Direct Probes of Linearly Polarized Gluons inside Unpolarized Hadrons

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We show that the unmeasured distribution of linearly polarized gluons inside unpolarized hadrons can be directly probed in jet or heavy quark pair production both in electron-hadron and hadron-hadron collisions. We present expressions for the simplest $\cos 2\phi$ asymmetries and estimate their maximal value in the particular case of electron-hadron collisions. Measurements of the linearly polarized gluon distribution in the proton should be feasible in future EIC or LHeC experiments.

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Although quarks and gluons are confined within hadrons, tests of their fundamental properties are possible through scattering processes. It has become clear that quarks are in general spin-polarized within unpolarized hadrons, with polarization directions and magnitudes that depend on their transverse momentum and flavor. This nontrivial feature of hadron structure shows itself through specific angular asymmetries in scattering processes [1–3] that have been studied in a number of experiments [4–7]. Quark spin-polarization is also supported by first-principle lattice QCD calculations [8]. What has received much less attention is that gluons can exhibit a similar property, i.e. they can be linearly polarized inside an unpolarized hadron. In this letter we propose measurements which are directly sensitive to this unexplored gluon distribution.

Thus far experimental and theoretical investigations of gluons inside hadrons have focussed on their momentum and helicity distributions. The gluon density $g(x)$ describing the distribution of unpolarized gluons with a collinear momentum fraction x in an unpolarized hadron, integrated over transverse momentum $\boldsymbol{p}_T$, has been extracted with considerable precision from measurements of high energy electron-proton collisions at HERA (DESY). This distribution enters the structure function F_L in inclusive deep inelastic scattering (DIS) at order α_s, and it drives the evolution of sea quark distributions at small values of x. The unintegrated gluon distribution $g(x, \boldsymbol{p}_T^2)$ enters less inclusive reactions where the transverse momentum of the gluons is taken into account, such as semi-inclusive deep inelastic scattering or dijet production in hadronic collisions. In these cases the gluons are not necessarily unpolarized, even if the parent hadron itself is unpolarized. In fact, because of their spin-orbit couplings, the gluons can obtain a linear polarization.

Information on linearly polarized gluons in a hadron is formally encoded in the hadron matrix element of a correlator of the gluon field strengths $F^{\mu\nu}(0)$ and $F^{\nu\sigma}(\lambda)$ evaluated at fixed light-front time $\lambda^+ = \lambda{\cdot}n = 0$, where n is a lightlike vector conjugate to the parent hadron's four-momentum P. Specifically, the gluon content of an unpolarized hadron at leading twist (omitting gauge links) for a gluon momentum $p = x\,P + p_T + p^- n$ is described by the correlator [9]

$$\begin{aligned}\Phi_g^{\mu\nu}(x,\boldsymbol{p}_T) \;&=\; \frac{n_\rho\, n_\sigma}{(p{\cdot}n)^2}\int \frac{d(\lambda{\cdot}P)\, d^2\lambda_T}{(2\pi)^3}\; e^{ip\cdot\lambda} \\ &\qquad \times \langle P|\,\mathrm{Tr}\left[\,F^{\mu\rho}(0)\,F^{\nu\sigma}(\lambda)\,\right]|P\rangle\,\big]_{\mathrm{LF}} \\ &= \frac{-1}{2x}\left\{g_T^{\mu\nu}\, f_1^g - \left(\frac{p_T^\mu p_T^\nu}{M^2} + g_T^{\mu\nu}\,\frac{\boldsymbol{p}_T^2}{2M^2}\right) h_1^{\perp\, g}\right\}, \quad (1)\end{aligned}$$

with $p_T^2 = -\boldsymbol{p}_T^2$, $g_T^{\mu\nu} = g^{\mu\nu} - P^\mu n^\nu/P{\cdot}n - n^\mu P^\nu/P{\cdot}n$. This defines the transverse momentum dependent distribution functions (TMDs) $f_1^g(x,\boldsymbol{p}_T^2)$ representing the unpolarized gluon distribution, at fixed light-front time whereas $h_1^{\perp\, g}(x,\boldsymbol{p}_T^2)$ is the quadrupole distribution for linearly polarized gluons in an unpolarized hadron. It is named $h_1^{\perp\, g}$, because of its resemblance to the transversely polarized quark distribution inside an unpolarized hadron $h_1^{\perp\, q}$ (also frequently referred to as Boer-Mulders function) [1]. There are notable differences though: the T-odd distribution $h_1^{\perp\, q}$ for quarks is a chiral-odd distribution (chirality-flip) and it is also odd in p_T (it enters as a rank 1 tensor). It is zero in the absence of initial or final state interactions (ISI/FSI) [10–12]. The T-even distribution $h_1^{\perp\, g}$ for gluons describes a $\Delta L = 2$ helicity-flip distribution, through a second rank tensor in the relative transverse momentum p_T (p_T-even). Since an imaginary phase is not required, it can in principle be nonzero in the absence of ISI or FSI. Nevertheless, as any TMD, $h_1^{\perp\, g}$ can receive contributions from ISI or FSI and therefore they can be process dependent; in other words, non-universal.

Thus far no experimental studies of the function $h_1^{\perp\, g}$ have been performed. It has been pointed out [13] that it contributes to the so-called dijet imbalance in hadronic collisions, which is commonly used to extract the average

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partonic intrinsic transverse momentum. In this case, it enters the observable as a convolution of two $h_1^{\perp\, g}$ functions, similarly to the double Boer-Mulders effect which leads to a large $\sin\theta\cos 2\phi$ term and the leading-twist violation of the Lam-Tung relation in lepton pair production [2, 3]. It is in principle possible to isolate the contribution from the $h_1^{\perp\, g}$ functions by appropriate weighting of the planar angular distribution, but that is likely too difficult to do in practice.

Given its unique nature, it would be very interesting to obtain an extraction of $h_1^{\perp\, g}$ in a simple manner. This is possible, since unlike $h_1^{\perp\, q}$, it does not need to appear in pairs. In this letter we will discuss several new ways to extract the linear gluon polarization from observables that involve only a single $h_1^{\perp\, g}$ and, in principle, allow a check of its dependence on ISI or FSI.

The first processes we will consider are semi-inclusive DIS to two heavy quarks or to two jets: $e\, p \to e'\, Q\, \bar{Q}\, X$ and $e\, p \to e'\,\mathrm{jet}\,\mathrm{jet}\, X$. Since these involve final-state heavy quarks or jets, they require high energy colliders, such as a future Electron-Ion Collider (EIC) or the Large Hadron electron Collider (LHeC) proposed at CERN.

First we consider the electroproduction of heavy quarks, $e(\ell)+h(P) \to e(\ell')+Q(K_1)+\bar{Q}(K_2)+X$, where the four-momenta of the particles are given within brackets, and the quark-antiquark pair in the final state is almost back-to-back in the plane perpendicular to the direction of the exchanged photon and hadron. The calculation proceeds along the lines explained in Refs. [13, 14]. We obtain for the cross section integrated over the angular distribution of the back-scattered electron $e(\ell')$:

$$\frac{d\sigma}{dy_1\, dy_2\, dy\, dx_{\scriptscriptstyle B}\, d^2\boldsymbol{q}_{\scriptscriptstyle T} d^2\boldsymbol{K}_\perp} = \delta(1-z_1-z_2) \times \frac{\alpha^2\alpha_s}{\pi s M_\perp^2}\, \frac{(1+y x_{\scriptscriptstyle B})}{y^5 x_{\scriptscriptstyle B}} \left[A + \frac{\boldsymbol{q}_{\scriptscriptstyle T}^2}{M^2}\, B\, \cos 2(\phi_{\scriptscriptstyle T}-\phi_\perp)\right]. \quad (2)$$

This expression involves the standard DIS variables: $Q^2=-q^2$, where q is the momentum of the virtual photon, $x_{\scriptscriptstyle B} = Q^2/2P\cdot q, y = P\cdot q/P\cdot \ell$ and $s=(\ell+P)^2 = 2\,\ell\cdot P = 2\,P\cdot q/y = Q^2/x_B y$. Furthermore, we have for the jet momenta $K_{i\perp}^2 = -\boldsymbol{K}_{i\perp}^2$ and introduced the rapidities y_i for the jet momenta (along photon-target direction). We denote the heavy (anti)quark mass with M_Q. For the partonic subprocess we have $p+q = K_1+K_2$, implying $z_1+z_2=1$, where $z_i = P\cdot K_i/P\cdot q$. We introduced the sum and difference of the transverse jet momenta, $K_\perp = (K_{1\perp}-K_{2\perp})/2$ and $q_{\scriptscriptstyle T} = K_{1\perp}+K_{2\perp}$ with $|q_{\scriptscriptstyle T}| \ll |K_\perp|$. In that situation, we can use the approximate transverse jet momenta $K_{1\perp}\approx K_\perp$ and $K_{2\perp}\approx -K_\perp$ denoting $M_{i\perp}^2 \approx M_\perp^2 = M_Q^2 + \boldsymbol{K}_\perp^2$. The azimuthal angles of $\boldsymbol{q}_{\scriptscriptstyle T}$ and $\boldsymbol{K}_\perp$ are denoted by $\phi_{\scriptscriptstyle T}$ and $\phi_\perp$, respectively. The functions A and B in general depend on $x_{\scriptscriptstyle B}, y, z(\equiv z_2), Q^2/M_\perp^2, M_Q^2/M_\perp^2$, and $\boldsymbol{q}_{\scriptscriptstyle T}^2$.

The explicit expressions for the angular independent part A involving only f_1^g will be presented elsewhere, along with other results. For the coefficient B of the $\cos 2(\phi_{\scriptscriptstyle T}-\phi_\perp)$ angular distribution we obtain

$$B^{eh\to eQ\bar{Q}X} = \sum_Q e_Q^2\, h_1^{\perp\, g}(x,\boldsymbol{q}_{\scriptscriptstyle T}^2)\mathcal{B}^{eg\to eQ\bar{Q}}\,, \quad (3)$$

with

$$\mathcal{B}^{eg\to eQ\bar{Q}} = \frac{1}{2}\,\frac{z(1-z)}{D^3}\left(1-\frac{M_Q^2}{M_\perp^2}\right) a(y) \times \left[\Big(2\, z(1-z)\, b(y) - 1\Big)\frac{Q^2}{M_\perp^2} + 2\,\frac{M_Q^2}{M_\perp^2}\right]\,, \quad (4)$$

where the denominator D is defined as

$$D \equiv D\left(z, \frac{Q^2}{M_\perp^2}\right) = 1 + z(1-z)\frac{Q^2}{M_\perp^2}\,, \quad (5)$$

and we have introduced the following functions of y,

$$a(y) = 2 - y(2-y)\,, \qquad b(y) = \frac{6-y(6-y)}{2-y(2-y)}\,. \quad (6)$$

One observes that the magnitude B of the $\cos 2\phi$ asymmetry, where $\phi = \phi_{\scriptscriptstyle T}-\phi_\perp$, is determined by $h_1^{\perp\, g}$ and that if Q^2 and/or M_Q^2 are of the same order as $K_\perp^2$, the coefficient B is not power suppressed. Since $h_1^{\perp\, g}$ is completely unknown, we estimate the maximum asymmetry that is allowed by the bound:

$$|h_1^{\perp\, g(2)}(x)| \le \frac{\langle p_{\scriptscriptstyle T}^2\rangle}{2M^2}\, f_1^g(x)\,, \quad (7)$$

that we derived from the spin density matrix given in [9] in the way presented in Ref. [15]. The superscript (2) denotes the $n=2$ transverse moment. Transverse moments of TMDs are defined as: $f^{(n)}(x) \equiv \int d^2\boldsymbol{p}_{\scriptscriptstyle T}\, \left(\boldsymbol{p}_{\scriptscriptstyle T}^2/2M^2\right)^n f(x,\boldsymbol{p}_{\scriptscriptstyle T}^2)$. If we select $Q^2 = M_Q^2 = K_\perp^2/4$, $y_1 = y_2$, the asymmetry ratio

$$\left|\frac{\int d^2\boldsymbol{q}_{\scriptscriptstyle T}\, \boldsymbol{q}_{\scriptscriptstyle T}^2\, \cos 2(\phi_{\scriptscriptstyle T}-\phi_\perp)\, d\sigma}{\int d^2\boldsymbol{q}_{\scriptscriptstyle T}\, \boldsymbol{q}_{\scriptscriptstyle T}^2\, d\sigma}\right| = \frac{\int d\boldsymbol{q}_{\scriptscriptstyle T}^2\; \boldsymbol{q}_{\scriptscriptstyle T}^4\, |B|}{2M^2 \int d\boldsymbol{q}_{\scriptscriptstyle T}^2\; \boldsymbol{q}_{\scriptscriptstyle T}^2\, A}\,, \quad (8)$$

is maximally around 13%, which we view as encouraging.

If one keeps the lepton plane angle ϕ_ℓ, there are other azimuthal dependences, such as a $\cos 2(\phi_\ell-\phi_{\scriptscriptstyle T})$, but the bound on the latter is at least a factor of 6 smaller than on $\cos 2(\phi_{\scriptscriptstyle T}-\phi_\perp)$.

The cross section for the process $e\,h \to e'\,\mathrm{jet}\,\mathrm{jet}\,X$ can be calculated in a similar way. The corresponding expressions can be obtained from Eqs. (3) and (4) with $M_Q = 0$. One can then also replace the rapidities of the outgoing particles, y_i, with the pseudo-rapidities $\eta_i = -\ln\left[\tan(\tfrac{1}{2}\theta_i)\right]$, θ_i being the polar angles of the final partons in the virtual photon-hadron cms frame.

We note that the measurement or reconstruction of the transverse momenta of the jets or heavy quarks is essential. Here it is assumed that the individual jet or heavy quark transverse momenta $K_{i\perp}$ can be reconstructed with an accuracy that is better than the magnitude of the sum of the transverse momenta $K_{1\perp} + K_{2\perp}$. Furthermore, we point out that an analogous asymmetry arises in QED, in the 'tridents' processes $\ell e(p) \to \ell\mu^+\mu^- e'(p' \,\mathrm{or}\, X)$ or $\mu^- Z \to \mu^- \ell\bar{\ell} Z$ [16–19]. This could be described by the distribution of linearly polarized photons inside a lepton, proton, or atom. QCD adds the twist that for gluons inside a hadron, ISI or FSI can considerably modify the result depending on the process.

Next we consider heavy quark production in hadronic collisions: $p\,p \to Q\,\bar{Q}\,X$. This does involve convolutions of two TMDs, but still in some cases only a single $h_1^{\perp\,g}$ appears, thereby avoiding a double suppression, assuming the function to be considerably smaller than the unpolarized gluon TMD. Here we have BNL's Relativistic Heavy Ion Collider (RHIC) and CERN's LHC in mind, but the measurements can equally well be done at Fermilab's Tevatron in $p\,\bar{p} \to Q\,\bar{Q}\,X$.

The cross section for the process

$$h_1(P_1)+h_2(P_2) \to Q(K_1)+\bar{Q}(K_2)+X \tag{9}$$

can be written in a way similar to the hadroproduction of two jets discussed in Ref. [13], in the following form

$$\begin{aligned}\frac{d\sigma}{dy_1 dy_2 d^2\boldsymbol{K}_{1\perp} d^2\boldsymbol{K}_{2\perp}} &= \frac{\alpha_s^2}{s M_\perp^2} \\ &\times \Big[A(\boldsymbol{q}_T^2) + B(\boldsymbol{q}_T^2)\boldsymbol{q}_T^2 \cos 2(\phi_T - \phi_\perp) \\ &\quad + C(\boldsymbol{q}_T^2)\boldsymbol{q}_T^4 \cos 4(\phi_T - \phi_\perp)\Big]. \end{aligned}\tag{10}$$

Besides $\boldsymbol{q}_T^2$, the terms A, B and C depend on other kinematic variables often not explicitly indicated, i.e. on z, $M_Q^2/M_\perp^2$ and on x_1, x_2, which are given by:

$$x_i = \frac{1}{\sqrt{s}}\left(M_{1\perp}\, e^{\sigma_i y_1} + M_{2\perp}\, e^{\sigma_i y_2}\right), \tag{11}$$

with $\sigma_1 = +1$ and $\sigma_2 = -1$.

The terms A, B, and C contain convolutions of various TMDs. Schematically,

$$\begin{aligned} A: &\quad f_1^q \otimes f_1^{\bar{q}},\ f_1^g \otimes f_1^g\,, \\ B: &\quad h_1^{\perp\,q} \otimes h_1^{\perp\,\bar{q}},\ \frac{M_Q^2}{M_\perp^2} f_1^g \otimes h_1^{\perp\,g}\,, \\ C: &\quad h_1^{\perp\,g} \otimes h_1^{\perp\,g}\,. \end{aligned}$$

Terms with higher powers in $M_Q^2/M_\perp^2$ are left out. The convolutions $h_1^{\perp\,q} \otimes h_1^{\perp\,\bar{q}}$ in B and $h_1^{\perp\,g} \otimes h_1^{\perp\,g}$ in C have already been addressed in [13] for dijet production. Here we will focus on B, for which we find:

$$B = \mathcal{B}^{q\bar{q}\to Q\bar{Q}} + (M_Q^2/M_\perp^2)\,\mathcal{B}^{gg\to Q\bar{Q}}\,, \tag{12}$$

where

$$\begin{aligned}\mathcal{B}^{q\bar{q}\to Q\bar{Q}} &= \frac{N^2-1}{N^2}\, z^2(1-z)^2 \left(1 - \frac{M_Q^2}{M_\perp^2}\right) \\ &\times\Big[\mathcal{H}^{q\bar{q}}(x_1,x_2,\boldsymbol{q}_T^2) + \mathcal{H}^{\bar{q}q}(x_1,x_2,\boldsymbol{q}_T^2)\Big], \end{aligned}\tag{13}$$

$$\mathcal{B}^{gg\to Q\bar{Q}} = \frac{N}{N^2-1}\, \mathcal{B}_1\, \mathcal{H}^{gg}(x_1,x_2,\boldsymbol{q}_T^2)\,, \tag{14}$$

with N being the number of quark colors and

$$\mathcal{B}_1 = z(1-z)\left(z^2 + (1-z)^2 - \frac{1}{N^2}\right)\left(1 - \frac{M_Q^2}{M_\perp^2}\right). \tag{15}$$

Instead of presenting the convolution expressions for $\mathcal{H}^{q\bar{q}}$ and $\mathcal{H}^{gg}$, we consider here the expressions which appear in the $\boldsymbol{q}_T^2/M^2$-weighted cross section (cf. [13]), for $M_1 = M_2 = M$. We encounter weighted integrals,

$$\begin{aligned}\pi \int d\boldsymbol{q}_T^2\, \left(\frac{\boldsymbol{q}_T^2}{M^2}\right)\, \boldsymbol{q}_T^2\, \mathcal{H}^{q\bar{q}}(x_1,x_2,\boldsymbol{q}_T^2) &= \\ 8\sum_{\text{flavors}} h_1^{\perp q(1)}(x_1)\, h_1^{\perp\bar{q}(1)}(x_2)\,, \end{aligned}\tag{16}$$

already discussed in [13], and

$$\begin{aligned}\pi \int d\boldsymbol{q}_T^2\, \left(\frac{\boldsymbol{q}_T^2}{M^2}\right)\, \boldsymbol{q}_T^2\, \mathcal{H}^{gg}(x_1,x_2,\boldsymbol{q}_T^2) &= \\ 4\left(h_1^{\perp g(2)}(x_1)\, f_1^g(x_2) + f_1^g(x_1)\, h_1^{\perp g(2)}(x_2)\right). \end{aligned}\tag{17}$$

The bound in Eq. (7) can again be used in the last expression to obtain the maximal asymmetry. Whether it is more important than the term in Eq. (16) depends strongly on the considered values of x_i and $M_Q^2/M_\perp^2$, and on whether one deals with $p\,p$ or $p\,\bar{p}$. In $p\,\bar{p}$ collisions and for $K_\perp^2$ not too large compared to M_Q^2, the contribution from $h_1^{\perp\,g}$ is expected to be the dominant one. The importance of the contribution from $h_1^{\perp\,q}$ can be assessed through a comparison to photon-jet production [14].

In Fig. 1 the origin of the factor $M_Q^2/M_\perp^2$ in the contribution of $h_1^{\perp\,g}$ to B is explained by a comparison to the ep case where it is absent.

We have not yet addressed the effects from the gauge link appearing in the correlator in Eq. (1). Even in relatively simple processes such as those discussed in this paper, they lead to different color flow factors for different diagrams in the partonic subprocess [20–22]. There are actually two different T-even four-gluon soft gluon pole matrix elements and hence functions $h_1^{\perp g(2)}(x)$, with color structures $f_{abe}\, f_{cde}$ and $d_{abe}\, d_{cde}$, respectively [21, 23]. In the $e\,q \to e'\,q\,Q\,\bar{Q}$ subprocess only the matrix element with the $f\,f$-structure appears, while in the $g\,g \to Q\,\bar{Q}$ subprocess the $d\,d$-structure dominates

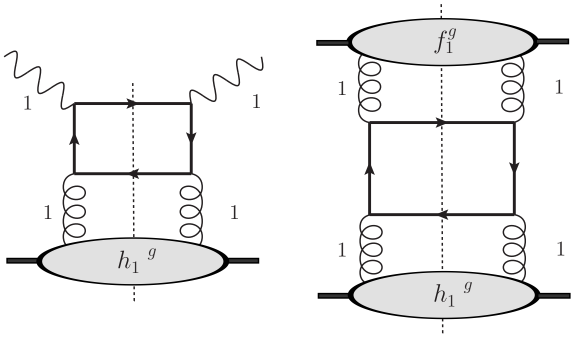

FIG. 1: *Examples of subprocesses contributing to the* $\cos 2\phi$ *asymmetries in* $e\,p \to e'\,Q\,\bar{Q}\,X$ *and* $p\,p \to Q\,\bar{Q}\,X$*, respectively. As the helicities of the photons and gluons indicate, the latter process requires helicity flip in quark propagators resulting in an* $M_Q^2/M_\perp^2$ *factor.*

(the $f\,f$-contribution is suppressed by $1/N^2$). Hence, although $h_1^{\perp\,g}$ does not require ISI or FSI to be nonzero, a process dependence is expected.

The four-gluon operator associated with $h_1^{\perp\,g(2)}$ leads to a process-dependent p_T broadening [24–26]. We thus expect the broadening Δp_T^2 in SIDIS, $(\Delta p_T^2)_{\rm DIS} \equiv \langle p_T^2 \rangle_{eA} - \langle p_T^2 \rangle_{ep}$, to be different from the one in hadron-hadron collisions, $(\Delta p_T^2)_{\rm hh} \equiv \langle p_T^2 \rangle_{pA} - \langle p_T^2 \rangle_{pp}$.

As a final point we comment on single transverse spin asymmetries (A_N) in heavy quark production processes. Spin-dependent ISI and FSI in A_N have been considered in Ref. [27]. It was concluded that no such effects arise in C-odd $c\,\bar{c}$ meson production, such as J/ψ production. Open charm, on the other hand, does allow one to probe ISI or FSI. The asymmetry A_N in D-meson production has been studied in [28–30] in terms of triple-gluon soft gluon pole correlators and in [31] in terms of the gluon Sivers effect. Because there are two different (f and d type) correlators to consider, the single-spin asymmetries in D and $\bar{D}$ meson production are found to be different. In contrast, in the unpolarized scattering case considered in this letter this problem does not arise, since only one operator contributes or dominates.

In summary, measurements of the azimuthal asymmetry of jet or heavy quark production in $e\,p$ and in $p\,p$ or $p\,\bar{p}$ collisions can directly probe the quadrupole distribution for linearly polarized gluons inside unpolarized hadrons. From a theoretical viewpoint these asymmetries are among the simplest TMD observables since the number of partonic subprocesses is in each case limited to just one type. This avoids having to consider complicated linear combinations of initial and final-state interactions. The relative simplicity of the proposed measurements (polarized beams are not required) suggests a promising prospect for the extraction of this gluon distribution in the future.

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