

Nucleon Tomography with Zero Jettiness

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We propose a novel strategy to systematically isolate the nucleon's intrinsic nonperturbative three-dimensional structure by employing zero jettiness to suppress initial-state radiation in transverse-momentum-dependent observables. Applying this method to transverse single spin asymmetries (SSAs) in $W^\pm$ and Z^0 boson production at Relativistic Heavy Ion Collider (RHIC), we demonstrate a substantial enhancement of the asymmetry signal (e.g., by 115% for W^- SSA at $q_\perp = 5$ GeV). We show that this enhancement yields a substantial net gain in experimental sensitivity—even after accounting for the statistical cost of the veto—facilitating a more definitive test of the predicted Sivers function sign change. We further explore its applicability to spin-dependent measurements at the Electron-Ion Collider. Our analysis is formulated within a joint resummation framework that systematically resums large logarithms associated with both the veto scale and the gauge boson's transverse momentum.

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Introduction—Understanding the three-dimensional (3D) structure of the nucleon is of fundamental importance to hadron physics, particularly with the upcoming Electron-Ion Colliders (EICs) in the United States [1,2] and China [3]. One of the central objectives of these future facilities is to map out the 3D motion of partons inside the nucleon through ep and eA collisions, as characterized by transverse-momentum-dependent (TMD) parton distributions [4–8]. TMD distributions not only provide a direct 3D imaging of the nucleon [9,10] but also offer insights into the nucleon's spin structure [9–14] and shed light on fundamental aspects of QCD factorization theorem [15–19]. They also encode crucial information on the color glass condensate—a unique form of nuclear matter at small x [20–29].

A significant challenge in probing the intrinsic transverse motion of partons is the contamination from initial state radiation (ISR), whose recoil effects introduce

perturbative contributions that obscure and complicate the disentanglement of the nonperturbative structure of TMDs [8]. To address this, we propose in this Letter a novel strategy: employing zero jettiness as a veto to suppress central ISR. This approach enhances experimental sensitivity to the inherent nonperturbative TMD dynamics, opening new avenues to separate these from perturbative contamination, a persistent difficulty that has limited precision extractions at collider energies. Notably, while semi-inclusive deep inelastic scattering (SIDIS) processes allow for experimental tuning of the hard scale Q , electroweak Drell-Yan processes have Q fixed by the gauge boson mass. Here, zero jettiness emerges as an effective tunable energy scale, offering novel flexibility in controlling TMD evolution effects.

The zero-jettiness observable, originally introduced as a global event shape to veto additional jets in an inclusive manner [30], has become a powerful tool in collider phenomenology [31–34]. This observable enhances signal sensitivity by effectively suppressing additional radiation. Recent advances in resummation techniques have enabled precise and systematically improvable predictions for zero-jettiness cross sections [35–41], thereby playing an essential role in interpreting LHC data.

Despite these advances, the application of the zero-jettiness veto to polarization-dependent TMD observables remains unexplored. Here, we present the first-ever extension of joint resummation [37,42–45] to polarization-

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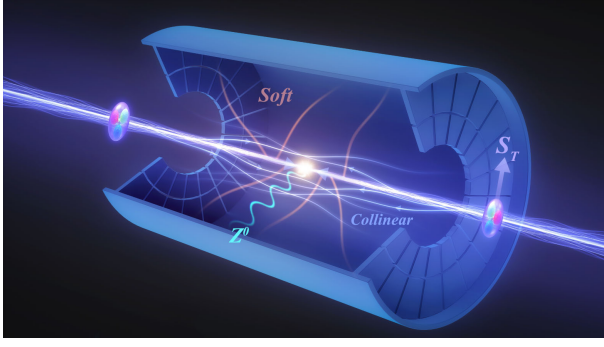


FIG. 1. Schematic illustration of the final state configuration for Z^0 production in polarized pp collisions with zero jettiness. The faint red lines represent the soft radiation vetoed by the cut, and the bright straight lines represent the collinear hard radiation that drives the scale evolution of PDFs and give rise to SSAs via the Qiu-Sterman mechanism.

dependent observables in Drell-Yan-like processes shown in Fig. 1, combining veto-induced logarithms and k_t logarithms. More specifically, we investigate the impact of the zero-jettiness veto on transverse SSAs in $W^\pm$ and Z^0 boson production at RHIC [46,47], which provides a critical test of the predicted sign change of the Sivers function [17]. Existing SSA measurements are strongly diluted by TMD evolution effects at high energies [48,49]. We demonstrate that imposing a central zero-jettiness veto significantly enhances the SSA signal generated by the Qiu-Sterman mechanism [50,51]. Crucially, we show that this enhancement provides a substantial net gain in experimental sensitivity, even after accounting for the statistical cost of the veto. Moreover, while our results focus on transverse spin asymmetries in Drell-Yan processes, the formalism is broadly applicable. This approach opens a new strategy to enhance sensitivity to intrinsic parton motion across a wide range of TMD observables, with relevance for RHIC, LHC, the future EIC, and beyond.

TMDs with zero jettiness—As TMD evolution is predominantly driven by soft gluon radiation, observables that suppress radiation in the central rapidity region are particularly effective for enhancing sensitivity to nonperturbative effects. Among the various event shape observables, beam thrust, also referred to as zero jettiness [30,52–56], stands out as a powerful tool for this purpose. For electroweak Drell-Yan processes, the zero jettiness variable is defined as

$$\tau \equiv \frac{2}{Q^2} \sum_i \min\{p_a \cdot l_i, p_b \cdot l_i\} = \sum_i \frac{|\vec{l}_{\perp,i}|}{Q} e^{-|y_i - y|}, \quad (1)$$

where Q and y are the mass and rapidity of the produced gauge boson, respectively. The sum runs over all particles i (excluding the gauge boson) with momentum l_i , transverse momentum $\vec{l}_{\perp,i}$ and rapidity y_i in the final state. The momenta of the incoming partons are defined as $p_a^\mu =$

$Qe^{+y}n_a^\mu/2$ and $p_b^\mu = Qe^{-y}n_b^\mu/2$, where $n_{a,b}^\mu = (1, 0, 0, \pm 1)$ are lightlike reference vectors. This observable is designed so that energetic particles collimated along the beam directions contribute minimally to τ , while central emissions contribute significantly. Imposing a veto $\tau < \tau_0$ therefore strongly suppresses central gluon radiation and effectively constrains ISR effects.

We examine the transverse momentum spectrum of the gauge boson across a wider range of scales, where the perturbative expansion is spoiled by two distinct types of large logarithms: $\ln(Q^2/q_\perp^2)$ and $\ln(1/\tau_0)$, thereby going beyond the standard TMD resummation framework, which accounts only for the former. While previous studies have explored the joint resummation of TMDs with zero jettiness or jet veto [37,42–45], we present a novel derivation that is both transparent and heuristic, aiming to enhance accessibility to a broader audience. Our formulation unifies the various soft-collinear effective theory (SCET) [57–61] regimes (SCET_I, SCET_{II}, and SCET₊) and is applicable in a more comprehensive kinematic region.

Sudakov form factor in TMD resummation—To establish the joint resummation framework, we first recall the standard TMD factorization formula for the unpolarized Drell-Yan cross section [4–7],

$$\frac{d\sigma_{UU}}{dyd^2\vec{q}_\perp} = \sigma_0 \sum_{q,q'} |V_{qq'}|^2 \int \frac{d^2\vec{b}_\perp}{(2\pi)^2} e^{i\vec{q}_\perp \cdot \vec{b}_\perp} H_{qq'}(Q, \mu) \times f_{q/p}(x_q, b, \mu, \zeta_q) f_{q'/p}(x_{q'}, b, \mu, \zeta_{q'}), \quad (2)$$

where $b \equiv |\vec{b}_\perp|$, σ_0 is the Born cross section, $|V_{qq'}|^2$ is the Cabibbo-Kobayashi-Maskawa (CKM) matrix element, H is the hard factor, and $f_{i/p}(x_i, b, \mu, \zeta_i)$ is the unpolarized TMD parton distribution function (PDF) depending on scales μ and ζ . Parton fractions are $x_q = Qe^y/\sqrt{s}$ and $x_{q'} = Qe^{-y}/\sqrt{s}$.

At next-to-leading logarithmic (NLL) in the Collins-Soper-Sterman framework [7,62], for $q_\perp \gg \Lambda_{\text{QCD}}$, the TMD PDF is related to the collinear PDF $f_{q/p}(x, \mu_b)$ at $\mu_b = 2e^{-\gamma_E}/b$:

$$f_{q/p}(x, b, \mu, \zeta) = e^{-S_p^{\text{TMD}}(b, \mu, \zeta)/2} f_{q/p}(x, \mu_b). \quad (3)$$

Substituting into Eq. (2) with $\mu = Q$ and $\zeta_q = \zeta_{q'} = Q^2$ and using azimuthal integration, we obtain

$$\frac{d\sigma_{UU}}{dyd^2\vec{q}_\perp} = \sigma_0 \sum_{q,q'} |V_{qq'}|^2 \int_0^\infty \frac{bdb}{2\pi} J_0(bq_\perp) \times f_q(x_q, \mu_b) f_{q'}(x_{q'}, \mu_b) e^{-S_p^{\text{TMD}}(b)}, \quad (4)$$

where J_0 is the Bessel function and the perturbative Sudakov factor $S_p^{\text{TMD}}(b) \equiv S_p^{\text{TMD}}(b, Q, Q^2)$ is given by

$$S_p^{\text{TMD}}(b) = \frac{C_F}{\pi} \int_{\mu_b^2}^{Q^2} \frac{d\mu^2}{\mu^2} \left(\ln \frac{Q^2}{\mu^2} - \frac{3}{2} \right) \alpha_s(\mu). \quad (5)$$

Sudakov form factor in joint resummation—Imposing a veto $\tau < \tau_0$ restricts the phase space for ISR, introducing logarithms of τ_0 and modifying the standard TMD Sudakov factor in Eq. (5), particularly below the scale $t \equiv \tau_0 Q^2$. The resulting NLL perturbative Sudakov factor for the joint resummation, incorporating both TMD and zero-jettiness logarithms, is given by

$$S_P(b) = \frac{C_F}{\pi} \left[\int_{\mu_b^2}^t \frac{d\mu^2}{\mu^2} \left(2 \ln \frac{t}{\mu^2} - \frac{3}{2} \right) - \int_{\mu_b^2}^{\tau_0 t} \frac{d\mu^2}{\mu^2} \ln \frac{\tau_0 t}{\mu^2} \right. \\ \left. + \int_t^{Q^2} \frac{d\mu^2}{\mu^2} \left(\ln \frac{Q^2}{\mu^2} - \frac{3}{2} \right) + \int_{\tau_0 t}^t \frac{d\mu^2}{\mu^2} \ln \frac{\mu^2}{\tau_0 t} \right] \alpha_s(\mu). \quad (6)$$

Here, the third integral, covering the range $[t, Q^2]$, corresponds to the standard TMD evolution above the scale t . The remaining terms encapsulate the modifications below the scale t due to the phase space restrictions imposed by the zero-jettiness veto. This result is derived independently within both perturbative QCD phase space analysis and the SCET framework. Details are provided in Supplemental Material [63], which includes Refs. [7,37,43,52].

Our calculation employs NLL accuracy for the Sudakov resummation. While the next-to-next-to-leading logarithmic (NNLL) joint resummation framework is available for the unpolarized cross section [37], the corresponding one-loop corrections for polarized TMDs under a zero-jettiness veto remain unknown. Extending the formalism to NNLL accuracy is therefore beyond the scope of this Letter.

Impact of zero jettiness on spin asymmetries—We now examine the phenomenological implications of this approach for probing the spin structure of the nucleon. As an example, we focus on the SSAs for $W^\pm/Z^0$ production in polarized pp collisions.

As is well known, the scale evolution of polarization-dependent TMDs is governed by the standard Collins-Soper equation [64], since the light cone divergence exhibits the same structure for both unpolarized and polarization dependent cases. Accordingly, the analysis presented in the previous section extends directly to the polarization-dependent case, yielding the same perturbative Sudakov factor for the Siverson distribution. In the small b limit (or equivalently, large $q_\perp$ region), the Siverson function can be further matched onto [51] a collinear twist-3 quark-gluon correlation function, known as the Qiu-Sterman function $T_{F,q}$ [50,65,66]. Recently, the collinear twist-3 factorization has been verified at one-loop order, marking a significant achievement [67,68]. A comprehensive study of the matching between the $k_\perp$ -odd TMDs and the collinear twist-3 correlations has been carried out in Refs. [69,70].

The transverse-spin-dependent cross section for $W^\pm$ boson production in polarized pp collisions can be written as [48]

$$\frac{d\sigma_{UT}(S_\perp)}{dy d^2\vec{q}_\perp} = -\sin(\phi_q - \phi_S) \sigma_0 \int_0^\infty \frac{b^2 db}{4\pi} J_1(bq_\perp) \\ \times \sum_{q,q'} |V_{qq'}|^2 T_{F,q}(x_q, x_q, \mu_b) f_{q'}(x_{q'}, \mu_b) e^{-S_P(b)}. \quad (7)$$

In general, the SSA A_N for $W^\pm$ boson production is given by

$$A_N = \frac{\int_0^{2\pi} d\phi_q 2 \sin(\phi_q - \phi_S) d\sigma_{UT}}{\int_0^{2\pi} d\phi_q d\sigma_{UU}}. \quad (8)$$

The SSA for Z^0 production takes a similar form, with appropriate modifications for the electroweak couplings.

Phenomenology—We now present numerical results to demonstrate the implications of the key results in Eq. (8). To ensure a smooth transition from the perturbative to the nonperturbative regions, the Sudakov factor is conventionally decomposed into two components,

$$S(b) = S_P(b_*) + S_{NP}(b), \quad (9)$$

where the scale μ_b in the perturbative Sudakov factor $S_P(b_*)$ is replaced by $\mu_{b_*} \equiv 2e^{-\gamma_E}/b_*$ with $b_* = b/\sqrt{1 + b^2/b_{\max}^2}$ and $b_{\max} = 1.5 \text{ GeV}^{-1}$ [62]. Note that, in our numerical evaluation, the Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) evolution of PDFs is frozen for $\mu_b^2 > t$ [52]. The parametrizations for the nonperturbative component of the Sudakov factor associated with the unpolarized quark TMDs are taken from [71],

$$S_{NP}^f(b) = 0.42 \ln \frac{Q}{Q_0} \ln \frac{b}{b_*} + 0.106 b^2, \quad (10)$$

with $Q_0 = \sqrt{2.4} \text{ GeV}$. For the Qiu-Sterman function $T_{F,q}$, the most economical choice is to parameterize it in terms of the unpolarized collinear PDF [72,73],

$$T_{F,q}(x, x, \mu) = \mathcal{N}_q(x) f_q(x, \mu), \quad (11)$$

where we adopt a parametrization for the function $\mathcal{N}_q(x)$ given in Table 2 in Ref. [49] (see also Refs. [74,75] for the state-of-the-art extractions of the Siverson functions). It is worth noting that the diagonal component of the scale evolution equation for the Qiu-Sterman function closely resembles the standard DGLAP kernel for the normal unpolarized quark PDF [76–86]. Consequently, this parametrization effectively addresses the scale dependence of the Qiu-Sterman function. In addition, we use the nonperturbative Sudakov factor for the Siverson distribution obtained from the same global fitting [49],

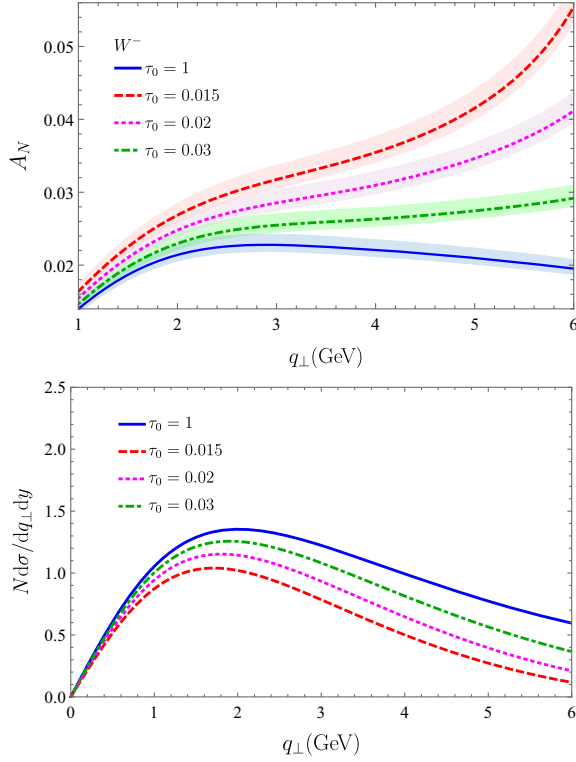


FIG. 2. Top: SSA for W^- production versus transverse momentum $q_\perp$. Bands correspond to theoretical uncertainties estimated via scale variation. Bottom: Corresponding normalized transverse momentum distribution. Results are for polarized pp collisions at the RHIC energy $\sqrt{s} = 500$ GeV and the rapidity $y = 0$, comparing various values of τ_0 (indicated by color and style).

$$S_{\text{NP}}^s(b) = 0.42 \ln \frac{Q}{Q_0} \ln \frac{b}{b_*} + 0.18b^2. \quad (12)$$

Moreover, the one-loop running coupling $\alpha_s(\mu)$ is used in our numerical estimations.

At RHIC energies, the SSAs for W^+ and Z^0 production are predicted to be small (see Supplemental Material [63]), so we focus on W^- production here. In Fig. 2, we present theoretical predictions for SSA (A_N) versus $q_\perp$ (top) and normalized $q_\perp$ distribution (bottom) at $\sqrt{s} = 500$ GeV and $y = 0$. Curves compare various τ_0 values, with $\tau_0 = 1$ as inclusive. Central predictions use the hard scale Q , with uncertainties from varying only this scale between $Q/2$ and $2Q$ in Eq. (6). As expected, SSA magnitude increases as τ_0 decreases, thus enhancing Siverts sensitivity.

The estimate of the experimental uncertainties, i.e., the statistical uncertainty of A_N , corresponding to a 68% confidence level (CL) for a given integrated luminosity $\mathcal{L}$ and beam polarization rate P is given by

$$\delta A_N = \frac{1}{P} \sqrt{\frac{1 - (A_N)^2}{\sigma \cdot \mathcal{L}}} \simeq \frac{1}{P} \frac{1}{\sqrt{\sigma \cdot \mathcal{L}}}, \quad (13)$$

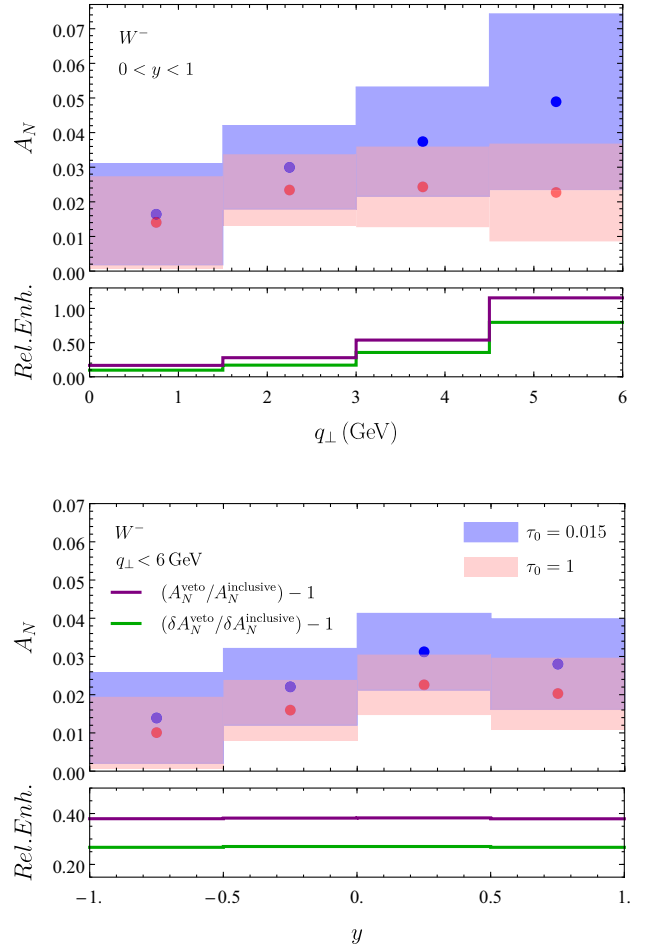


FIG. 3. Projected SSA for W^- production in polarized pp collisions at $\sqrt{s} = 500$ GeV. Top half: Distributions versus transverse momentum $q_\perp$ (integrated over $0 < y < 1$). The upper panel shows A_N for the inclusive case ($\tau_0 = 1$, red) and with a zero-jettiness veto ($\tau_0 = 0.015$, blue). Bands represent estimated statistical uncertainties (68% CL). The lower panel shows the relative enhancement of the asymmetry (purple) compared to the relative increase in statistical uncertainty (green). Bottom half: Distributions versus rapidity y (integrated over $q_\perp < 6$ GeV). Panels are as described for the top half.

where the second (approximated) equality holds due to the expected small value of A_N . For projected performance at RHIC, we assume polarized pp collisions at the center-of-mass energy $\sqrt{s} = 500$ GeV with an integrated luminosity of $\mathcal{L} = 780 \text{ pb}^{-1}$ and effective polarization $P = 53\%$ [87].

Figure 3 presents the projected A_N for W^- production, comparing the inclusive ($\tau_0 = 1$, red) and vetoed ($\tau_0 = 0.015$, blue) scenarios at $\sqrt{s} = 500$ GeV. The top half shows the distributions versus transverse momentum $q_\perp$ (integrated over $0 < y < 1$), and the bottom half shows the distributions versus rapidity y (integrated over $q_\perp < 6$ GeV). In both halves, the upper panels plot the A_N values, including estimated statistical error bands (68% CL) based on Eq. (13). A quantitative comparison

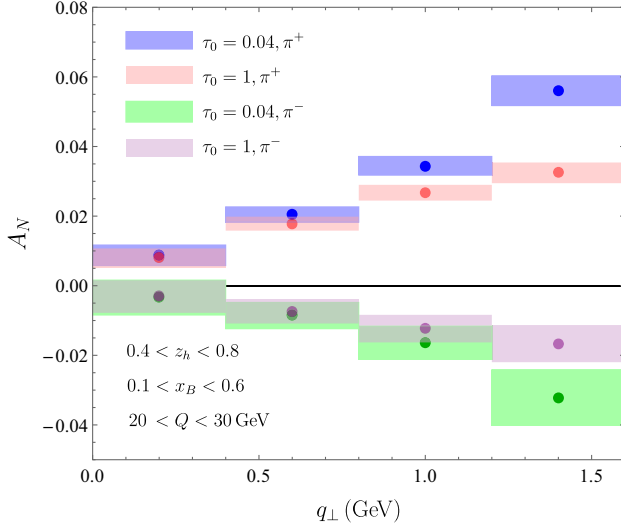


FIG. 4. SSAs for π^+ (blue and red) and π^- (green and purple) production in SIDIS at the EIC ($\sqrt{s} = 100$ GeV) as a function of $q_\perp$. Predictions compare the inclusive case ($\tau_0 = 1$, red and purple bands) with a one-jettiness veto ($\tau_0 = 0.04$, blue and green bands). Kinematic selections are $0.4 < z_h < 0.8$, $0.1 < x_B < 0.6$, and $20 < Q < 30$ GeV. Statistical uncertainty bands (68% CL) assume a projected integrated luminosity of $\mathcal{L} = 100 \text{ fb}^{-1}$ and beam polarization $P = 70\%$ [2,89].

in the $q_\perp$ distribution shows that the enhancement in the asymmetry is comparable to the estimated combined statistical uncertainty, indicating that the net gain is statistically significant and measurable at RHIC. To visualize this trade-off directly, the lower panels plot the relative enhancement of the asymmetry, $(A_N^{\text{veto}}/A_N^{\text{inclusive}}) - 1$ (purple line), against the relative increase in statistical uncertainty, $(\delta A_N^{\text{veto}}/\delta A_N^{\text{inclusive}}) - 1$ (green line). Across the entire kinematic range in both $q_\perp$ and y , the relative gain in the asymmetry (purple line) significantly exceeds the relative cost in precision (green line). For example, in the $q_\perp \approx 5$ GeV bin (top-right panel), the asymmetry increases by $\sim 115\%$, while the statistical uncertainty increases by only $\sim 80\%$. This confirms that the zero-jettiness veto yields a substantial net improvement in sensitivity. These results demonstrate that implementing a zero-jettiness veto offers a promising strategy to enhance the sensitivity of SSA measurements to the predicted sign flip of the Sivers function.

Our formalism can be naturally extended to describe the produced hadron's transverse momentum distribution in the SIDIS process. The relevant factorization framework associated with one-jettiness in deep inelastic scattering (DIS) has been previously studied in the context of inclusive observables [36,41,53,88]. In Fig. 4, we apply our method to SSAs for π^+ and π^- production at the future EIC, plotting the predictions with projected statistical uncertainty bands (68% CL) based on Eq. (13) (adapted for EIC parameters). The results, with and without imposing a veto

of $\tau_0 = 0.04$, demonstrate a clear enhancement in the asymmetries at moderate $q_\perp$ when the veto is applied. The statistical separation is clear: the error bands for the π^+ vetoed and inclusive cases show no overlap for $q_\perp \gtrsim 1$ GeV, confirming a measurable improvement in sensitivity.

Conclusions and outlook—In this Letter, we establish a novel method for probing the nucleon's spin structure by employing a zero-jettiness veto to suppress TMD evolution effects. This approach directly addresses the long-standing challenge of disentangling nonperturbative dynamics from perturbative radiation. To this end, we derive a k_T -resummation formalism for zero-jettiness events and demonstrate its phenomenological impact by analyzing SSAs in $W^\pm/Z^0$ production at RHIC. We find that the veto leads to a remarkable enhancement of the asymmetry, enabling a more definitive test of the predicted sign change of the Sivers function.

This approach broadens the potential of TMD studies across a variety of experimental settings. The ability to systematically tune the veto scale, or to employ other veto observables like jet transverse momentum or vector angularities [90], provides a powerful new lever to isolate intrinsic parton motion and constrain the nonperturbative Sudakov factor [91]. The framework's versatility extends to studying TMDs in di-jet and electron-jet production at the EIC, while offering particular advantages for TMD measurements at the LHC, where evolution effects are more pronounced. These developments pave the way for a new era of precision in mapping the nucleon's 3D structure.

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Data availability—The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

- [1] A. Accardi *et al.*, *Eur. Phys. J. A* **52**, 268 (2016).
- [2] R. Abdul Khalek *et al.*, *Nucl. Phys.* **A1026**, 122447 (2022).
- [3] D. P. Anderle *et al.*, *Front. Phys. (Beijing)* **16**, 64701 (2021).
- [4] J. C. Collins and D. E. Soper, *Nucl. Phys.* **B193**, 381 (1981); **B213**, 545(E) (1983).
- [5] J. C. Collins and D. E. Soper, *Nucl. Phys.* **B194**, 445 (1982).
- [6] X.-d. Ji, J.-p. Ma, and F. Yuan, *Phys. Rev. D* **71**, 034005 (2005).

- [7] J. Collins, *Foundations of Perturbative QCD* (Cambridge University Press, Cambridge, England, 2011), Vol. 32, ISBN 978-1-009-40184-5, 978-1-009-40183-8, 978-1-009-40182-1.
- [8] R. Boussarie *et al.*, arXiv:2304.03302.
- [9] P. J. Mulders and R. D. Tangerman, Nucl. Phys. **B461**, 197 (1996); **B484**, 538(E) (1997).
- [10] P. J. Mulders and J. Rodrigues, Phys. Rev. D **63**, 094021 (2001).
- [11] D. W. Sivers, Phys. Rev. D **41**, 83 (1990).
- [12] J. C. Collins, Nucl. Phys. **B396**, 161 (1993).
- [13] Z.-t. Liang and C. Boros, Phys. Rev. Lett. **79**, 3608 (1997).
- [14] X.-d. Ji, J.-P. Ma, and F. Yuan, Nucl. Phys. **B652**, 383 (2003).
- [15] S. J. Brodsky, D. S. Hwang, and I. Schmidt, Phys. Lett. B **530**, 99 (2002).
- [16] S. J. Brodsky, D. S. Hwang, and I. Schmidt, Nucl. Phys. **B642**, 344 (2002).
- [17] J. C. Collins, Phys. Lett. B **536**, 43 (2002).
- [18] J. C. Collins and A. Metz, Phys. Rev. Lett. **93**, 252001 (2004).
- [19] X.-d. Ji and F. Yuan, Phys. Lett. B **543**, 66 (2002).
- [20] L. D. McLerran and R. Venugopalan, Phys. Rev. D **49**, 2233 (1994).
- [21] L. D. McLerran and R. Venugopalan, Phys. Rev. D **49**, 3352 (1994).
- [22] F. Dominguez, C. Marquet, B.-W. Xiao, and F. Yuan, Phys. Rev. D **83**, 105005 (2011).
- [23] A. Metz and J. Zhou, Phys. Rev. D **84**, 051503 (2011).
- [24] E. Akcakaya, A. Schäfer, and J. Zhou, Phys. Rev. D **87**, 054010 (2013).
- [25] J. Zhou, Phys. Rev. D **89**, 074050 (2014).
- [26] P. Kotko, K. Kutak, C. Marquet, E. Petreska, S. Sapeta, and A. van Hameren, J. High Energy Phys. **09** (2015) 106.
- [27] D. Boer, M. G. Echevarria, P. Mulders, and J. Zhou, Phys. Rev. Lett. **116**, 122001 (2016).
- [28] I. Balitsky and A. Tarasov, J. High Energy Phys. **06** (2016) 164.
- [29] T. Altinoluk and R. Boussarie, J. High Energy Phys. **10** (2019) 208.
- [30] I. W. Stewart, F. J. Tackmann, and W. J. Waalewijn, Phys. Rev. Lett. **105**, 092002 (2010).
- [31] J. Gaunt, M. Stahlhofen, F. J. Tackmann, and J. R. Walsh, J. High Energy Phys. **09** (2015) 058.
- [32] R. Boughezal, C. Focke, W. Giele, X. Liu, and F. Petriello, Phys. Lett. B **748**, 5 (2015).
- [33] R. Boughezal, C. Focke, X. Liu, and F. Petriello, Phys. Rev. Lett. **115**, 062002 (2015).
- [34] R. Boughezal, J. M. Campbell, R. K. Ellis, C. Focke, W. Giele, X. Liu, F. Petriello, and C. Williams, Eur. Phys. J. C **77**, 7 (2017).
- [35] I. W. Stewart, F. J. Tackmann, and W. J. Waalewijn, J. High Energy Phys. **09** (2010) 005.
- [36] Z.-B. Kang, X. Liu, and S. Mantry, Phys. Rev. D **90**, 014041 (2014).
- [37] G. Lustermans, J. K. L. Michel, F. J. Tackmann, and W. J. Waalewijn, J. High Energy Phys. **03** (2019) 124.
- [38] S. Alioli, A. Broggio, and M. A. Lim, J. High Energy Phys. **01** (2022) 066.
- [39] S. Alioli, G. Bell, G. Billis, A. Broggio, B. Dehnadi, M. A. Lim, G. Marinelli, R. Nagar, D. Napoletano, and R. Rahn, Phys. Rev. D **109**, 094009 (2024).
- [40] M. Knobbe, D. Reichelt, and S. Schumann, J. High Energy Phys. **09** (2023) 194.
- [41] H. Cao, Z.-B. Kang, X. Liu, and S. Mantry, Phys. Rev. D **110**, 014045 (2024).
- [42] A. Jain, M. Procura, and W. J. Waalewijn, J. High Energy Phys. **04** (2012) 132.
- [43] M. Procura, W. J. Waalewijn, and L. Zeune, J. High Energy Phys. **02** (2015) 117.
- [44] P. F. Monni, L. Rottoli, and P. Torrielli, Phys. Rev. Lett. **124**, 252001 (2020).
- [45] Y. Makris, F. Ringer, and W. J. Waalewijn, J. High Energy Phys. **02** (2021) 070.
- [46] Z.-B. Kang and J.-W. Qiu, Phys. Rev. Lett. **103**, 172001 (2009).
- [47] A. Metz and J. Zhou, Phys. Lett. B **700**, 11 (2011).
- [48] Z.-B. Kang, B.-W. Xiao, and F. Yuan, Phys. Rev. Lett. **107**, 152002 (2011).
- [49] M. G. Echevarria, Z.-B. Kang, and J. Terry, J. High Energy Phys. **01** (2021) 126.
- [50] J.-w. Qiu and G. F. Sterman, Phys. Rev. Lett. **67**, 2264 (1991).
- [51] X. Ji, J.-W. Qiu, W. Vogelsang, and F. Yuan, Phys. Rev. Lett. **97**, 082002 (2006).
- [52] I. W. Stewart, F. J. Tackmann, and W. J. Waalewijn, Phys. Rev. D **81**, 094035 (2010).
- [53] Z.-B. Kang, S. Mantry, and J.-W. Qiu, Phys. Rev. D **86**, 114011 (2012).
- [54] T. T. Jouttenus, I. W. Stewart, F. J. Tackmann, and W. J. Waalewijn, Phys. Rev. D **88**, 054031 (2013).
- [55] Z.-B. Kang, X. Liu, S. Mantry, and J.-W. Qiu, Phys. Rev. D **88**, 074020 (2013).
- [56] J. R. Gaunt, M. Stahlhofen, and F. J. Tackmann, J. High Energy Phys. **04** (2014) 113.
- [57] C. W. Bauer, S. Fleming, D. Pirjol, and I. W. Stewart, Phys. Rev. D **63**, 114020 (2001).
- [58] C. W. Bauer and I. W. Stewart, Phys. Lett. B **516**, 134 (2001).
- [59] C. W. Bauer, D. Pirjol, and I. W. Stewart, Phys. Rev. D **65**, 054022 (2002).
- [60] C. W. Bauer, S. Fleming, D. Pirjol, I. Z. Rothstein, and I. W. Stewart, Phys. Rev. D **66**, 014017 (2002).
- [61] M. Beneke, A. P. Chapovsky, M. Diehl, and T. Feldmann, Nucl. Phys. **B643**, 431 (2002).
- [62] J. C. Collins, D. E. Soper, and G. F. Sterman, Nucl. Phys. **B250**, 199 (1985).
- [63] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/rvgc-sgv7> for a detailed derivation of the Sudakov factor with a zero-jettiness cut imposed, as well as additional numerical results.
- [64] X.-d. Ji, J.-P. Ma, and F. Yuan, Phys. Lett. B **597**, 299 (2004).
- [65] J.-w. Qiu and G. F. Sterman, Nucl. Phys. **B378**, 52 (1992).
- [66] J.-w. Qiu and G. F. Sterman, Phys. Rev. D **59**, 014004 (1999).
- [67] D. Rein, M. Schlegel, P. Tollkühn, and W. Vogelsang, Phys. Rev. D **112**, 114024 (2025).

- [68] D. Rein, M. Schlegel, P. Tollkühn, and W. Vogelsang, *Phys. Rev. Lett.* **135**, 251901 (2025).
- [69] J. Zhou, F. Yuan, and Z.-T. Liang, *Phys. Rev. D* **81**, 054008 (2010).
- [70] F. Rein, S. Rodini, A. Schäfer, and A. Vladimirov, *J. High Energy Phys.* **01** (2023) 116.
- [71] P. Sun, J. Isaacson, C. P. Yuan, and F. Yuan, *Int. J. Mod. Phys. A* **33**, 1841006 (2018).
- [72] P. Sun and F. Yuan, *Phys. Rev. D* **88**, 034016 (2013).
- [73] M. G. Echevarria, A. Idilbi, Z.-B. Kang, and I. Vitev, *Phys. Rev. D* **89**, 074013 (2014).
- [74] A. Bacchetta, F. Delcarro, C. Pisano, and M. Radici, *Phys. Lett. B* **827**, 136961 (2022).
- [75] M. Bury, A. Prokudin, and A. Vladimirov, *J. High Energy Phys.* **05** (2021) 151.
- [76] Z.-B. Kang and J.-W. Qiu, *Phys. Rev. D* **79**, 016003 (2009).
- [77] J. Zhou, F. Yuan, and Z.-T. Liang, *Phys. Rev. D* **79**, 114022 (2009).
- [78] W. Vogelsang and F. Yuan, *Phys. Rev. D* **79**, 094010 (2009).
- [79] V. M. Braun, A. N. Manashov, and B. Pirnay, *Phys. Rev. D* **80**, 114002 (2009); **86**, 119902(E) (2012).
- [80] J. P. Ma and H. Z. Sang, *J. High Energy Phys.* **04** (2011) 062.
- [81] A. Schafer and J. Zhou, *Phys. Rev. D* **85**, 117501 (2012).
- [82] J. P. Ma and Q. Wang, *Phys. Lett. B* **715**, 157 (2012).
- [83] Z.-B. Kang and J.-W. Qiu, *Phys. Lett. B* **713**, 273 (2012).
- [84] P. Sun and F. Yuan, *Phys. Rev. D* **88**, 114012 (2013).
- [85] J. Zhou, *Phys. Rev. D* **92**, 074016 (2015).
- [86] O. del Rio, A. Prokudin, I. Scimemi, and A. Vladimirov, *Phys. Rev. D* **110**, 016003 (2024).
- [87] The total integrated luminosity $\mathcal{L} = 780 \text{ pb}^{-1}$ combines data from the 2017 RHIC run (350 pb^{-1} with polarization $P \approx 55\%$) and the 2022 run (430 pb^{-1} with $P \approx 50\%$). An effective polarization of $P = 53\%$ is used in our calculation.
- [88] D. Kang, C. Lee, and I. W. Stewart, *Phys. Rev. D* **88**, 054004 (2013).
- [89] P. Nadel-Turonski, *Acta Phys. Pol. B Proc. Suppl.* **18**, 1 (2025).
- [90] P. Bijl, S. Niedenzu, and W. J. Waalewijn, *Phys. Rev. D* **109**, 014011 (2024).
- [91] J. Collins and T. Rogers, *Phys. Rev. D* **91**, 074020 (2015).