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Gluon TMDs at the LHC and RHIC

M. Schlegel
Institute for Theoretical Physics, Tuebingen University
Auf der Morgenstelle 14, D-72076 Tuebingen, Germany

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Summary. — It is argued that transverse-momentum–dependent (TMD) gluonic parton distributions can be relevant for physics probed in proton collisions at high energies. The production of two real photons in proton collisions is discussed within the framework of TMD factorization at small transverse pair momenta where information on gluon TMDs can be extracted. Spin-dependent gluon TMDs can be studied in polarized proton collisions at RHIC. On the other hand, the effect of the distribution of linearly polarized gluons on unpolarized proton collisions at the LHC may be utilized to determine the parity of the Higgs boson.

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1. – Introduction

The spin structure of the nucleon has attracted an enormous amount of interest in recent years. One of the theoretical tools to study the partonic spin degrees of freedom of hadrons in QCD is provided by the so-called transverse momentum dependent (TMD) factorization [1]. This kind of factorization is well suited to describe hard processes with small final state transverse momenta. Integral parts of TMD factorization are TMD parton distributions that depend on small "intrinsic" transverse parton momenta. This sensitivity on the partonic motion enables additional non-perturbative correlations between the hadron/nucleon spin and the parton spin that can be accessed experimentally (cf. the recent experimental review [2]). In this context two processes can reveal information on quark TMDs, *i.e.* semi-inclusive deep-inelastic lepton scattering off nucleons (SIDIS) and lepton pair production in hadron collisions (Drell-Yan process, DY). Of particular interest are transverse single-spin asymmetries that are caused by a naive time-reversal odd (T-odd) TMD parton distribution which describes unpolarized quarks in a transversely polarized nucleon, the so-called Sivers function [3].

Gluonic TMD parton distributions (a first parameterization was given in [4]) have been studied to a lesser extend partly because they don't appear in the TMD factorization formulas for SIDIS and DY. Hence, these objects remain unknown. They were found to

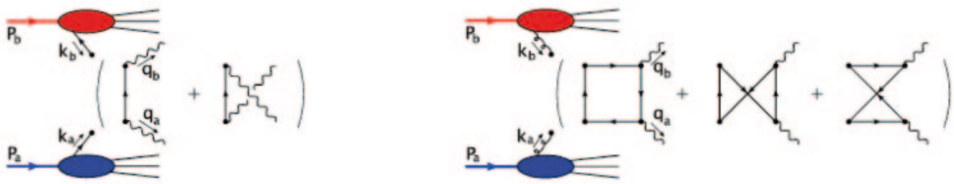

Fig. 1. – Left: $q\bar{q} \to \gamma\gamma$-channel. Right: $gg \to \gamma\gamma$-channel at $\mathcal{O}(\alpha_s^2)$.

appear in the TMD framework for proton collisions with hadronic final states [5, 6]. However TMD factorization has been shown to be violated for these processes [7]. It was also suggested to study gluon TMDs in heavy-quark production in lepton nucleon collisions [8]. TMD factorization is valid in this process and measurements can be realized at a future electron-ion collider (EIC).

Since it is a necessity to have a colorless final state in proton collisions in order to establish TMD factorization the production of a photon pair has been investigated in refs. [9, 10]. Here, gluon distributions enter the description of the cross section by means of gluon fusion to a photon pair via a quark box on the parton level at $\mathcal{O}(\alpha_s^2)$. It was demonstrated in [9] that gluons may give sizeable contributions already at RHIC energies $\sqrt{\mathcal{S}} = 500\,\text{GeV}$. In [10] it was found that a specific gluon TMD, the distribution of linearly polarized gluons $h_1^{\perp g}$, can be useful for LHC physics, to be specific, for the determination of the parity of a Higgs boson. It is the aim of these proceedings to discuss and summarize the results of refs. [9, 10].

2. – Photon pair production in the TMD framework

The diphoton process $h(P_a) + h(P_b) \to \gamma(q_a) + \gamma(q_b) + X$ is analyzed in the center-of-mass (c.m.) frame of the incoming hadrons where hadron a moves along the positive z-direction. The spatial orientation of the photons is fixed by two parameters θ, ϕ which are physical angles in the so-called Collins-Soper (CS) frame [1]. The Lorentz transformation between the c.m. frame and CS frame has been worked out in ref. [11]. The explicit form of the external momenta of the hadrons and photons is given in terms of the pair momentum $q = q_a + q_b$, its mass $Q^2 = q^2$, the variables $x_{a/b} = q^2/(2P_{a/b} \cdot q)$, and the c.m. energy $\mathcal{S} = (P_a + P_b)^2$. The partonic Mandelstam variables expressed in the c.m. frame read $s = 2k_a \cdot k_b = Q^2$, $t = -2k_a \cdot q_a = -Q^2 \sin^2 \frac{\theta}{2}$ and $u = -2k_b \cdot q_a = -Q^2 \cos^2 \frac{\theta}{2}$.

It is instructive to calculate the simplest subprocess within the TMD formalism first, *i.e.* photon pair production via quark-antiquark annihilation $q\bar{q} \to \gamma\gamma$ (cf. fig. 1). The calculation can be performed along the lines of ref. [11] where results were discussed for the DY cross section in the TMD framework at leading twist. Indeed, one can directly compare the results for the DY cross section with the $q\bar{q}$-channel of the photon pair production,

$$\text{(1)} \qquad \left(\frac{\mathrm{d}\sigma^{q\bar{q}\to\gamma\gamma}}{\mathrm{d}^4 q \mathrm{d}\Omega}\right)(q_T \ll Q) = \frac{2}{\sin^2\theta}\left(\frac{\mathrm{d}\sigma^{q\bar{q}\to l^+ l^-}}{\mathrm{d}^4 q \mathrm{d}\Omega}\right)(q_T \ll Q | Q_q \to Q_q^2).$$

The various (leading twist) spin observables in DY on the r.h.s. can be read off from ref. [11]. $Q_q \to Q_q^2$ indicates an obvious replacement of the quark charges in the DY expressions by the squared quark charges. It is important to note that both sides of (1) incorporate the *same* TMDs since initial state interactions are relevant only for both processes.

In a next step the gluonic channel shown in the right panel of fig. 1 is calculated. If the general partonic amplitude $gg \to \gamma\gamma$ is denoted by $\mathcal{M}^{ij}_{\lambda_1\lambda_2}$ (λ are the helicities of the photons, i and j transverse gluon indices) the formula for the TMD factorized hadronic cross section reads,

$$\begin{aligned}(2)\quad \left(\frac{\mathrm{d}\sigma^{gg\to\gamma\gamma}}{\mathrm{d}^4q\mathrm{d}\Omega}\right)(q_T \ll Q) = {} & \frac{1}{128(2\pi)^2 x_a x_b \mathcal{S}^2}\int \mathrm{d}^2k_{aT}\mathrm{d}^2k_{bT}\delta^{(2)}(\vec{k}_{aT}+\vec{k}_{bT}-\vec{q}_T)\\ & \times\Gamma^{+i;+j}(x_a,\vec{k}_{aT})\Gamma^{-k;-l}(x_b,\vec{k}_{bT})\sum_{\lambda_1,\lambda_2}\mathcal{M}^{ik}_{\lambda_1\lambda_2}\left(\mathcal{M}^{jl}_{\lambda_1\lambda_2}\right)^*.\end{aligned}$$

The gluon correlators Γ describe the TMD distributions of gluons in the nucleon that may or may not be polarized. If the partonic amplitude $\mathcal{M}$ is evaluated only at leading order, the gluon correlator reduces to a naive hadronic matrix element of two gluonic field strength tensors that was originally discussed in refs. [4, 12]. However, it was argued in refs. [13, 1] that the TMD correlators generally receive modifications in all-order factorization theorems due to a soft factors and a rapidity cut-off that determine the evolution of these objects. Instead of giving an exact definition of the TMD correlator the parameterization in terms of eight gluon TMDs is presented here [4, 12],

$$\begin{aligned}(3)\quad 2M^2\Gamma^{+i;+j}_U &= 2M^2\Gamma^{-i;-j}_U = M^2\delta^{ij}f_1^g + \left(k_T^i k_T^j - \frac{1}{2}\vec{k}_T^2\delta^{ij}\right)h_1^{\perp g},\\ 8M^3\Gamma^{+i;+j}_T &= -8M^3\Gamma^{-i;-j}_T = -4M^2\delta^{ij}\epsilon_T^{rs}k_T^r S_T^s f_{1T}^{\perp g} + 4M^2 i\epsilon_T^{ij}\left(\vec{k}_T\cdot\vec{S}_T\right)g_{1T}^{\perp g}\\ &\quad + M^2\left(S_T^{\{i}\epsilon_T^{j\}r}k_T^r + k_T^{\{i}\epsilon_T^{j\}r}S_T^r\right)h_{1T}^g + 2\left(k_T^{\{i}\epsilon_T^{j\}r}k_T^r\right)\left(\vec{k}\cdot\vec{S}_T\right)h_{1T}^{\perp g},\end{aligned}$$

with M being the nucleon mass, $\epsilon_T^{ij} \equiv \epsilon^{-+ij}$ and $a_T^{\{i}\epsilon_T^{j\}r} \equiv a_T^i\epsilon_T^{jr} + a_T^j\epsilon_T^{ir}$. Each of the TMDs in (3) is a function of x and $\vec{k}_T^2$. The transverse polarization of the nucleon is denoted by a spin vector $\vec{S}_T = |\vec{S}_T|(\cos\phi_{a/b}, \sin\phi_{a/b})$.

The leading order diagrams for $gg \to \gamma\gamma$ are shown in fig. 1. Due to gauge invariance those boxes are both UV- and IR-finite, hence no regularization procedure is needed. The amplitudes $\mathcal{M}^{ij}_{\lambda_1\lambda_2}$ in eq. (2) can be conveniently expressed in terms of the auxiliary vectors $\varepsilon_\pm = (1, \pm i)\mathrm{e}^{\mp i\phi}$ and $N_{gg} = \sum_q Q_q^2$,

$$\begin{aligned}(4)\qquad \mathcal{M}^{ik}_{\pm\pm} &= -2\alpha_s\alpha_{\mathrm{em}}N_{gg}\left(\varepsilon_-^i\varepsilon_-^k + \varepsilon_+^i\varepsilon_+^k - \varepsilon_\mp^i\varepsilon_\pm^k + f_s\varepsilon_\pm^i\varepsilon_\mp^k\right),\\ \mathcal{M}^{ik}_{\pm\mp} &= 2\alpha_s\alpha_{\mathrm{em}}N_{gg}\left(f_t\varepsilon_\mp^i\varepsilon_\mp^k + f_u\varepsilon_\pm^i\varepsilon_\pm^k + \varepsilon_-^i\varepsilon_+^k + \varepsilon_+^i\varepsilon_-^k\right).\end{aligned}$$

Furthermore, $f_s \equiv L(t,u)$, $f_t \equiv L(s,u)$, $f_u \equiv L(s,t)$, with L defined as

$$\begin{aligned}(5)\qquad L(x,y) = {} & 1 - \frac{x-y}{x+y}\left[\ln\left|\frac{x}{y}\right| - \mathrm{i}\pi\theta\left(-\frac{x}{y}\right)\right]\\ & + \frac{1}{2}\frac{x^2+y^2}{(x+y)^2}\left[\pi^2 + \left(\ln\left|\frac{x}{y}\right| - \mathrm{i}\pi\theta\left(-\frac{x}{y}\right)\right)^2\right].\end{aligned}$$

Inserting eqs. (4) and (3) into eq. (2) then leads to the following result for the unpolarized and single transverse-spin polarized cross sections in terms of the CS frame angles:

$$\begin{aligned}(6)\quad \frac{\mathrm{d}\sigma_{UU}^{gg}}{\mathrm{d}^4q\mathrm{d}\Omega} &= \sigma_0^{gg}\Big[\mathcal{F}_1(\theta)\mathcal{C}\left[f_1^g f_1^g\right] + \mathcal{F}_2(\theta)\mathcal{C}\left[w_5 h_1^{\perp g} h_1^{\perp g}\right] + \cos(2\phi)\Big\{\mathcal{F}_3(\theta)\Big(\mathcal{C}\left[w_1 h_1^{\perp g} f_1^g\right] \\ &\quad + \mathcal{C}\left[w_2 f_1^g h_1^{\perp g}\right]\Big)\Big\} + \cos(4\phi)\left\{\mathcal{F}_4(\theta)\mathcal{C}\left[w_4 h_1^{\perp g} h_1^{\perp g}\right]\right\}\Big],\\ (7)\quad \frac{\mathrm{d}\sigma_{TU}^{gg}}{\mathrm{d}^4q\mathrm{d}\Omega} &= \sigma_0^{gg}|\vec{S}_T|\sin\phi_a\Big[\mathcal{F}_1(\theta)\mathcal{C}\left[w_3 f_{1T}^{\perp g} f_1^g\right] \\ &\quad + \mathcal{F}_2(\theta)\left(\mathcal{C}\left[w_6 h_{1T}^g h_1^{\perp g}\right] + \mathcal{C}\left[w_7 h_{1T}^{\perp g} h_1^{\perp g}\right]\right) + \ldots\Big],\end{aligned}$$

where $\sigma_0^{gg} \equiv \alpha_{\mathrm{em}}^2\alpha_s^2 N_{gg}^2/(128\pi^2 x_a x_b \mathcal{S}^2)$, and where the ellipses denote additional terms that vanish upon ϕ-integration. We have defined $\mathcal{F}_1(\theta) = f_s^2 + |f_t|^2 + |f_u|^2 + 5$, $\mathcal{F}_2(\theta) = 2(f_s - 1)$, $\mathcal{F}_3(\theta) = f_s + \Re[f_u + f_t] - 1$, $\mathcal{F}_4(\theta) = f_u f_t^* + f_t f_u^* + 2$, and

$$(8)\quad \mathcal{C}\left[w f_1 f_2\right] \equiv \int \mathrm{d}^2k_{aT}\mathrm{d}^2k_{bT}\delta^{(2)}(\vec{k}_{aT} + \vec{k}_{bT} - \vec{q}_T)w(\vec{k}_{aT}, \vec{k}_{bT})f_1(x_a, \vec{k}_{aT}^2)f_2(x_b, \vec{k}_{bT}^2).$$

Defining $(ab)_\pm \equiv (a_1b_1 \pm a_2b_2)/2M^2$ and $[ab]_\pm \equiv (a_1b_2 \pm a_2b_1)/2M^2$, the following weights appear in (6) and (7) with $\vec{q}_T = (q_T, 0)$:

$$\begin{aligned}(9)\qquad & w_1 = -2(k_{aT}k_{aT})_-, \qquad w_2 = -2(k_{bT}k_{bT})_-, \qquad w_3 = \frac{1}{M}k_{aT,2},\\ & w_4 = (k_{aT}k_{bT})_-^2 - [k_{aT}k_{bT}]_+^2, \qquad w_5 = [k_{aT}k_{bT}]_-^2 - (k_{aT}k_{bT})_+^2,\\ & w_6 = \frac{1}{2M}\left((k_{bT}k_{bT})_+ k_{aT,2} - 2(k_{aT}k_{bT})_+ k_{bT,2}\right),\\ & w_7 = -\frac{2}{M}(k_{aT}k_{bT})_+[k_{aT}k_{bT}]_- k_{aT,1}.\end{aligned}$$

It is evident from eq. (6) that azimuthal modulations are generated in the gluon sector by the function $h_1^{\perp,g}$, the distribution of linearly polarized gluons in an unpolarized nucleon [4]. In particular an experimental measurement of the $\cos(4\phi)$-modulation gives a clean access to this function since this modulation is absent in the quark sector. Furthermore, it is possible to determine the sign of the function $h_1^{\perp,g}$ by measuring the $\cos(2\phi)$-modulation. This observable also receives contributions from the quark Boer-Mulders function [14] in the TMD formalism (cf. eq. (1)). The similarity of the quark sector to the lepton pair production makes the photon pair production particularly powerful when considered in combination with the Drell-Yan process where quark contributions can be cleanly determined.

The size of the effects in eq. (6) have been numerically estimated in [9] under the assumption of a saturation of a positivity bound for the function $h_1^{\perp,g}$ [4]. It was found that the resulting azimuthal modulations in eq. (6) are feasible already at RHIC energies. It is interesting to note that recent model calculations give support to the assumption of a saturation of the positivity bound and a rise at small-x [15, 16]. No

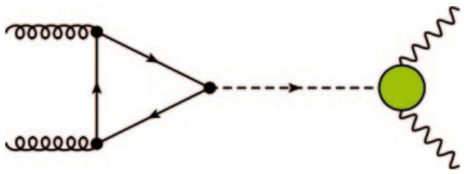

Fig. 2. – Higgs boson production and decay via $gg \to H \to \gamma\gamma$.

equivalent model calculation exist for the gluon Sivers function $f_{1T}^{\perp,g}$ in eq. (7). Although a gluon Sivers effect in photon pair production is dominant for a saturated gluon Sivers function [9] this TMD might be smaller in reality. It is up to an experimental investigation to decide this question. Also note that weighted cross sections of the form $\langle F\rangle_{XX} \equiv \int \mathrm{d}^2q_T \mathrm{d}\phi F(q_T,\phi)(\mathrm{d}\sigma_{XX}/\mathrm{d}^4q\mathrm{d}\Omega)$ may help in disentangling the various terms in (6) and (7). For instance, $\langle q_T^4 \cos(4\phi)\rangle_{UU} \propto h_1^{\perp(2),g}(x_a)h_1^{\perp(2),g}(x_b)$, $\langle q_T^2 \cos(2\phi)\rangle_{UU} \propto h_1^{\perp(2),g}(x_a)f_1^g(x_b)$, and $\langle 1\rangle_{TU} \propto f_{1T}^{\perp(1),g}(x_a)f_1^g(x_b)$, with k_T moments of $f_{1T}^{\perp,g}$ and $h_1^{\perp,g}$. The structure proportional to $\mathcal{F}_2$ does not contribute to the first two azimuthally independent moments $\langle(q_T^2)^n\rangle_{UU}$ for $n = 0, 1$. The consequence on this feature is that linearly polarized gluons will not contribute to the q_T-integrated photon pair cross section. On the other hand it leads to a characteristic double node in the q_T-distribution of this term caused by linearly polarized gluons. This will be discussed in the next section.

3. – Linearly polarized gluons and Higgs boson properties

If the c.m. energy of the protons becomes large enough, *e.g.*, at the LHC, Higgs bosons can be produced by gluon fusion via a top-quark loop, with a decay channel into a photon pair (cf. fig. 2). In this way the channel $gg \to H \to \gamma\gamma$ becomes a competing process to the quark box shown in fig. 1. The Higgs contribution was calculated in the TMD formalism in ref. [10]. It was found that the Higgs channel modifies the azimuthally independent cross section (6). To be specific, the prefactors $\mathcal{F}_1$ and $\mathcal{F}_2$ receive additional contributions from the Higgs boson diagram in fig. 2 of the following form:

$$\begin{aligned} \mathcal{F}_1^{H/A} &= f_s^2 + |f_t|^2 + |f_u|^2 + 5 - 2(f_s \mp 1)\Re[f_{H/A}] + 2|f_{H/A}|^2, \\ \mathcal{F}_2^{H/A} &= 2(f_s - 1) \pm 2(f_s \mp 1)\Re[f_{H/A}] \mp 2|f_{H/A}|^2, \end{aligned} \tag{10}$$

where H indicates a scalar Higgs boson predicted by the Standard Model (SM) and A a pseudoscalar Higgs predicted by models beyond the SM. The coupling of the pseudoscalar Higgs to the top-quark is assumed to be of the simple form $g_t i\gamma_5$ instead of 1. The objects $f_{H/A}$ are proportional to the amplitude pictured by fig. 2 and defined as

$$f_{H/A} = \frac{\sqrt{2}G_F s^2}{128\pi^2 N_{gg}} \frac{A_{gg\to H/A}(s)A_{H/A\to\gamma\gamma}(s)}{s - m_H^2 + i\Gamma_H m_H}, \tag{11}$$

where G_F is the Fermi constant, m_H, Γ_H the Higgs mass and width, respectively. $A_{gg\to H/A}$ describes the gluon fusion to a Higgs boson via a top-quark and $A_{H/A\to\gamma\gamma}$ the Higgs decay into a photon pair. With $\tau = s/(4m_t^2)$ and $\tilde{\tau} = s/(4m_W^2)$ (m_t, m_W the top and W mass) the fusion and decay triangular diagrams read,

$$\begin{aligned} A_{gg\to H} &= 2(\tau + (\tau - 1)J(\tau))/\tau^2, \qquad A_{gg\to A} = g_t 2J(\tau)/\tau, \\ A_{HH\to\gamma\gamma} &= 4A_{gg\to H}/3 - (2\tilde{\tau}^2 + 3\tilde{\tau} + 3(2\tilde{\tau} - 1)J(\tilde{\tau}))/\tilde{\tau}^2), \\ A_{A\to\gamma\gamma} &= 4A_{gg\to A}/3, \end{aligned} \tag{12}$$

with

$$J(\tau) = \begin{cases} -\dfrac{1}{4}\left(\ln\left(\dfrac{1+\sqrt{1-1/\tau}}{1-\sqrt{1-1/\tau}}\right) - i\pi\right)^2, & \tau > 1, \\ \arcsin^2\left(\sqrt{\tau}\,\right), & \tau \leq 1. \end{cases} \tag{13}$$

Note that the scalar Higgs can additionally decay via a W-triangular loop whereas only a decay via a top-loop is assumed for the pseudoscalar Higgs.

It is evident from eq. (11) that the Higgs production dominates only if the pair mass of the photons is in the vicinity of the Higgs mass, $s = Q^2 = m_H^2$. In this case the pole of the Higgs propagator is probed. On the other hand the box dominates for $s \neq m_H^2$ due to the suppression of the Fermi constant G_F in eq. (11). Interestingly, the prefactor $\mathcal{F}_2^{H/A}$ of the contribution caused by the distribution of linearly polarized gluons $h_1^{\perp,g}$ in (10) is sensitive to the parity of the Higgs boson. In particular at $s = m_H^2$ one has approximately $\mathcal{F}_2^{H/A}/\mathcal{F}_1^{H/A} \simeq \mp 1$. Hence, this feature could in principle be used to experimentally determine the Higgs parity by measuring the q_T-distribution at small q_T. The double node characteristics of the contribution $\mathcal{C}[w_5 h_1^{\perp,g} h_1^{\perp,g}]$ in (6) may be helpful for this measurement. Naturally, the feasibility of this procedure depends not only on the size of the function $h_1^{\perp,g}$ but also on the decay width Γ_H of the Higgs boson and consequently on the experimental ability to resolve the pair mass of the photons in the vicinity $Q^2 \simeq m_H^2$.

4. – Conclusions

Photon pair production in proton collisions is a promising process to study the gluonic spin substructure of the nucleon encoded in terms of gluon TMDs. This process does not suffer from the theoretical shortcomings that were found in other processes with hadronic final states, that is, TMD factorization is expected to hold in photon pair production. While on the one hand it is challenging to realize measurements of this process due to the detection of prompt photons it would be rewarding on the other hand. In combination with Drell-Yan data information on all gluon TMDs could be extracted from photon pair data at RHIC, in particular on the gluon Sivers function $f_{1T}^{\perp,g}$ and the distribution of linearly polarized gluons, $h_1^{\perp,g}$. The latter can also be studied at the LHC and may be used in the context of the search for the Higgs boson.

REFERENCES

[1] Collins J. C., *Foundations of perturbative QCD*, in *Cambridge Monographs on Particle Physics, Nuclear Physics and Cosmology*, Vol. **32** (Cambridge University Press) 2011; Ji X. D., Ma J. P. and Yuan F., *Phys. Rev. D*, **71** (2005) 034005; *Phys. Lett. B*, **597** (2004) 299; Collins J. C. and Metz A., *Phys. Rev. Lett.*, **93** (2004) 252001; Collins J. C., Rogers T. C. and Stasto A. M., *Phys. Rev. D*, **77** (2008) 085009; Collins J. C. and Soper D. E., *Nucl. Phys. B*, **193** (1981) 381; *Nucl. Phys. B*, **213** (1983) 545(E).

[2] Barone V., Bradamante F. and Martin A., *Prog. Part. Nucl. Phys.*, **65** (2010) 267.

[3] Sivers D. W., *Phys. Rev. D*, **41** (1990) 83; **43** (1991) 261.

[4] Mulders P. J. and Rodrigues J., *Phys. Rev. D*, **63** (2001) 094021; Ji X. D., Ma J. P. and Yuan F., *JHEP*, **0507** (2005) 020.

[5] Boer D. and Vogelsang W., *Phys. Rev. D*, **69** (2004) 094025; Bacchetta A. *et al.*, *Phys. Rev. Lett.*, **99** (2007) 212002.

[6] Yuan F., *Phys. Rev. D*, **78** (2008) 014024; Kang Z. B. and Qiu J. W., *Phys. Rev. D*, **78** (2008) 034005; Yuan F. and Zhou J., *Phys. Lett. B*, **668** (2008) 216; Kang Z. B. *et al.*, *Phys. Rev. D*, **78** (2008) 114013; Anselmino M. *et al.*, *Phys. Rev. D*, **70** (2004) 074025; Boer D., Mulders P. J. and Pisano C., *Phys. Rev. D*, **80** (2009) 094017.

[7] Collins J. C. and Qiu J. W., *Phys. Rev. D*, **75** (2007) 114014; Rogers T. C. and Mulders P. J., *Phys. Rev. D*, **81** (2010) 094006; Bomhof C. J., Mulders P. J. and Pijlman F., *Phys. Lett. B*, **596** (2004) 277.

[8] Boer D. *et al.*, *Phys. Rev. Lett.*, **106** (2011) 132001.

[9] Qiu J.-W., Schlegel M. and Vogelsang W., *Phys. Rev. Lett.*, **107** (2011) 062001.

[10] Boer D., den Dunnen W. J., Pisano C., Schlegel M. and Vogelsang W., arXiv: 1109.1444 [hep-ph].

[11] Arnold S., Metz A. and Schlegel M., *Phys. Rev. D*, **79** (2009) 034005.

[12] Meissner S., Metz A. and Goeke K., *Phys. Rev. D*, **76** (2007) 034002.

[13] Aybat S. M. and Rogers T. C., *Phys. Rev. D*, **83** (2011) 114042.

[14] Boer D. and Mulders P. J., *Phys. Rev. D*, **57** (1998) 5780.

[15] Metz A. and Zhou J., *Phys. Rev. D*, **84** (2011) 051503.

[16] Dominguez F., Qiu J.-W., Xiao B.-W. and Yuan F., arXiv:1109.6293 [hep-ph].