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Probing quark orbital angular momentum in electron-proton collisions

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In these proceedings, we present our proposal to extract quark orbital angular momentum (OAM) through exclusive production in electron-(longitudinally-polarized) proton collisions. Our analysis demonstrates that the azimuthal angular correlation between the transverse momentum of the scattered electron and the recoil proton is a sensitive probe of quark OAM. We also provide a numerical estimate of the asymmetry associated with this correlation for the kinematics accessible at EIC and EicC, aiming to demonstrate the feasibility of measuring our observable. Our work paves the way for the first measurement of quark OAM in relation to the Jaffe-Manohar spin sum rule.

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PoS(DIS2024)241

1. Introduction

Generalized Transverse Momentum-dependent Distributions (GTMDs) are the most comprehensive (2-) parton correlation functions for describing the structure of nucleons. Studying GTMDs is motivated by several factors. Functions such as form factors (FFs), one-dimensional Parton Distributions (PDFs), and their three-dimensional generalizations—Transverse Momentum-dependent parton Distributions (TMDs) and Generalized Parton Distributions (GPDs)—are specific kinematical projections of GTMDs. GTMDs therefore act as a “mother distribution" [1]. Second, GTMDs encapsulate more physics than TMDs and GPDs, as some GTMDs do not survive in the TMD/GPD limits (e.g., $F_{1,4}$ and $G_{1,1}$ GTMDs). Exploring the physics lost in these limits is particularly interesting. Third, for specific kinematics, the Fourier transform of GTMDs can be related to Wigner distributions [2]. Wigner distributions, encountered in the phase-space formulations of quantum mechanics and in other branches of physics such as Atomic, Molecular, and Optical physics, allow for five-dimensional imaging of nucleons. Furthermore, certain GTMDs can reveal non-trivial correlations between the orbital motion of partons and the spin of nucleons [3, 4]

$$\langle \vec{L}_{q/g} \cdot \vec{S}_N \rangle \sim F_{1,4}^{q/g} , \qquad \langle \vec{L}_{q/g} \cdot \vec{S}_{q/g} \rangle \sim G_{1,1}^{q/g} . \tag{1}$$

These compelling reasons underscore the importance of studying GTMDs, making them the “ultimate" functions we aim to access in experiments. In the next section, we identify a physical process that will enable the first measurement of quark orbital angular momentum (OAM) by exploiting connections with GTMDs; see the first equation in Eq. (1).

2. Main results

In Ref. [5], we identified for the first time a physical process sensitive to the quark GTMDs, known as exclusive double Drell-Yan in pion-nucleon collisions (see Fig. 1). In that process, we also identified an observable sensitive to quark OAM and spin-orbit correlation. However, there are two practical challenges: the count rate of the process is low (cross section is $\sim \alpha_{\rm em}^4$), and we are sensitive to the GTMDs only in the ERBL region. This is problematic because the OAM density $L^{q/g}(x,\xi) \sim \int d^2k_\perp F_{1,4}^{q/g}(x, k_\perp, \xi, \Delta_\perp)$ relates to OAM $L^{q/g}$ only at $\xi = 0$ (ξ is skewness). Thus, the challenge lies in extrapolating the distribution $F_{1,4}^{q/g}$, which we can access only in the ERBL region in that process, down to the forward limit where the OAM equation is applicable. Our

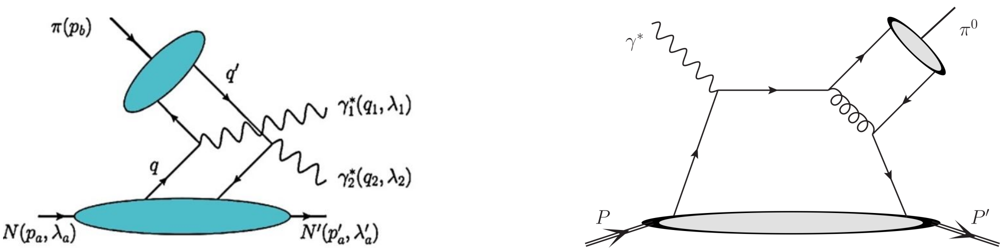

Figure 1: Left figure: Exclusive pion-nucleon double Drell-Yan process. Right figure: Exclusive pion production in electron-proton collisions.

PoS(DIS2024)241

current work [6] is motivated by these challenges. We propose examining exclusive π^0 production in electron-proton collisions to overcome the issues associated with double Drell-Yan. See Fig. 1 for a sketch of the process.

First, let us define the kinematics of the process under consideration: $e(l) + p(p, \lambda) \longrightarrow \pi^0(l_\pi) + e(l') + p(p', \lambda')$. The standard kinematic variables are defined as follows: $Q^2 = -q^2 = -(l - l')^2$, representing the photon's virtuality. The incoming (outgoing) electron's momentum is denoted as $l^\mu (l'^\mu)$. The $\gamma^* p$ center-of-mass energy is given by $W^2 = (p + q)^2$. In our calculation, we neglect the pion mass ($l_\pi^2 \approx 0$) to simplify the analysis. The skewness variable is defined as $\xi = (p^+ - p'^+)/(p^+ + p'^+) = -\Delta^+/2P^+ = x_B/2 - x_B$, where $x_B = Q^2/2p \cdot q$ represents the Bjorken scaling variable. The momentum transfer squared can be expressed as $t = (p - p')^2 = -4\xi^2 M^2 + \Delta_\perp^2/1 - \xi^2$, with M being the proton mass. We will also use z to denote the longitudinal momentum fraction of the π^0 carried by the quark.

To leading order in perturbative QCD calculations, there are four Feynman diagrams contributing. The schematic structure of the scattering amplitude A is given by

$$A \propto \int dx \int d^2k_\perp \, H(x, \xi, z, k_\perp, \Delta_\perp) \, f^q(x, \xi, k_\perp, \Delta_\perp) \int dz \phi_\pi(z), \tag{2}$$

where $H(..)$ denotes the hard part, $f^q(..)$ is the soft part from the proton side, and $\phi_\pi(..)$ is the pion's distribution amplitude (DA). Next, we perform a collinear twist expansion of the hard part in powers of $k_\perp$ and $\Delta_\perp$:

$$H(k_\perp, \Delta_\perp) = H(k_\perp = 0, \Delta_\perp = 0) + \left.\frac{\partial H(k_\perp, \Delta_\perp = 0)}{\partial k_\perp^\mu}\right|_{k_\perp = 0} k_\perp^\mu + \left.\frac{\partial H(k_\perp = 0, \Delta_\perp)}{\partial \Delta_\perp^\mu}\right|_{\Delta_\perp = 0} \Delta_\perp^\mu + \ldots \tag{3}$$

The leading-order term, or the twist-2 term, $H(k_\perp = 0, \Delta_\perp = 0)$, vanishes for our process, and the non-zero contribution arises from the next-to-leading order, or the twist-3 term. Care must be taken during this expansion to ensure that electromagnetic gauge invariance is preserved. To maintain gauge invariance, we employ the special propagator technique. When we pick up one factor of $k_\perp$ from the hard part, the amplitude is proportional to certain weighted moments of GTMDs. When we pick up one factor of $\Delta_\perp$ from the hard part, the amplitude is proportional to GPDs. Consequently, the scattering amplitudes are of twist-3 nature and are a convolution of moments of GTMDs (in other words, quark OAM) and GPDs. We note that we have omitted the $k_\perp$-dependence in the pion DA in Eq. (3) because it is of twist-3 nature. More specifically, if we include $k_\perp$ dependence in the pion DA, the leading power contribution would come from the coupling between the twist-3 DA and the transversity GPDs, as computed by the GK group [7]. For simplicity, we neglect this contribution in our work.

Scattering amplitudes depend on various angular correlations emerging from structures such as $\epsilon_\perp \times \Delta_\perp$ (involving moments of $F_{1,1}$ and $G_{1,1}$), $\epsilon_\perp \cdot S_\perp$ (involving moments of $F_{1,2}$ and $G_{1,2}$), and $\epsilon_\perp \cdot \Delta_\perp$ (involving moments of $F_{1,4}$ and $G_{1,4}$), where $\epsilon_\perp$ is the polarization vector of the virtual photon and $S_\perp$ is defined as $S_\perp^\mu = (0^+, 0^-, -i, \lambda)$. For the full expressions, see Ref. [6]. The main result of our work is the following expression for the polarization-dependent cross section:

$$\frac{d\sigma}{dt dQ^2 dx_B d\phi} = \frac{(N_c^2 - 1)^2 \alpha_{em}^2 \alpha_s^2 f_\pi^2 \xi^3 \Delta_\perp^2}{2N_c^4 (1 - \xi^2) Q^{10} (1 + \xi)} \left[1 + (1 - y)^2\right] \\ \times \lambda \sin(2\phi) \, 2a \, \mathrm{Re}\left[\left(i\mathcal{F}_{1,4} + i\mathcal{G}_{1,4}\right)\left(\mathcal{F}_{1,1}^* + \mathcal{G}_{1,1}^*\right)\right]\Big\}, \tag{4}$$

PoS(DIS2024)241

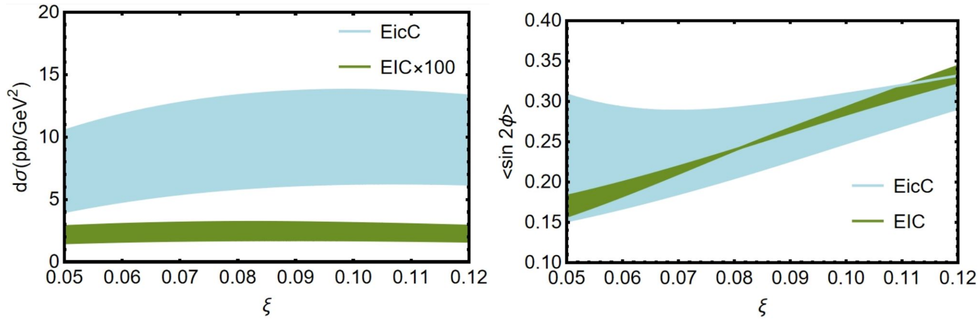

Figure 2: The unpolarized cross section is depicted in the left plot for EIC kinematics with $Q^2 = 10\,\text{GeV}^2$ and $\sqrt{s_{ep}} = 100\,\text{GeV}$, and for EicC kinematics with $Q^2 = 3\,\text{GeV}^2$ and $\sqrt{s_{ep}} = 16\,\text{GeV}$. The unpolarized cross section for the EIC case is scaled by a factor of 100. The right plot shows the average value of $\langle\sin(2\phi)\rangle$ as a function of ξ. The variable t is integrated over the range $[-0.5, \text{GeV}^2, -\frac{4\xi^2 M^2}{1-\xi^2}]$. For a detailed description of how we generated the error bands, we refer to our paper, Ref. [6].

where $a = \frac{2(1-y)}{1+(1-y)^2}$ and $f_\pi = 131$ MeV represents the π^0 decay constant. This expression clearly shows that the moment of the quark GTMD $F_{1,4}$, denoted as $\mathcal{F}_{1,4}$, and consequently the quark OAM, leaves a distinct signature through an azimuthal angular correlation of $\sin 2\phi$ in the longitudinal single target-spin asymmetry. This angle represents the correlation between the scattered lepton angle and the recoil proton ($\phi = \phi_{l_\perp} - \phi_{\Delta_\perp}$). We have also calculated the unpolarized cross section for this process, which is given by:

$$\frac{d\sigma}{dt dQ^2 dx_B d\phi} = \frac{(N_c^2-1)^2 \alpha_{em}^2 \alpha_s^2 f_\pi^2 \xi^3 \Delta_\perp^2}{2N_c^4 (1-\xi^2) Q^{10} (1+\xi)} \left[1 + (1-y)^2\right]$$
$$\times \left\{ \left[|\mathcal{F}_{1,1} + \mathcal{G}_{1,1}|^2 + |\mathcal{F}_{1,4} + \mathcal{G}_{1,4}|^2 + 2\frac{M^2}{\Delta_\perp^2} |\mathcal{F}_{1,2} + \mathcal{G}_{1,2}|^2 \right] \right.$$
$$\left. + \cos(2\phi) a \left[-|\mathcal{F}_{1,1} + \mathcal{G}_{1,1}|^2 + |\mathcal{F}_{1,4} + \mathcal{G}_{1,4}|^2 \right] \right\}.$$

The unpolarized cross section exhibits a distinctive $\cos 2\phi$ modulation. There are several surprises in this equation: First, we find that the quark Sivers function, through its relation with GTMD $F_{1,2}$, as $\text{Im}\,[F_{1,2}]\big|_{\Delta=0} = -f_{1T}^\perp$, leaves a trace in the unpolarized cross section. This is surprising because it is commonly known that the Sivers function can only be accessed through transversely polarized targets. What we observe here is very similar to how and why the gluon Sivers function also enters this process, as was first pointed out in Ref. [8]. Second, we find that we can probe the quark worm-gear function through its relation with GTMD $G_{1,2}$, as $\text{Re}\,[G_{1,2}]\big|_{\Delta=0} = g_{1T}$, using an unpolarized target. The distinctive method to isolate these helicity-flip terms is by approaching the limit $\Delta_\perp \to 0$, where these terms persist relative to others. Before presenting the numerical results below, we emphasize once again that since both unpolarized and polarized cross sections contribute at twist-3, the unpolarized cross section does not wash away the asymmetry.

For details on how we modeled the various non-perturbative functions, we refer to our paper Ref. [6]. Here, we discuss the consequences of certain endpoint singularities and discontinuities in the distributions that may pose challenges for factorization in this process. Before proceeding, we

PoS(DIS2024)241

present the full expression of one of the moments of the GTMDs appearing in our cross section:

$$\mathcal{F}_{1,4} = \int_{-1}^{1} dx \frac{x\xi \int d^2k_\perp k_\perp^2 F_{1,4}^{u+d}(x,\xi,\Delta_\perp,k_\perp)}{M^2(x+\xi-i\epsilon)^2(x-\xi+i\epsilon)^2} \times \int_0^1 dz \frac{\phi_\pi(z)(1+z^2-z)}{z^2(1-z)^2}\,. \tag{5}$$

First, we observe that the hard part diverges as z approaches 0 or 1, indicating an endpoint singularity that typically challenges collinear factorization. From a phenomenological perspective, regularization can be achieved by modifying the integration limits of z to $\int_{\langle p_\perp^2\rangle/Q^2}^{1-\langle p_\perp^2\rangle/Q^2} dz$ [9], where $\langle p_\perp^2\rangle$ denotes the mean squared transverse momentum of the quark inside the pion. We adopt $\langle p_\perp^2\rangle = 0.04$ GeV^2 for our numerical calculations, based on a fit to CLAS data (details in Ref. [6]). Additionally, for simplicity, we use the asymptotic form of the pion's DA, $\phi_\pi(z) = 6z(1-z)$. Second, the discontinuity of the derivative of quark GPDs at the endpoints $x = \pm\xi$ (as observed, for example, in Refs. [10]), along with the presence of double poles in the integrals at $x = \pm\xi$, may potentially introduce a divergent component in the cross section. To mitigate this issue, we shift the double pole from $\frac{1}{(x-\xi+i\epsilon)^2}$ to $\frac{1}{(x-\xi-\langle p_\perp^2\rangle/Q^2+i\epsilon)^2}$ (and similarly for the negative x region). These adjustments underscore the necessity of accounting for the transverse motion of quarks within pions and protons to ensure that factorization holds. A few additional points about the kinematic region of interest before discussing the numerical results. First, we emphasize the importance of the large skewness region to suppress gluon contributions. Second, we prioritize the large momentum transfer (t) region to minimize contributions from the Primakoff process (which scales as $\propto 1/t$). Interestingly, a similar azimuthal asymmetry with a $\sin 2\phi$ modulation, resembling the modulation observed in this study, arises from the interference between the Primakoff process and the contribution from the gluon GTMD $F_{1,4}$. This interference serves as a background for our process and is observable under these kinematic conditions, as detailed in Ref. [11].

In the left plot of Fig. 2, we display the unpolarized cross section as a function of skewness for achievable kinematics at EIC and EicC. We observe a significant magnitude of the unpolarized cross section at EicC energy, whereas it appears relatively small at EIC energy. This difference is likely due to the unpolarized cross section scaling as $1/Q^{10}$. However, it is worth noting that EIC may have the capability to probe smaller ξ regions at the same Q^2 as EicC, potentially offering greater leverage in constraining quark OAM in the small x region. Thus, there is complementarity between these two electron-ion colliders that can be envisioned. In the right plot of Fig. 2, we illustrate the $\sin 2\phi$ weighted asymmetry, defined as $\langle \sin(2\phi)\rangle = \frac{\int \frac{d\Delta\sigma}{d\mathcal{P}.\mathcal{S}.}\sin(2\phi)\, d\mathcal{P}.\mathcal{S}.}{\int \frac{d\sigma}{d\mathcal{P}.\mathcal{S}.} d\mathcal{P}.\mathcal{S}.}$, where $d\Delta\sigma = \sigma(\lambda=1) - \sigma(\lambda=-1)$, as a function of skewness. The asymmetries are substantial for both EIC and EicC kinematics, highlighting the azimuthal asymmetry $\sin 2\phi$ in exclusive π^0 production as a promising avenue for probing the quark OAM distribution.

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3. Summary

We propose extracting the quark OAM associated with the Jaffe-Manohar spin sum rule by measuring the azimuthal angular correlation $\sin 2\phi$ in exclusive π^0 production at EIC and EicC. This observable serves as a clean and sensitive probe of quark OAM for several reasons. Firstly, the azimuthal asymmetry is not a power correction, as both the unpolarized and longitudinal

PoS(DIS2024)241

polarization-dependent cross sections contribute at twist-3. Secondly, this process provides direct access to the quark GTMD $F_{1,4}$ in the DGLAP region for the first time. In addition to uncovering quark OAM, our work highlights another significant finding: the quark components of $F_{1,2}$ and $G_{1,2}$, equivalent to the Sivers function and the worm-gear function respectively, contribute to the unpolarized cross-section of this process in the forward limit. This result is particularly noteworthy since traditionally, the Sivers function and the worm-gear function are probed exclusively through transversely polarized targets. We computed the differential cross section within the collinear higher-twist expansion framework. Our numerical results reveal a substantial azimuthal asymmetry, crucially dependent on the quark OAM distribution. In the kinematic range accessible to EIC and EicC, our observable can be thoroughly investigated, paving the way for the first experimental extraction of the Jaffe-Manohar quark OAM distribution in the future.

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PoS(DIS2024)241