

Nucleon Energy Correlators as a Probe of Light-Quark Dipole Operators at the Electron-Ion Collider

Yingsheng Huang^{1,2,*}, Xuan-Bo Tong^{3,4,†} and Hao-Lin Wang^{5,6,‡}

¹*School of Physics, Central South University, Changsha 410083, China*

²*Department of Physics and Astronomy, University of Utah, Salt Lake City, Utah 84112, USA*

³*Department of Physics, University of Jyväskylä, P.O. Box 35, 40014 University of Jyväskylä, Finland*

⁴*Helsinki Institute of Physics, P.O. Box 64, 00014 University of Helsinki, Finland*

⁵*State Key Laboratory of Nuclear Physics and Technology, Institute of Quantum Matter, South China Normal University, Guangzhou 510006, China*

⁶*Guangdong Basic Research Center of Excellence for Structure and Fundamental Interactions of Matter, Guangdong Provincial Key Laboratory of Nuclear Science, Guangzhou 510006, China*

 (Received 25 August 2025; revised 9 February 2026; accepted 16 March 2026; published 2 April 2026)

We propose nucleon energy correlators (NECs) as a novel framework to probe electroweak light-quark dipole operators in deep inelastic scattering with an unpolarized nucleon. These operators encode chirality-flipping interactions, whose effects are usually quadratically suppressed in unpolarized cross sections. We construct a chiral-odd quark NEC that accesses quark transverse spin via azimuthal angle asymmetries in the energy flow of the target fragmentation region. These asymmetries serve as clean and powerful observables, enabling linear constraints on the quark dipole couplings. Unlike existing methods, our approach requires neither polarized nucleon beams nor final-state hadron identification, relying instead on fully inclusive calorimetric measurements. This Letter establishes one of the first applications of energy correlator observables to new physics searches and opens a promising direction for precision studies of chirality-flipping effects at electron-ion colliders.

DOI: [10.1103/7y5n-cscx](https://doi.org/10.1103/7y5n-cscx)

Introduction—The search for new physics (NP) beyond the standard model (SM) has been a major focus of collider experiments. The lack of direct signals suggests that NP may emerge at an energy scale Λ well beyond the current experimental reach. The standard model effective field theory (SMEFT) offers a systematic framework to describe such NP effects via higher-dimensional operators built from SM fields [1–3]. Among these, dimension-6 electroweak (EW) dipole operators involving light fermions are particularly intriguing for their chirality-flipping structure [4] [5–20]. Such effects are highly suppressed in the SM but essential for understanding anomalous magnetic moments (AMM) and P , T violating electric dipole moments (EDM) of leptons and baryons [21–36]. However, constraints on these operators at colliders remain weak because the interference between dipole and SM amplitudes, scaling

as $\mathcal{O}(\Lambda^{-2})$, typically vanishes in unpolarized observables [9–16]. Consequently, the leading effects arise at $\mathcal{O}(\Lambda^{-4})$, resulting in reduced sensitivity and contamination from dimension-8 operators.

A new strategy to overcome this issue was recently proposed in [17,18], which leverages the transverse spin of fermions. Since transverse spin states are superpositions of helicity eigenstates, they allow interferences between opposite-helicity components, restoring sensitivity to the dipole-SM interference at $\mathcal{O}(\Lambda^{-2})$. Building on this insight, several observables have been designed for the Electron-Ion Collider (EIC) and future lepton colliders [17–20,37]. While transverse polarization of electron beams can be generated and controlled in accelerators, accessing transversely polarized quarks is more difficult due to color confinement. Existing approaches require nonperturbative correlations between quark spin and a transverse reference vector, either via nucleon polarization or via the transverse momenta of final-state hadrons, encoded in transversity parton distribution functions (PDFs) and dihadron fragmentation functions [17,19]. Nevertheless, these methods either demand polarized nucleon beams, which compromise luminosity, or rely on semi-inclusive measurements that require particle identification and multiparticle tracking.

To address these limitations, we propose nucleon energy correlators (NECs) as an innovative approach to probe

*Contact author: yingsheng.huang@outlook.com

†Contact author: xutong@jyu.fi

‡Contact author: whaolin@m.scnu.edu.cn

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International](https://creativecommons.org/licenses/by/4.0/) license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

light-quark EW dipole operators in inclusive deep inelastic scattering (DIS) with an unpolarized nucleon. Recently introduced in [38] as an extension of energy correlators from final-state jets to nucleon structure, NECs enable the imaging of partonic dynamics through correlations between initial-state partons and final-state energy flux in the target fragmentation region (TFR) [38–47] (see Ref. [48] for a review). By exploiting these correlations, we construct a chiral-odd quark NEC that accesses the quark transverse spin inside an unpolarized nucleon. This NEC ensures a nonzero interference between the dipole and SM amplitudes, generating unique azimuthal asymmetries in the energy distribution in DIS. We demonstrate that these asymmetries exhibit clean and strong sensitivity to the quark dipole couplings. Crucially, our method works without polarized nucleon beams and requires only inclusive calorimetric measurements, making it an exceptional observable for probing chirality-flipping interactions at both existing and future electron-ion colliders.

Quark transversity NEC—Let us begin by considering quark NECs in an unpolarized nucleon, which moves rapidly along the $+\hat{z}$ direction with momentum P . The associated NECs are encoded in the following correlation function [38,43]:

$$\mathcal{M}^{[\Gamma]}(x, \theta, \phi) = \int \frac{d\eta^-}{4\pi} e^{-ixP^+\eta^-} \times \langle P | \bar{\psi}(\eta^-) \mathcal{L}_n^\dagger(\eta^-) \Gamma \mathcal{E}(\theta, \phi) \mathcal{L}_n(0) \psi(0) | P \rangle, \quad (1)$$

which describes the conditional probability of finding a quark with the longitudinal momentum fraction x inside the target nucleon, given the observation of an energy flux at the solid angle (θ, ϕ) in the TFR. θ and ϕ denote the polar and azimuthal angles, respectively. Formally, the quark NECs can be viewed as the collinear PDFs modified by the insertion of the energy flow operator $\mathcal{E}(\theta, \phi)$, which is defined through its action on hadronic states [49,50]

$$\mathcal{E}(\theta, \phi) | X \rangle = \sum_{i \in X} \frac{E_i}{E_N} \delta(\theta_i^2 - \theta^2) \delta(\phi_i - \phi) | X \rangle, \quad (2)$$

where E_i denotes the energy of hadron i , normalized by the nucleon energy E_N . Here, we use the light-cone coordinates for a vector $a^\mu = (a^+, a^-, \vec{a}_\perp)$ with $a^\pm = (a^0 \pm a^3)/\sqrt{2}$, and $\mathcal{L}_n$ is the light-cone gauge link.

The correlation function $\mathcal{M}^{[\Gamma]}$ depends on the choice of the Dirac matrix Γ , which selects specific quark spin components. For instance, the leading-twist unpolarized quark NEC is obtained with $\Gamma = \gamma^+$:

$$f_1^q(x, \theta^2) = \mathcal{M}^{[\gamma^+]}(x, \theta, \phi), \quad (3)$$

which is chiral even and rotationally invariant in the transverse plane [38].

While prior studies have primarily focused on the chiral-even NECs [38–44] (see a systematic classification in [43]), we extend the formalism to the chiral-odd components. In particular, the quark transversity NEC can be obtained by using the projector $\Gamma^\alpha = i\sigma^{\alpha+}\gamma_5$ with $\alpha = 1, 2$, analogous to that used for transversity PDFs. The interpretation of quark transverse spin becomes explicit by rewriting $\Gamma^\alpha = \gamma^+(Q_+^\alpha - Q_-^\alpha)$, where $Q_\pm^\alpha = \frac{1}{2}(1 \mp \gamma^5 \gamma_\perp^\alpha)$ projects onto quarks with transverse spin aligned or antialigned with the $\alpha_\perp$ direction [51]. However, unlike transversity PDFs, which requires a transversely polarized nucleon [52], NECs can access the quark transverse spin even in an *unpolarized* nucleon, owing to the azimuthal dependence of the observed energy flux.

To explicitly illustrate this feature, we introduce a unit vector $\hat{n}_T^\mu = (0, 0, \cos \phi, \sin \phi)$ to specify the transverse direction of the observed energy flux. By rotational covariance and parity property, the quark transversity NEC $h_1^{t,q}(x, \theta^2)$ is parametrized as

$$\epsilon_\perp^{\alpha\rho} \hat{n}_{T,\rho} h_1^{t,q}(x, \theta^2) = \mathcal{M}^{[i\sigma^{\alpha+}\gamma_5]}(x, \theta, \phi), \quad (4)$$

where $\epsilon_\perp^{\mu\nu} = \epsilon^{\mu\nu-+}$ with $\epsilon^{0123} = 1$. This equation shows that $h_1^t(x, \theta^2)$ captures azimuthal correlations between the quark's transverse polarization and the energy flow direction. Physically, it describes the difference in probability of finding a quark polarized in the two opposite directions perpendicular to the transverse direction of the energy flow. Being chiral-odd, the transversity NEC is off diagonal in the helicity basis and naturally induces interference between amplitudes of opposite quark helicities. This provides a novel mechanism for generating chirality-flipping effects, enabling probes of dipole operators in unpolarized nucleon scatterings.

Probing the quark dipole operators with transversity NEC—To demonstrate how the transversity NEC probes the SMEFT dipole couplings, we consider the energy pattern within the TFR in inclusive DIS with an unpolarized nucleon, $l + p \rightarrow l' + X$. Experimentally, this observable is determined from the angular distribution of the total energy deposited in the calorimeters. Following Refs. [38,43], we define the energy pattern cross section as

$$\Sigma(\theta, \phi) = \sum_{i \in X} \int d\sigma^{l+p \rightarrow l'+X} \frac{E_i}{E_N} \delta(\theta^2 - \theta_i^2) \delta(\phi - \phi_i). \quad (5)$$

We work in the Breit frame, where the momentum transfer ($q = l' - l$) is aligned along the $-\hat{z}$ direction, and the nucleon moves in the opposite direction. As shown in Fig. 1, the polar angle θ is measured with respect to the nucleon beam, and the azimuthal angle ϕ is defined relative to the lepton scattering plane. The standard DIS kinematic variables are also used: the photon virtuality $Q^2 = -q^2$, the Bjorken variable $x_B = Q^2/(2P \cdot q)$, and the inelasticity

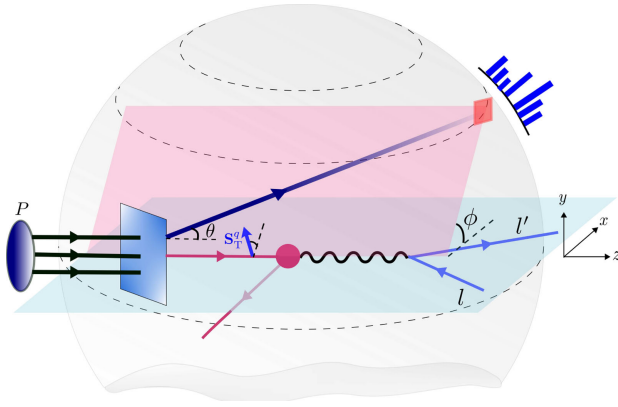


FIG. 1. Illustration of the NEC mechanism for probing the SMEFT quark dipole operator in DIS. The blue transverse plate highlights the role of the transversity NEC, which selects the quark transverse spin S_T^q in an unpolarized nucleon by measuring an energy flux from the target remnants. The magenta blob denotes the dipole interaction mediated by γ or Z . The interference with the SM amplitude is implied.

$y = Q^2/(sx_B)$ with $\sqrt{s}$ denoting the center-of-mass energy of the ep collision.

Let us first consider the case of an unpolarized electron beam. We are particularly interested in the following azimuthal modulations of the energy pattern:

$$\Sigma(\theta, \phi) = \Sigma_{UU}(\theta) + \Sigma_{UU}^{\sin\phi}(\theta) \sin\phi + \Sigma_{UU}^{\cos\phi}(\theta) \cos\phi. \quad (6)$$

We focus on the TFR, characterized by small polar angles ($\theta P^+ \ll Q$), where the energy flux originates from the fragmentation of target remnants. Following the general arguments in [39,43], the energy pattern cross section in this region factorizes into hard partonic scattering coefficients and the nonperturbative NECs, defined by Eq. (1). Because NECs encapsulate all the information regarding the energy flow, including its azimuthal ϕ dependence, the ϕ -independent cross section Σ_{UU} factorizes with unpolarized NECs, such as f_1 given in Eq. (3). By contrast, the ϕ -dependent terms $\Sigma_{UU}^{\sin\phi}$ and $\Sigma_{UU}^{\cos\phi}$ are expected to factorize with the transversity NEC h_1^T , which encodes azimuthal correlations as seen in Eq. (4). Remarkably, after factorization, the hard scatterings remain the same as in standard inclusive DIS process [38,39,43].

Within the SM, the energy pattern is isotropic in the transverse plane. Since quark helicity flip in DIS hard scattering is suppressed by light quark masses, only the chiral-even unpolarized quark NEC $f_1(x, \theta^2)$ contributes to $\Sigma(\theta, \phi)$ in the massless limit, while the chiral-odd contributions from the transversity NEC $h_1^T(x, \theta^2)$ vanish to all orders in α_s . The only nonzero contribution arises from the azimuthally independent cross section Σ_{UU} . At leading order in α_s , the contribution from the photon exchanges takes the following form [38]:

$$\frac{d\Sigma_{UU}}{dx_B dQ^2} = \frac{2\pi\alpha_{\text{em}}^2}{Q^4} (y^2 - 2y + 2) \sum_q Q_q^2 f_1^q(x_B, \theta^2), \quad (7)$$

where $\alpha_{\text{em}} = e^2/(4\pi)$ is the fine structure constant and Q_q is the quark electric charge.

Although the $\sin\phi$ and $\cos\phi$ azimuthal modulations are absent within the SM, they could arise from chirality-flipping NP effects, such as the following dimension-6 SMEFT dipole interactions [3]:

$$\begin{aligned} \mathcal{L} = & \frac{1}{\Lambda^2} [C_{uW}(\bar{Q}\sigma^{\mu\nu}u)\tau^I\tilde{H}W_{\mu\nu}^I + C_{uB}(\bar{Q}\sigma^{\mu\nu}u)\tilde{H}B_{\mu\nu} \\ & + C_{dW}(\bar{Q}\sigma^{\mu\nu}d)\tau^I\tilde{H}W_{\mu\nu}^I + C_{dB}(\bar{Q}\sigma^{\mu\nu}d)\tilde{H}B_{\mu\nu}] + \text{H.c.}, \end{aligned} \quad (8)$$

where τ^I are the Pauli matrices, and $\tilde{H}_I = \epsilon_{IJ}H_J^*$. The field strength tensors $W_{\mu\nu}^I$ and $B_{\mu\nu}$ correspond to the $SU(2)_L$ and $U(1)_Y$ gauge fields, respectively. Here, Q denotes the left-handed quark doublet, while u and d represent the right-handed up- and down-type quark singlets. We restrict our attention to operators involving only the first generation and define their dipole couplings to photon and Z -boson c_i as

$$\begin{aligned} c_{q\gamma} &= (v/\sqrt{2}\Lambda^2)(c_W C_{qB} \pm s_W C_{qW}), \\ c_{qZ} &= (v/\sqrt{2}\Lambda^2)(-s_W C_{qB} \pm c_W C_{qW}). \end{aligned} \quad (9)$$

The plus (minus) sign is for the up (down) quark, $c_W = \cos\theta_W$ and $s_W = \sin\theta_W$ with θ_W being the Weinberg angle, and v is the Higgs vacuum expectation value.

As illustrated in Fig. 1, these azimuthal modulations originate from the interplay between the SMEFT dipole operator and the quark transversity NEC $h_1^T(x, \theta^2)$. Being inherently chiral odd, the dipole interaction induces a quark helicity flip during the hard scattering with the virtual photon or Z boson. As a result, the dominant contributions emerge from the interference between the helicity-conserving SM amplitude and the helicity-flipping SMEFT amplitude at $\mathcal{O}(\Lambda^{-2})$. Such interference necessarily couples to a chiral-odd quark NEC, namely, the transversity NEC. As revealed in Eq. (4), the quark transverse spin is accessed through intrinsic azimuthal correlations with the energy flux direction $\vec{n}_T$, thereby yielding distinct $\sin\phi$ and $\cos\phi$ modulations. Furthermore, these asymmetries are uniquely generated by the dipole operators, as other dimension-6 SMEFT operators either require transverse electron polarization or are suppressed by the electron Yukawa coupling.

We find that the $\sin\phi$ modulation is linearly determined by the real parts of the quark dipole couplings $c_{q\gamma}$ and c_{qZ} through the interference between photon and Z -boson exchanges. The corresponding cross section is given by

$$\frac{d\Sigma_{UU}^{\sin\phi}}{dx_B dQ^2} = \frac{4\pi\alpha_{\text{em}}^2}{ec_W s_W} \frac{y\sqrt{1-y}}{Q(Q^2 + m_Z^2)} \sum_q h_1^{t,q}(x_B, \theta^2) \times \left\{ \left[\frac{2-y}{y} g_A^q g_V^e + g_V^q g_A^e \right] \frac{\text{Re}[c_{q\gamma}]}{c_W s_W} - Q_q g_A^e \text{Re}[c_{qZ}] \right\}, \quad (10)$$

where $g_V^f = T_f^3/2 - Q_f s_W^2$ and $g_A^f = T_f^3/2$ are the vector and axial-vector couplings with T^3 as the third component of weak isospin. Similarly, the $\cos\phi$ modulation is sensitive to the imaginary parts of the dipole couplings and can be obtained from Eq. (10) by replacing $\text{Re}[c_{q\gamma,Z}]$ with $-\text{Im}[c_{q\gamma,Z}]$. The antiquark components in Eqs. (7) and (10) can be obtained by taking $g_A^q \rightarrow -g_A^q$. Considering the EIC kinematics, we have neglected pure Z-boson exchanges.

Our analysis can be readily generalized to include a longitudinally polarized electron beam, yielding the cross section $\Sigma_{LU}^{\sin\phi(\cos\phi)}$ factorized in terms of the transversity NEC h_1^t . Owing to the parity violation induced by electron polarization, the azimuthal modulations can be generated via $\gamma\gamma$ interference, without Q^2/m_Z^2 suppression present in the unpolarized case $\Sigma_{UU}^{\sin\phi(\cos\phi)}$. As indicated later, this leads to an enhanced sensitivity to the dipole couplings. Complete analytical results are provided in Supplemental Material [53].

Calorimetric asymmetries and the nonperturbative input—To quantify the effects of the SMEFT dipole operators, we introduce azimuthal angle asymmetries of the energy pattern. For an unpolarized electron beam, these are purely calorimetric asymmetries defined with respect to the lepton scattering plane,

$$A_{UU}^u = \frac{\pi \Sigma(u > 0) - \Sigma(u < 0)}{2 \Sigma(u > 0) + \Sigma(u < 0)}, \quad (11)$$

where u is taken to be either $\sin\phi$ or $\cos\phi$, isolating the real or imaginary part of the dipole couplings, respectively. $\Sigma(u > 0)$ represents the energy pattern cross section integrated over the region with $u > 0$:

$$\Sigma(u > 0) \equiv \int_{\theta_{\min}^2}^{\theta_{\max}^2} d\theta^2 w(\theta) \int_0^{2\pi} d\phi \Sigma(\theta, \phi) \Theta(u), \quad (12)$$

and similarly for $u < 0$. The polar-angle weight function $w(\theta)$ is introduced for later use. For a longitudinally polarized electron beam, one can analogously define the asymmetries A_{LU}^u by combining the calorimetric asymmetry with the longitudinal spin asymmetry, as we specify in [53].

Numerical evaluation of these asymmetries requires nonperturbative input for the quark NECs f_1 and h_1^t . While the unpolarized quark NEC f_1 can be extracted directly from Σ_{UU} in the TFR [38], one can also measure

the transversity NEC h_1^t without assuming SMEFT interactions by tagging an additional hadron in the current fragmentation region. Such measurements are feasible at facilities including HERA, JLab, and the EIC.

Although direct experimental determinations of NECs are not yet available, a remarkable correspondence has been established between θ -weighted moments of the NECs and $k_\perp$ -weighted moments of transverse-momentum-dependent parton distributions (TMDs), including the unpolarized quark NEC [42]. This correspondence can be straightforwardly extended to the transversity NEC. Specifically, we have

$$E_N \int d\theta^2 |\sin\theta| f_1(x, \theta^2) = \int \frac{d^2 k_\perp}{2\pi} |k_\perp| \mathcal{F}_1(x, k_\perp^2),$$

$$E_N \int d\theta^2 |\sin\theta| h_1^t(x, \theta^2) = \int \frac{d^2 k_\perp}{2\pi} \frac{k_\perp^2}{M} \mathcal{H}_1^t(x, k_\perp^2), \quad (13)$$

where $\mathcal{F}_1$ and $\mathcal{H}_1^t$ are the unpolarized and Boer-Mulder quark TMDs [63], respectively. These relations allow us to infer the NECs from existing global fits of TMDs. In our analysis, we adopt the parametrization from [64], which assumes a Gaussian form for the $k_\perp$ dependence. The evolution effects are accounted for as described in [53].

To utilize the above relations, in Eq. (12) we choose the θ weight as $w(\theta) = |\sin\theta|$. The limits of the θ integration are determined by both theoretical and experimental considerations. The upper bound $\theta_{\max}$ is imposed to ensure the validity of the factorization in the TFR [39,43]. This corresponds to a transverse momentum cutoff $k_\perp^{\max} \sim Q\theta_{\max}$ in the $k_\perp$ -integrals in Eq. (13). Due to rapid falloff of the Gaussian TMDs at large $k_\perp$, this cutoff has negligible numerical impact on the integral. The lower bound $\theta_{\min}$ is determined by the coverage of the calorimeter in the nucleon forward region. It imposes a minimal transverse-momentum threshold given by $k_\perp^{\min} \approx E_{\min}\theta_{\min}$, where $E_{\min}$ characterizes the energy resolution of calorimeters [65]. Practically, we adopt $k_\perp^{\min} \approx 0$ GeV and $k_\perp^{\max} = 5$ GeV. With these inputs, we have evaluated numerically both $A_{UU}^{\sin\phi}$ and $A_{LU}^{\sin\phi}$ for the EIC kinematics, and the detailed results can be found in Supplemental Material [53].

Projected sensitivity at the EIC—To estimate the sensitivity to the dipole operators at the EIC, we perform a χ^2 analysis based on the predicted asymmetries. We select a representative EIC configuration [65] with a center-of-mass energy of $\sqrt{s} = 105$ GeV, which corresponds to the optimal integrated luminosity of $\mathcal{L} = 100 \text{ fb}^{-1}$. The analysis is carried out in Q bins for $Q \in [10, 60]$ GeV and $x \in [0.01, 0.5]$, incorporating an inelasticity cut $0.1 \leq y \leq 0.9$. We assume the experimental value consistent with the leading-twist SM prediction, and it is negligible. With this setup, the statistical uncertainty $\delta A_i \simeq \pi/(2\sqrt{\mathcal{E}_{\text{total}}})$ is dominant over the systematic

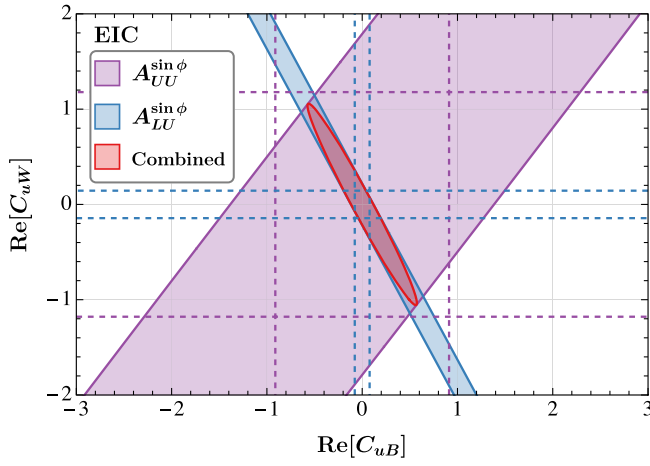


FIG. 2. Projected 68% C.L. constraints at the EIC on the two Wilson coefficients $\text{Re}[C_{uB}]$ and $\text{Re}[C_{uW}]$, assuming $\Lambda = 1$ TeV. The limits are derived from the azimuthal asymmetries $A_{UU}^{\sin\phi}$ (purple), $A_{LU}^{\sin\phi}$ (blue), and their combination (red). The shaded regions represent the simultaneous two-coefficient constraints, while dashed lines show the limits for a single coefficient.

uncertainties. Here, $\mathcal{E}_{\text{total}} = (\Sigma(u > 0) + \Sigma(u < 0)) \times \mathcal{L}$ is the accumulated energy deposition.

Figure 2 presents the projected constraints on the real parts of dimensionless Wilson coefficients C_{uW} and C_{uB} at 68% C.L. using $\sin\phi$ asymmetry. Assuming single-operator dominance, $A_{UU}^{\sin\phi}$ constrain the Wilson coefficients almost at the $\mathcal{O}(1)$ level, as shown by the purple dashed lines. However, when both coefficients are considered simultaneously, the resulting constraint ellipse is elongated into a purple band, indicating a strong correlation between C_{uW} and C_{uB} .

We observe that this correlation is significantly alleviated by incorporating the $A_{LU}^{\sin\phi}$ asymmetry. For the EIC, we assume 70% polarization of the electron beam [65]. Under this setup, single-operator constraints can reach the $\mathcal{O}(0.01) \sim \mathcal{O}(0.1)$ level. Importantly, while $A_{LU}^{\sin\phi}$ retains strong correlations between coefficients, the orientation of the corresponding blue band is different from that of $A_{UU}^{\sin\phi}$. Their combination thus greatly improves the parameter space resolution, confining the allowed region to a narrow area highlighted in red. Similar improvements are also observed in the constraints on C_{dW} - C_{dB} and $c_{q\gamma}$ - c_{qZ} (see Supplemental Material [53]). The constraints on the imaginary parts can be obtained analogously from the $\cos\phi$ asymmetry, yielding identical sensitivities. The constraints are stable against different nonperturbative inputs of NEC, as we analyze in [53]. These results underscore the strong capability of the proposed asymmetries to unveil light-quark dipole interactions at the EIC.

Besides the future EIC, the absence of a polarized nucleon beam requirement for these asymmetries also opens up opportunities to use available data from HERA. In Fig. 3, we compare the projected sensitivity

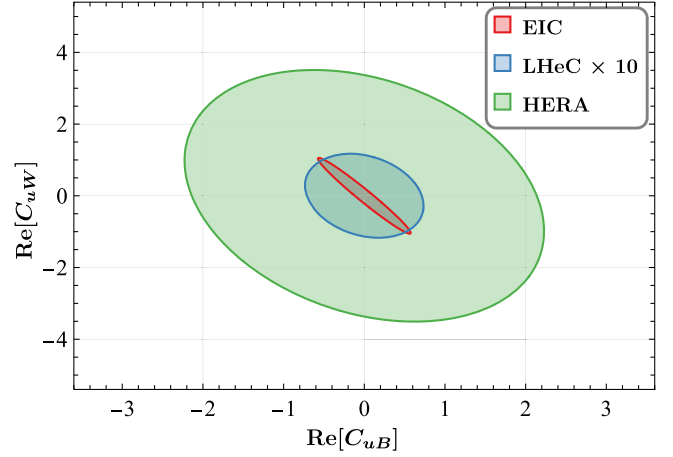


FIG. 3. Comparison of projected 68% C.L. constraints on $\text{Re}[C_{uB}]$ and $\text{Re}[C_{uW}]$ from HERA (green), the EIC (red), and the LHeC (blue), assuming $\Lambda = 1$ TeV. The LHeC limits are scaled by a factor of 10 for visibility. A detailed description of our analysis setup is provided in [53].

from the combined-asymmetry analysis among the EIC, the HERA [66], and the LHeC [67]. Although HERA benefits from higher Q^2 coverage, its lower luminosity results in slightly weaker sensitivity compared to the EIC. However, it provides better discrimination between operators. With even higher energy and better luminosity, the LHeC could outperform both HERA and EIC by an order of magnitude.

Finally, let us comment on the photon-quark dipole coupling $c_{q\gamma}$ [given in Eq. (9)], which can be constrained by measuring hadronic magnetic and electric dipole moments at low energy. While high-precision neutron EDM measurements yield strong bounds on $\text{Im}[c_{q\gamma}]$ [6,68], existing hadronic AMM measurements do not well constrain $\text{Re}[c_{q\gamma}]$, and thus $\text{Re}[C_{qW}]$ and $\text{Re}[C_{qB}]$. These real parts encode the P - and T -conserving NP effects and may produce sizable signatures without conflicting with neutron-EDM bounds. As shown in Fig. 3, our proposal provides sensitive limits on the real components; for explicit bounds on $\text{Re}[c_{q\gamma}]$, we refer the reader to Supplemental Material [53]. Consequently, this Letter is dedicated to searching for these P - and T -conserving NP effects at ep colliders.

Conclusion and discussion—In summary, we have shown that nucleon energy correlators provide a powerful framework to probe EW light-quark dipole operators in inclusive DIS with an unpolarized nucleon. By exploiting azimuthal correlations of the energy flux in the TFR, we introduce a chiral-odd NEC sensitive to the transverse spin of quarks, even within an unpolarized target. Coupling to the interference between the dipole and SM amplitudes, the transversity NEC produces characteristic $\sin\phi$ and $\cos\phi$ asymmetries in the energy pattern, providing clean access to the dipole couplings. Our combined analysis with both unpolarized and longitudinally polarized electron beams

demonstrates that these observables can deliver precise and complementary constraints at the EIC.

This Letter initiates an exploration of energy-correlator observables in the search for new physics (for an effort in other directions, see Ref. [69]). Compared with proposals based on transversity PDFs or dihadron production [17,19], our NEC-based approach eliminates the need for nucleon polarization and relies entirely on inclusive calorimetric measurements, without particle identification or hadron reconstruction. These features simplify experimental implementation and reduce statistical uncertainties. They also enable the full utilization of high luminosity from the future EIC and LHeC as well as the use of existing HERA data. Furthermore, the resulting constraints may benefit from extensions to heavy-ion beams.

While the present Letter employs inclusive NECs, improved angular resolution could be achieved through track-based measurements of charged-hadron energy flux [70–73], which would be valuable for precision studies on the dipole couplings. Furthermore, the chiral-odd NEC framework can be extended to polarized nucleons, with all such NECs sharing identical hard-scattering kernels, thus offering robust and complementary constraints. Beyond DIS, NECs could also be applied to processes such as Drell-Yan production at the LHC. Detailed investigations along these lines are left for future work.

Note added—On the day this Letter was submitted to arXiv, a preprint [74] appeared in which the author proposed using a chiral-odd fracture function to probe Yukawa couplings at the LHC. In contrast to NECs, fracture functions typically involve measuring a specific type of hadron with momentum $\vec{P}_h$ in the TFR (see, e.g., Refs. [75–79]), which requires particle identification and tracking. While [74] focused on the $\vec{P}_h$ -integrated fracture functions, we note that the $\vec{P}_h$ -dependent fracture functions (also referred to as extended fracture functions [80]) are formally related to NECs through an inclusive energy sum rule recently established in [43].

Acknowledgments—We thank Yu Jia, Yi Liao, Xiaohui Liu, Jianping Ma, and Hongxi Xing for helpful discussions and comments. X.-B. T. is supported by the Research Council of Finland, the Centre of Excellence in Quark Matter and Projects No. 338263 and No. 359902, and by the European Research Council (ERC, Grant Agreements No. ERC-2023-101123801 GlueSatLight and No. ERC-2018-ADG-835105 YoctoLHC). H.-L. W. is supported by Grant No. NSFC-12035008.

The content of this article does not reflect the official opinion of the European Union and responsibility for the information and views expressed therein lies entirely with the authors.

Data availability—The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

- [1] W. Buchmuller and D. Wyler, *Nucl. Phys.* **B268**, 621 (1986).
- [2] C. Arzt, M. B. Einhorn, and J. Wudka, *Nucl. Phys.* **B433**, 41 (1995).
- [3] B. Grzadkowski, M. Iskrzynski, M. Misiak, and J. Rosiek, *J. High Energy Phys.* **10** (2010) 085.
- [4] See, e.g., the quark dipole operator $(\bar{Q}\sigma^{\mu\nu}d)HB_{\mu\nu}$ in Eq. (8).
- [5] R. Alonso, E. E. Jenkins, A. V. Manohar, and M. Trott, *J. High Energy Phys.* **04** (2014) 159.
- [6] J. Kley, T. Theil, E. Venturini, and A. Weiler, *Eur. Phys. J. C* **82**, 926 (2022).
- [7] J. Aebischer, W. Dekens, E. E. Jenkins, A. V. Manohar, D. Sengupta, and P. Stoffer, *J. High Energy Phys.* **07** (2021) 107.
- [8] M. Ardu and N. Valori, *Phys. Rev. D* **113**, 015035 (2026).
- [9] R. Boughezal, E. Mereghetti, and F. Petriello, *Phys. Rev. D* **104**, 095022 (2021).
- [10] L. Allwicher, D. A. Faroughy, F. Jaffredo, O. Sumensari, and F. Wilsch, *J. High Energy Phys.* **03** (2023) 064.
- [11] E. da Silva Almeida, N. Rosa-Agostinho, O. J. P. Éboli, and M. C. Gonzalez-Garcia, *Phys. Rev. D* **100**, 013003 (2019).
- [12] X. Li, B. Yan, and C. P. Yuan, *Phys. Rev. D* **111**, 073007 (2025).
- [13] R. Gauld, U. Haisch, and J. Weiss, *SciPost Phys.* **18**, 148 (2025).
- [14] G. Li, X. Li, and B. Yan, *Phys. Lett. B* **870**, 139931 (2025).
- [15] R. Escribano and E. Masso, *Nucl. Phys.* **B429**, 19 (1994).
- [16] G. Kopp, D. Schaile, M. Spira, and P. M. Zerwas, *Z. Phys. C* **65**, 545 (1995).
- [17] R. Boughezal, D. de Florian, F. Petriello, and W. Vogelsang, *Phys. Rev. D* **107**, 075028 (2023).
- [18] X.-K. Wen, B. Yan, Z. Yu, and C. P. Yuan, *Phys. Rev. Lett.* **131**, 241801 (2023).
- [19] X.-K. Wen, B. Yan, Z. Yu, and C. P. Yuan, *arXiv:2408.07255*.
- [20] X.-K. Wen, B. Yan, Z. Yu, and C. P. Yuan, *Phys. Rev. D* **112**, 053004 (2025).
- [21] M. B. Hecht, C. D. Roberts, and S. M. Schmidt, *Phys. Rev. C* **64**, 025204 (2001).
- [22] M. Pitschmann, C.-Y. Seng, C. D. Roberts, and S. M. Schmidt, *Phys. Rev. D* **91**, 074004 (2015).
- [23] T. Liu, Z. Zhao, and H. Gao, *Phys. Rev. D* **97**, 074018 (2018).
- [24] T. Fukuyama, *Int. J. Mod. Phys. A* **27**, 1230015 (2012).
- [25] T. Bhattacharya, V. Cirigliano, and R. Gupta, *Proc. Sci. LATTICE2012* (2012) 179 [arXiv:1212.4918].
- [26] T. Chupp, P. Fierlinger, M. Ramsey-Musolf, and J. Singh, *Rev. Mod. Phys.* **91**, 015001 (2019).
- [27] R. Alarcon *et al.*, in Snowmass 2021 (2022), *arXiv:2203.08103*.
- [28] A. Keshavarzi, K. S. Khaw, and T. Yoshioka, *Nucl. Phys.* **B975**, 115675 (2022).
- [29] G. A. González-Sprinberg, in 17th International Workshop on Tau Lepton Physics (2024), *arXiv:2406.15286*.

- [30] B. Abi *et al.*, *Phys. Rev. Lett.* **126**, 141801 (2021).
- [31] D. P. Aguillard *et al.* (Muon g-2 Collaboration), *Phys. Rev. Lett.* **135**, 101802 (2025).
- [32] R. Aliberti *et al.*, *Phys. Rep.* **1143**, 1 (2025).
- [33] M. Ablikim *et al.* (BESIII Collaboration), *arXiv:2506.19180*.
- [34] ACME Collaboration, *Nature (London)* **562**, 355 (2018).
- [35] X. Fan, T. G. Myers, B. A. D. Sukra, and G. Gabrielse, *Phys. Rev. Lett.* **130**, 071801 (2023).
- [36] C. Abel *et al.*, *Phys. Rev. Lett.* **124**, 081803 (2020).
- [37] H.-L. Wang, X.-K. Wen, H. Xing, and B. Yan, *Phys. Rev. D* **109**, 095025 (2024).
- [38] X. Liu and H. X. Zhu, *Phys. Rev. Lett.* **130**, 091901 (2023).
- [39] H. Cao, X. Liu, and H. X. Zhu, *Phys. Rev. D* **107**, 114008 (2023).
- [40] H. Cao, H. T. Li, and Z. Mi, *Phys. Rev. D* **109**, 096004 (2024).
- [41] H.-Y. Liu, X. Liu, J.-C. Pan, F. Yuan, and H. X. Zhu, *Phys. Rev. Lett.* **130**, 181901 (2023).
- [42] X. Liu and H. X. Zhu, *arXiv:2403.08874*.
- [43] K.-B. Chen, J.-P. Ma, and X.-B. Tong, *J. High Energy Phys.* **08** (2024) 227.
- [44] H. Mäntysaari, Y. Tawabutr, and X.-B. Tong, *Phys. Rev. D* **112**, 114027 (2025).
- [45] X. L. Li, X. Liu, F. Yuan, and H. X. Zhu, *Phys. Rev. D* **108**, L091502 (2023).
- [46] Y. Guo, X. Liu, F. Yuan, and H. X. Zhu, *Research* **2025**, 0552 (2025).
- [47] Y. Guo, X. Liu, and F. Yuan, *Phys. Rev. D* **112**, 054006 (2025).
- [48] I. Moulton and H. X. Zhu, *arXiv:2506.09119*.
- [49] N. A. Sveshnikov and F. V. Tkachov, *Phys. Lett. B* **382**, 403 (1996).
- [50] C. W. Bauer, S. P. Fleming, C. Lee, and G. F. Sterman, *Phys. Rev. D* **78**, 034027 (2008).
- [51] R. L. Jaffe, in *Ettore Majorana International School of Nucleon Structure: 1st Course: The Spin Structure of the Nucleon* (1996), pp. 42–129, *arXiv:hep-ph/9602236*.
- [52] R. L. Jaffe and X.-D. Ji, *Nucl. Phys.* **B375**, 527 (1992).
- [53] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/7y5n-cscx> for additional analytical results, evolution details, numerical setup, uncertainty analysis, and comparisons with other probes, which includes Refs. [54–62].
- [54] J. C. Collins, *Phys. Rev. D* **57**, 3051 (1998); **61**, 019902(E) (2000).
- [55] Z.-B. Kang and J.-W. Qiu, *Phys. Lett. B* **713**, 273 (2012).
- [56] X. Artru and M. Mekhfi, *Z. Phys. C* **45**, 669 (1990).
- [57] W. Vogelsang, *Phys. Rev. D* **57**, 1886 (1998).
- [58] Z. Lu and I. Schmidt, *Phys. Rev. D* **81**, 034023 (2010).
- [59] H. Abramowicz *et al.* (H1, ZEUS Collaborations), *Eur. Phys. J. C* **75**, 580 (2015).
- [60] C. Diaconu, T. Haas, M. Medinnis, K. Rith, and A. Wagner, *Annu. Rev. Nucl. Part. Sci.* **60**, 101 (2010).
- [61] E. Christova, D. Kotlorz, and E. Leader, *Phys. Rev. D* **102**, 014035 (2020).
- [62] B. Zhang, Z. Lu, B.-Q. Ma, and I. Schmidt, *Phys. Rev. D* **78**, 034035 (2008).
- [63] R. Boussarie *et al.*, *arXiv:2304.03302*.
- [64] V. Barone, S. Melis, and A. Prokudin, *Phys. Rev. D* **81**, 114026 (2010).
- [65] R. Abdul Khalek *et al.*, *Nucl. Phys.* **A1026**, 122447 (2022).
- [66] M. Klein and R. Yoshida, *Prog. Part. Nucl. Phys.* **61**, 343 (2008).
- [67] P. Agostini *et al.* (LHeC, FCC-he Study Group Collaborations), *J. Phys. G* **48**, 110501 (2021).
- [68] J. Kumar and E. Mereghetti, *J. High Energy Phys.* **09** (2024) 028.
- [69] L. Ricci and M. Riemann, *Phys. Rev. D* **106**, 114010 (2022).
- [70] H. Chen, I. Moulton, X. Zhang, and H. X. Zhu, *Phys. Rev. D* **102**, 054012 (2020).
- [71] Y. Li, I. Moulton, S. S. van Velzen, W. J. Waalewijn, and H. X. Zhu, *Phys. Rev. Lett.* **128**, 182001 (2022).
- [72] M. Jaarsma, Y. Li, I. Moulton, W. J. Waalewijn, and H. X. Zhu, *J. High Energy Phys.* **12** (2023) 087.
- [73] K. Lee and I. Moulton, *arXiv:2308.00746*.
- [74] J. K. L. Michel, *arXiv:2508.05914*.
- [75] M. Anselmino, V. Barone, and A. Kotzinian, *Phys. Lett. B* **699**, 108 (2011).
- [76] K. B. Chen, J. P. Ma, and X. B. Tong, *Phys. Rev. D* **108**, 094015 (2023).
- [77] K.-B. Chen, J.-P. Ma, and X.-B. Tong, *J. High Energy Phys.* **05** (2024) 298.
- [78] K. B. Chen, J. P. Ma, and X. B. Tong, *J. High Energy Phys.* **11** (2021) 038.
- [79] X. P. Chai, K. B. Chen, J. P. Ma, and X. B. Tong, *J. High Energy Phys.* **10** (2019) 285.
- [80] M. Grazzini, L. Trentadue, and G. Veneziano, *Nucl. Phys.* **B519**, 394 (1998).