the pertinent spin-dependent cross section (numerator of the asymmetry A_{LT}) for direct photon production in nucleon–nucleon scattering and express it in terms of twist-3 parton distributions. In Section 2, we briefly review the factorization formalism as it applies to the first non-leading term of an expansion in powers of $1/l_{\gamma\perp}$, with $l_{\gamma\perp}$ denoting the transverse momentum of the observed photon. We also list the twist-3 correlation functions that appear in the factorization formula and provide the complete twist-3 spin-dependent cross section. We conclude the Letter in Section 3.

2. Calculation of double spin-dependent cross section

For definiteness we consider the process

$$N(P, S_\perp) + N(\bar{P}, \Lambda) \to \gamma(l_\gamma) + X, \qquad (1)$$

where we indicated the 4-momenta and the polarization states of the particles. Furthermore, we define the Mandelstam variables through $S = (P + \bar{P})^2$, $T = (P - l_\gamma)^2$, and $U = (\bar{P} - l_\gamma)^2$. On the partonic level one has $\hat{s} = xx'S$, $\hat{t} = xT$, and $\hat{u} = x'U$, with x and x' representing the longitudinal momentum fractions of active partons coming from the transversely polarized nucleon and the longitudinally polarized nucleon, respectively.

The generic form of the (spin-dependent) cross section reads

$$d\sigma(l_{\gamma\perp}, S_\perp, \Lambda)$$
$$= H^{(0)} \otimes f_2 \otimes f_2 + \frac{1}{l_{\gamma\perp}} H^{(1)} \otimes f_2 \otimes f_3 + \mathcal{O}\left(\frac{1}{l_{\gamma\perp}^2}\right), \qquad (2)$$

where f_2 (f_3) indicates a twist-2 (twist-3) parton distribution, while $H^{(0)}$ and $H^{(1)}$ are the corresponding perturbatively calculable coefficient functions. The direct products denote convolutions in the fractional parton momenta. The second term in Eq. (2) gives rise to a leading non-vanishing contribution to the double spin asymmetry A_{LT}. The lowest-order contributions to $H^{(1)}$ are found

from the Born diagrams for the partonic process (see Fig. 1(a1), (b1)), plus $\mathcal{O}(g)$ corrections in which an extra gluon is exchanged between the remnants of the transversely polarized nucleon and the hard partonic scattering process (see Fig. 1(a2), (b2)). In order to isolate the twist-3 contributions to the Born diagrams at this order, we employ a collinear expansion in the parton momenta.

To proceed further, let us briefly review the relevant twist-3 correlation functions involved in our calculation. For the A_{LT} asymmetry we consider twist-3 correlators in the transversely polarized nucleon, together with ordinary twist-2 helicity distributions g_1 in the longitudinally polarized nucleon. The so-called D-type (twist-3) quark–gluon–quark functions (here denoted by G_D and $\tilde{G}_D$) have been introduced a long time ago and defined through [3,50]

$$\int \frac{dy^-}{2\pi} \frac{dy_1^-}{2\pi} e^{-ixP^+ y^-} e^{i(x-x_1)P^+ y_1^-}$$
$$\times \langle P, S_\perp | \bar{\psi}_\beta(y^-) iD_\perp^\mu(y_1^-) \psi_\alpha(0) | P, S_\perp \rangle$$
$$= \frac{M}{2P^+} \left[G_D(x, x_1) i\varepsilon_\perp^{\mu\nu} S_{\perp\nu} \slashed{n} + \tilde{G}_D(x, x_1) S_\perp^\mu \gamma_5 \slashed{n} \right]_{\alpha\beta}, \qquad (3)$$

where the gauge links between the field operators have been suppressed. The hadron momentum P is proportional to the light cone vector $n = (1^+, 0^-, \vec{0}_\perp)$, whose conjugate light-cone vector is $\bar{n} = (0^+, 1^-, \vec{0}_\perp)$. The nucleon mass is denoted by M. The gluon field enters through the covariant derivative $D_\perp^\mu = \partial_\perp^\mu - igA_\perp^\mu$. The variables x, x_1 are the momentum fractions of the hadron carried by the quarks, implying that the gluon momentum fraction is given by $x_g = x - x_1$. We also recall the relation between the D-type functions in Eq. (3) and the twist-3 quark-quark correlator g_T [3],

$$xg_T(x) = \int_0^1 dx_1 \left[G_D(x, x_1) + \tilde{G}_D(x, x_1) \right]. \qquad (4)$$

In the case of A_{LT} in inclusive DIS, for instance, G_D and $\tilde{G}_D$ appear with the same hard scattering coefficient such that the final result is proportional to g_T.

Making use of the field strength tensor offers an alternative way of defining gauge invariant quark–gluon–quark correlators [9,10],

$$\int \frac{dy^-}{2\pi} \frac{dy_1^-}{2\pi} e^{-ixP^+y^-} e^{i(x-x_1)P^+y_1^-}$$
$$\times \langle P, S_\perp | \bar{\psi}_\beta(y^-) g F_\perp^{+\mu}(y_1^-) \psi_\alpha(0) | P, S_\perp \rangle$$
$$= \frac{M}{2} \big[T_F(x,x_1) \varepsilon_\perp^{\nu\mu} S_{\perp\nu} \slashed{n} + \tilde{T}_F(x,x_1) i S_\perp^\mu \gamma_5 \slashed{n} \big]_{\alpha\beta}, \tag{5}$$

where T_F and $\tilde{T}_F$ are the so-called F-type functions. Note that our definition of these functions differs by a factor $2\pi M$ from the conventions used in Ref. [46,48]. It was found that the F-type functions directly enter the QCD-description of transverse SSAs in various processes [10–13].

The D-type functions and the F-type functions are not independent as they can be related to each other by means of the equation of motion [12,51],

$$G_D(x,x_1) = P \frac{1}{x-x_1} T_F(x,x_1), \tag{6}$$

$$\tilde{G}_D(x,x_1) = P \frac{1}{x-x_1} \tilde{T}_F(x,x_1) + \delta(x-x_1) \tilde{g}(x), \tag{7}$$

where P indicates the principal value prescription, and the function $\tilde{g}$ is given by [12,51]

$$\int \frac{dy^-}{2\pi} e^{-ixP^+y^-} \langle P, S_\perp | \bar{\psi}_\beta(y^-)$$
$$\times \left(i D_\perp^\mu + g \int_0^\infty d\zeta^- F_\perp^{+\mu}(\zeta^-) \right) \psi_\alpha(0) | P, S_\perp \rangle$$
$$= \frac{M}{2} \big[\tilde{g}(x) S_\perp^\mu \gamma_5 \slashed{n} \big]_{\alpha\beta}. \tag{8}$$

Because of the relations in Eqs. (6), (7) we may either use the D-type or the F-type functions for our calculation. However, it is mandatory to also include $\tilde{g}$ in order to completely describe the spin-dependent cross section. (See also [4,52–55] and references therein for a discussion about independent twist-3 correlation functions.) It is worthwhile to point out that the function $\tilde{g}$ is related to the twist-2 TMD g_{1T} describing the distribution of longitudinally polarized quarks in a transversely polarized nucleon [14,56],

$$\tilde{g}(x) = \int d^2\vec{k}_\perp \frac{\vec{k}_\perp^2}{2M^2} g_{1T}(x, \vec{k}_\perp^2), \tag{9}$$

where we used the definition of g_{1T} as given in [32,38].

As mentioned above, the twist expansion is the key step in our calculations. The relevant technical ingredients required for such a twist-3 analysis have been well developed in the last few decades [9–13,49,50]. In the twist expansion, a set of non-perturbative matrix elements of the hadron state is analyzed according to the power counting of the associated contributions. At the twist-3 level, the following matrix elements can contribute [50]:

$$\langle \bar{\psi} \partial_\perp \psi \rangle, \qquad \langle \bar{\psi} A_\perp \psi \rangle, \qquad \langle \bar{\psi} \partial_\perp A^+ \psi \rangle. \tag{10}$$

We found it most convenient to work in the light-cone gauge ($A^+ = 0$) in which only the first two matrix elements survive. The matrix element $\langle \bar{\psi} \partial_\perp \psi \rangle$ can be transformed into the gauge invariant matrix element $\tilde{g}$. Moreover, by partial integration, the matrix element $\langle \bar{\psi} A_\perp \psi \rangle$ can be expressed as $\langle \bar{\psi} \partial^+ A_\perp \psi \rangle$ and further be related to the gauge invariant matrix elements T_F and $\tilde{T}_F$. These functions may then be rewritten in terms of the D-type functions by means of Eqs. (6), (7).

Our goal is to perturbatively calculate the hard scattering coefficients associated with these two soft parts. The corresponding

partonic scattering processes are illustrated in Fig. 1. It turns out that the hard parts are not always real. Imaginary contributions occur whenever an internal parton line in the hard scattering goes on-shell. When this happens, we use the distribution identity

$$\frac{1}{x \pm i\epsilon} = P \frac{1}{x} \mp i\pi \delta(x). \tag{11}$$

While for the related calculations of transverse SSAs only the δ-function contribution matters, in this calculation for A_{LT} we are left with the principal value part. To be more precise, the imaginary parts always cancel between the different cut diagrams. In order to combine corresponding cut diagrams we use the symmetry properties

$$T_F(x,x_1) = T_F(x_1,x), \qquad \tilde{T}_F(x,x_1) = -\tilde{T}_F(x_1,x),$$
$$G_D(x,x_1) = -G_D(x_1,x), \qquad \tilde{G}_D(x,x_1) = \tilde{G}_D(x_1,x) \tag{12}$$

of the F-type functions and the D-type functions.

Making use of these ingredients, the calculation is rather straightforward, and we obtain the following result for the spin-dependent cross section:

$$l_\gamma^0 \frac{d\sigma(S_\perp, \Lambda)}{d^3 l_\gamma}$$
$$= -2M \frac{\alpha_s \alpha_{em}}{S} \Lambda \vec{S}_\perp \cdot \vec{l}_{\gamma\perp} \sum_{i=qg,q\bar{q}} \sum_{a,b} e_a^2 \int_{x'_{min}}^1 \frac{dx'}{x'} \frac{1}{x'S+T} \frac{1}{x\hat{u}} g_1^b(x')$$
$$\times \left\{ \left[\tilde{g}^a(x) - x \frac{d}{dx} \tilde{g}^a(x) \right] H_i^{\tilde{g}} + \int_0^1 dx_1 \big[G_D^a(x,x_1) H_i^{G_D} \right.$$
$$\left. + \tilde{G}_D^a(x,x_1) H_i^{\tilde{G}_D} \big] \right\}, \tag{13}$$

where $x = -x'U/(x'S+T)$ and $x'_{min} = -T/(S+U)$. The hard coefficient functions for the $qg \to q\gamma$ partonic channel are

$$H_{qg}^{\tilde{g}} = \frac{N_c}{N_c^2-1} \left[\frac{\hat{s}^2 - \hat{t}^2}{\hat{s}\hat{t}} \right],$$

$$H_{qg}^{G_D} = \frac{N_c}{N_c^2-1} \left[\frac{\hat{s}^2 - \hat{t}^2}{(1-\xi)\hat{s}\hat{t}} \right] + \frac{1}{N_c} \left[\frac{\hat{u}(\hat{s}^2 + 2\hat{t}^2)}{\hat{s}\hat{t}^2} + \frac{\hat{u}}{(1-\xi)\hat{t}} \right],$$

$$H_{qg}^{\tilde{G}_D} = \frac{N_c}{N_c^2-1} \left[\frac{(\xi-2)(\hat{s}^2-\hat{t}^2)}{\xi(1-\xi)\hat{s}\hat{t}} \right]$$
$$+ \frac{1}{N_c} \left[\frac{\hat{u}(\hat{s}^2+2\hat{t}^2)}{\hat{s}\hat{t}^2} + \frac{(\xi-2)\hat{u}}{\xi(1-\xi)\hat{t}} + \frac{2\hat{u}}{\xi\hat{s}} \right], \tag{14}$$

while the coefficient functions for the $q\bar{q} \to g\gamma$ channel read

$$H_{q\bar{q}}^{\tilde{g}} = \frac{1}{N_c^2} \left[\frac{\hat{t}^2 + \hat{u}^2}{\hat{t}\hat{u}} \right],$$

$$H_{q\bar{q}}^{G_D} = \frac{\hat{t}^2 + \hat{u}^2}{(1-\xi)\hat{t}\hat{u}} + \frac{2C_F}{N_c} \left[\frac{\hat{s}^2(\hat{t}-\hat{u})}{\hat{t}^2\hat{u}} - \frac{(\xi-2)(\hat{t}-\hat{u})}{(1-\xi)\hat{t}} \right],$$

$$H_{q\bar{q}}^{\tilde{G}_D} = \frac{(\xi-2)(\hat{t}^2+\hat{u}^2)}{\xi(1-\xi)\hat{t}\hat{u}} + \frac{2C_F}{N_c} \left[\frac{\hat{s}^2(\hat{t}-\hat{u})}{\hat{t}^2\hat{u}} - \frac{\xi(\hat{t}-\hat{u})}{(1-\xi)\hat{t}} \right], \tag{15}$$

with $\xi = x_g/x$. Note that the spin-dependent cross section relevant for the A_{LT} asymmetry is characterized by the correlation $\Lambda \vec{S}_\perp \cdot \vec{l}_{\gamma\perp}$ instead of the correlation $\vec{S}_\perp \times \vec{l}_{\gamma\perp}$ that arises in the case of transverse SSAs. We have expressed the cross section in terms of the D-type functions plus $\tilde{g}$. By doing so the contributions from the derivative term and the non-derivative term associated

with the correlator $\tilde{g}$ can be combined into the same compact form that was found for the SSAs for direct photon production and inclusive pion production [10,11,13]. We also point out that, unlike related previous studies of other processes [7,8], the coefficient functions for G_D and for $\tilde{G}_D$ are different. Therefore, in the present case it does not pay off to involve g_T in the final result. This also means that through direct photon production, in combination with other reactions, one should in principle be able to study a complete set of twist-3 correlators for a transversely polarized nucleon.

3. Summary

By using collinear twist-3 factorization we derived the longitudinal-transverse double spin asymmetry A_{LT} for direct photon production in nucleon–nucleon scattering. Our study can be considered as the counterpart of the calculation of the transverse SSA in the same process [10]. Measuring this observable might also open a new window to test the higher-twist approach, for instance by looking at the transverse momentum behavior of A_{LT}. Also, due to the derivative term of the correlator $\tilde{g}$, the asymmetry may be largest in the large x_F region — the same kinematics for which the largest SSAs in hadron production have been observed. Moreover, if the gluon helicity distribution (in the small x region) is very small (see, e.g., Ref. [57]), one can expect A_{LT} to be larger for $p\bar{p}$ scattering than for pp-scattering.

The result for A_{LT} depends on a complete set of collinear twist-3 parton correlators for a transversely polarized nucleon. It requires numerical studies to find out to what extent these correlators could, at least in principle, be separately explored by variation of the kinematics. Most likely one would need the combined information from different processes to address all three twist-3 correlation functions. Besides the twist-3 effect investigated in this Letter, a collinear twist-3 correlator for the longitudinally polarized hadron, coupling to the transversity distribution, may also contribute to A_{LT}. This part as well as numerical estimates are left for future study. One can further extend our work to the A_{LT} asymmetry in other processes such as single inclusive hadron production and jet production [58].

Acknowledgements

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