Threshold resummation for polarized (semi-)inclusive deep inelastic scattering

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We explore the effects of the resummation of large logarithmic perturbative corrections to double-longitudinal spin asymmetries for inclusive and semi-inclusive deep inelastic scattering in fixed-target experiments. We find that the asymmetries are overall rather robust with respect to the inclusion of the resummed higher-order terms. Significant effects are observed at fairly high values of $x$, where resummation tends to decrease the spin asymmetries. This effect turns out to be more pronounced for semi-inclusive scattering. We also investigate the potential impact of resummation on the extraction of polarized valence quark distributions in dedicated high-$x$ experiments.

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I. INTRODUCTION

Longitudinal double-spin asymmetries in inclusive and semi-inclusive deep inelastic scattering have been prime sources of information on the nucleon’s spin structure for several decades. They may be used to extract the helicity parton distributions of the nucleon,

$$\Delta f(x, Q^2) \equiv f^+(x, Q^2) - f^-(x, Q^2),$$

where $f^+$ and $f^-$ are the distributions of parton $f = q, \bar{q}, g$ with positive and negative helicity, respectively, when the parent nucleon has positive helicity. $x$ denotes the momentum fraction of the parton and $Q$ the hard scale at which the distribution is probed. Inclusive polarized deep inelastic scattering (DIS), $\ell p \rightarrow \ell X$, offers access to the combined quark and antiquark distributions for a given flavor, $\Delta q + \Delta\bar{q}$, whereas in semi-inclusive deep inelastic scattering (SIDIS), $\ell p \rightarrow \ell hX$, one exploits the fact that a produced hadron $h$ (like a $\pi^+$) may for instance have a quark of a certain flavor as a valence quark, but not the corresponding antiquark [1]. In this way, it becomes possible to separate quark and antiquark distributions in the nucleon from one another, as well as to better determine the distributions for the various flavors. HERMES [2] and recent COMPASS [3] measurements have marked significant progress concerning the accuracy and kinematic coverage of polarized SIDIS measurements. The inclusive measurements have improved vastly as well [4–8]. Some modern analyses of spin-dependent parton distributions include both inclusive and semi-inclusive data [9–11]. In addition, high-precision data for polarized SIDIS will become available from experiments to be carried out at the Jefferson Lab after the CEBAF upgrade to a 12 GeV beam [12]. Here the focus will be on the large-$x$ regime.

A good understanding of the theoretical framework for the description of spin asymmetries in lepton scattering is vital for a reliable extraction of polarized parton distributions. In a recent paper [13] we have investigated the effects of QCD threshold resummation on hadron multiplicities in SIDIS in the HERMES and COMPASS kinematic regimes. SIDIS is characterized by two scaling variables, Bjorken-$x$ and a variable $z$ given by the energy of the produced hadron over the energy of the virtual photon in the target rest frame. Large logarithmic corrections to the SIDIS cross section arise when the corresponding partonic variables become large, corresponding to scattering near a phase space boundary, where real-gluon emission is suppressed. This is typically the case for the presently relevant fixed-target kinematics. Threshold resummation addresses these logarithms to all orders in the strong coupling. In [13] we found fairly significant resummation effects on the spin-averaged multiplicities. Since the spin-dependent cross section is subject to similar logarithmic corrections as the unpolarized one, it is worthwhile to explore the effects of resummation on the spin asymmetries. This is the goal of the present paper. Our calculations will be carried out both for inclusive DIS and for SIDIS. We note that previous work [14, 15] has addressed the large-$x$ resummation for the inclusive spin-dependent structure function $g_1$, with a focus on the moments of $g_1$ and their $Q^2$-dependence. In this paper we are primarily concerned with spin asymmetries and with semi-inclusive scattering.

Our work will use the framework developed in [13]. In Section [1] we briefly review the basic terms and definitions relevant for longitudinal spin asymmetries, and we describe the extension of threshold resummation to the polarized case. In Section [II] our phenomenological results are presented. We compare our resummed inclusive and semi-inclusive spin asymmetries with available HERMES, COMPASS and Jefferson Lab data. We also discuss the relevance of resummation for the extraction of $\Delta u/u$ and $\Delta d/d$ at large values of $x$.

II. RESUMMATION FOR LONGITUDINAL SPIN ASYMMETRIES IN DIS AND SIDIS

A. Leading and next-to-leading order expressions

We first consider the polarized SIDIS process $\ell(k)\bar{p}(P) \rightarrow \ell(k')h(P_h)X$ with longitudinally polarized
beam and target and with an unpolarized hadron in the final state. The corresponding double-spin asymmetry is given by a ratio of structure functions \[ F_1(x, z, Q^2) \]:

\[ A^h_1(x, z, Q^2) \approx \frac{g_h^b(x, z, Q^2)}{F_1^h(x, z, Q^2)}, \]  

where \( Q^2 = -q^2 \) with \( q \) the momentum of the virtual photon, \( x = Q^2/(2P \cdot q) \) is the usual Bjorken variable, and \( z \equiv P \cdot P_h/P \cdot q \) the corresponding hadronic scaling variable associated with the fragmentation process.

Using factorization, the polarized structure function \( g_h^b \), which appears in the numerator of Eq. \( \text{(2)} \), can be written as

\[ 2g_h^b(x, z, Q^2) = \sum_{f, f'} q_{f' \rightarrow q} \int \frac{d\hat{x}}{\hat{x}} \int \frac{d\tilde{z}}{\tilde{z}} \Delta f \left( \frac{x}{\hat{x}}, \mu^2 \right) \times D^h_{f'} \left( \frac{z}{\tilde{z}}, \mu^2 \right) \Delta C_{f'f} \left( \hat{x}, \tilde{z}, Q^2, \frac{z}{\hat{x}}, \frac{\alpha_s(\mu^2)}{\hat{x}^2} \right), \]  

\( \text{(3)} \)

where \( \Delta f(\xi, \mu^2) \) denotes the polarized distribution function for parton \( f \) of Eq. \( \text{(1)} \), whereas \( D^h_{f'}(\zeta, \mu^2) \) is the corresponding fragmentation function for parton \( f' \) going to the observed hadron \( h \). The \( \Delta C_{f'f} \) are spin-dependent coefficient functions. We have set all factorization and renormalization scales equal and collectively denoted them by \( \mu \). In \( \text{(3)} \), \( \hat{x} \) and \( \tilde{z} \) are the partonic counterparts of the hadronic variables \( x \) and \( z \). Setting for simplicity \( \mu = Q \), we use the short-hand-notation

\[ 2g_h^b(x, z, Q^2) \equiv \sum_{f, f'} q_{f' \rightarrow q} [\Delta f \otimes \Delta C_{f'f} \otimes D^h_{f'}] \left( x, z, Q^2 \right) \]  

\( \text{(4)} \)

for the convolutions in \( \text{(3)} \). A corresponding expression for the “transverse” unpolarized structure function \( 2F_1^h \) can be written by replacing the polarized parton distributions with the unpolarized ones, and using unpolarized coefficient functions which we denote here by \( C_{f'f} \).

The spin-dependent hard-scattering coefficient functions \( \Delta C_{f'f} \) in \( \text{(3)} \) can be computed in perturbation theory:

\[ \Delta C_{f'f} = \Delta C_{f'f}^{(0)} + \frac{\alpha_s(\mu^2)}{2\pi} \Delta C_{f'f}^{(1)} + O(\alpha_s^2). \]  

\( \text{(5)} \)

At leading order (LO), we have

\[ \Delta C_{qq}(\hat{x}, \tilde{z}) = \Delta C_{qq}(\hat{x}, \tilde{z}) = e_q^2 \delta(1-\hat{x})\delta(1-\tilde{z}), \]  

\( \text{(6)} \)

with the quark’s fractional charge \( e_q \). All other coefficient functions vanish. The same result holds for the LO coefficient function for the spin-averaged structure function \( 2F_1^h \).

Hence the asymmetry in Eq. \( \text{(2)} \) reduces to

\[ A^h_1 = \sum_q e_q^2 [\Delta q(x, Q^2)D^h_q(z, Q^2) + \Delta \bar{q}(x, Q^2)D^h_{\bar{q}}(z, Q^2)] \]  

\[ \sum_q e_q^2 [g(x, Q^2)D^h_g(z, Q^2) + \bar{q}(x, Q^2)D^h_{\bar{g}}(z, Q^2)] \]  

\( \text{(7)} \)

At next-to-leading order (NLO), Eq. \( \text{(3)} \) becomes

\[ 2g_h^b(x, z, Q^2) = \sum_q e_q^2 \left\{ \Delta q(x, Q^2)D^h_q(z, Q^2) + \bar{q}(x, Q^2)D^h_{\bar{q}}(z, Q^2) \right\} \left. \right. \]  

\[ + \frac{\alpha_s(Q^2)}{2\pi} \left\{ \left[ \Delta q \otimes D^h_q + \Delta \bar{q} \otimes D^h_{\bar{q}} \right] \otimes \Delta C^{(1)}_{qq} \right. \]  

\[ + \left. \left[ \Delta q + \Delta \bar{q} \right] \otimes \Delta C^{(1)}_{gq} \right. \]  

\[ + \Delta g \otimes \Delta C^{(1)}_{qg} \left\} \right. \left( z, x, Q^2 \right), \]  

\( \text{(8)} \)

where the symbol \( \otimes \) denotes the convolution defined in Eqs. \( \text{(3)} \). \( \text{(11)} \). The explicit expressions for the spin-dependent NLO coefficients \( \Delta C^{(1)}_{f'f} \) have been derived in \( \text{[16, 17]} \). The corresponding spin-averaged NLO coefficient functions \( C^{(1)}_{f'f} \) may be found in \( \text{[13, 16, 21]} \).

In the case of inclusive polarized DIS, the longitudinal spin asymmetry \( A_1 \) is given in analogy with \( \text{(2)} \) by

\[ A_1(x, Q^2) \approx \frac{g_1(x, Q^2)}{F_1(x, Q^2)}. \]  

\( \text{(9)} \)

The inclusive structure functions \( g_1 \) and \( F_1 \) have expressions analogous to their SIDIS counterparts, except for the fact that they do not contain any fragmentation functions, of course. The unpolarized and polarized NLO coefficient functions for inclusive DIS may be found at many places; see, for example \( \text{[20, 22]} \).

B. Threshold resummation

As was discussed in \( \text{[13]} \), the higher-order terms in the spin-averaged SIDIS coefficient function \( C_{qq} \) introduce large terms near the “partonic threshold” \( \hat{x} \to 1 \), \( \tilde{z} \to 1 \). The same is true for the spin-dependent \( \Delta C_{qq} \).

At NLO, choosing again for simplicity the scale \( \mu = Q \), one has

\[ \Delta C_{qq}^{(1)}(\hat{x}, \tilde{z}) \sim e_q^2 C_F \left[ \right. \right. \]  

\[ + 2\delta(1-\hat{x}) \left( \frac{\ln(1-\hat{x})}{1-\hat{x}} \right) + 2\delta(1-\tilde{z}) \left( \frac{\ln(1-\tilde{z})}{1-\tilde{z}} \right) \]  

\[ + \frac{2}{(1-\hat{x})+(1-\tilde{z})} - 8\delta(1-\hat{x})\delta(1-\tilde{z}) \]  

\( \text{(10)} \)

where the “+”-distribution is defined as usual. The expression on the right-hand side is in fact identical to the one for the unpolarized coefficient function near threshold \( \text{[13]} \). At the \( k \)th order of perturbation theory, the coefficient function contains terms of the form

\[ \alpha_k^q(1-\tilde{z}) \left( \frac{\ln^{k-1}(1-\tilde{z})}{1-\tilde{z}} \right)^+ + \alpha_k^q(1-\hat{x}) \left( \frac{\ln^{k-1}(1-\hat{x})}{1-\hat{x}} \right)^+, \]

or
“mixed” distributions $\alpha_k^k \left( \frac{\ln^m(1-x)}{1-x} + \frac{\ln^n(1-z)}{1-z} \right)$ with $m + n = 2k - 2$, plus terms less singular by one or more logarithms. Again, each of these terms will appear equally in the unpolarized and in the polarized coefficient function. The reason for this is that the terms are associated with emission of soft gluons [13], which does not care about spin. Threshold resummation addresses the large logarithmic terms to all orders in the strong coupling. The resummation for the case of SIDIS was carried out in [13]. Given these results and the equality of the spin-averaged and spin-dependent coefficient function, the resummation for both cases is performed in [13].

Having the resummation for both $g_\lambda^h$ and $F_\lambda^h$, we obtain resummed predictions for the experimentally relevant spin asymmetry $A_{\lambda}^h$.

In [13, 23, 24] threshold resummation for SIDIS was derived using an eikonal approach, for which exponentiation of the threshold logarithms is achieved in Mellin space. One takes Mellin moments of $g_\lambda^h$ separately in the two independent variables $x$ and $z$ [13, 25]:

$$\tilde{g}_\lambda^h(N, M, Q^2) \equiv \int_0^1 dx x^{N-1} \int_0^1 dz z^{M-1} g_\lambda^h(x, z, Q^2).$$

With this definition, Eq. (4) takes the form (again at scale $\mu = Q$)

$$2g_\lambda^h(N, M, Q^2) = \sum_{f, f', q, q'} \Delta \tilde{f}(N) \times \Delta \tilde{C}_{f f'}(N, M, \alpha_s(Q^2)) \tilde{D}_{f' f}^\lambda(Q^2),$$

(11)

where the moments of the polarized parton distributions and the fragmentation functions are defined as

$$\Delta \tilde{f}(N) = \int_0^1 dx x^{N-1} \Delta f(x, Q^2),$$

$$\Delta \tilde{C}_{f f'}(N, M, \alpha_s(Q^2)) \tilde{D}_{f f}^\lambda(Q^2),$$

(12)

and the double Mellin moments of the polarized coefficient functions are

$$\Delta \tilde{C}_{f f'}(N, M, \alpha_s(Q^2)) = \int_0^1 d\hat{x} \hat{x}^{N-1} \int_0^1 d\hat{z} \hat{z}^{M-1} \Delta C_{f f'}(\hat{x}, \hat{z}, 1, \alpha_s(Q^2)) \times \Delta \tilde{C}_{f f'}(N, M, \alpha_s(Q^2)) \text{,}$$

(13)

Large $\hat{x}$ and $\hat{z}$ in $\Delta \tilde{C}_{f f'}$ correspond to large $N$ and $M$ in $\Delta \tilde{C}_{f f'}$, respectively.

The resummed spin-dependent coefficient function is identical to the spin-averaged one of [13] and reads to next-to-leading logarithmic (NLL) accuracy in the MS-scheme:

$$\Delta \tilde{C}_{qq}(N, M, \alpha_s(Q^2)) = c^2 H_{qq}(\alpha_s(Q^2)) \times \exp \left[ 2 \int \frac{Q^2}{k_+^2} \frac{dk_+^2}{k_+^2} A_q(\alpha_s(k_+^2)) \ln \left( \frac{k_+}{Q N M} \right) \right],$$

(14)

where $N \equiv n^{\gamma_E} M \equiv M^{\gamma_E}$, with $\gamma_E$ the Euler constant, and

$$A_q(\alpha_s) = \frac{\alpha_s}{\pi} A_q^{(1)}(\alpha_s) + \left( \frac{\alpha_s}{\pi} \right)^2 A_q^{(2)} + \ldots$$

(15)

is a perturbative function. The coefficients required to NLL read

$$A_q^{(1)} = C_F, \quad A_q^{(2)} = \frac{1}{2} C_F \left[ C_A \left( \frac{67}{18} - \frac{\pi^2}{6} \right) - \frac{5}{9} N_f \right],$$

(16)

where $C_F = 4/3, C_A = 3$ and $N_f$ is the number of active flavors. Furthermore,

$$H_{qq}(\alpha_s) = 1 + \frac{\alpha_s}{\pi} C_F \left[ -8 + \frac{\pi^2}{3} \right] + \mathcal{O}(\alpha_s^2).$$

(17)

The explicit NLL expansion of the exponent in (15) is given by [13]

$$\int \frac{Q^2}{k_+^2} \frac{dk_+^2}{k_+^2} A_q(\alpha_s(k_+^2)) \ln \left( \frac{k_+}{Q N M} \right) \approx h_{q(1)}^1 \left( \frac{\lambda_{NM}}{2} \right) + h_{q(2)}^1 \left( \frac{\lambda_{NM}}{2}, \frac{Q^2}{\mu^2}, \frac{Q^2}{\mu_F^2} \right),$$

(18)

where

$$\lambda_{NM} \equiv b_0 \alpha_s(\mu^2) \left( \log N + \log M \right),$$

$$h_{q(1)}^1(\lambda) = A_q^{(1)} b_0 \ln(1 - 2\lambda),$$

$$h_{q(2)}^1(\lambda, Q^2/\mu^2, Q^2/\mu_F^2) = -A_q^{(2)} b_0 \left[ 2\lambda + \ln(1 - 2\lambda) \right] + \frac{A_q^{(1)} b_0}{2\pi} \left[ 2\lambda + \ln(1 - 2\lambda) \right] + \frac{A_q^{(1)} b_1}{2\pi b_0} \left[ 2\lambda + \ln(1 - 2\lambda) \right] \ln \frac{Q^2}{\mu^2} - \frac{A_q^{(1)} b_0}{\pi b_0} \lambda \ln \frac{Q^2}{\mu_F^2},$$

(19)

with

$$b_0 = \frac{11C_A - 4T_R N_f}{12\pi},$$

$$b_1 = \frac{17C_A^3 - 10C_A T_R N_f - 6C_F T_R N_f}{24\pi^2}.$$

(20)

The functions $h_{q(1)}^1, h_{q(2)}^1$ collect all leading-logarithmic and NLL terms in the exponent, which are of the form
\[ \alpha_s^k \ln^n N \ln^m M \] with \( n + m = k + 1 \) and \( n + m = k \), respectively. Note that we have restored the full dependence on the factorization and renormalization scales in the above expressions.

The polarized moment-space structure function \( g_1^{h, \text{res}} \) resummed to NLL is obtained by inserting the resummed coefficient function into Eq. (12). To get the physical hadronic structure function \( g_1^{h, \text{res}} \) one needs to take the Mellin inverse of the moment-space expression. As in Ref. [13], we choose the required integration contours in complex \( N, M \)-space according to the minimal prescription of Ref. [26], in order to properly deal with the singularities arising from the Landau pole due to the divergence of the perturbative running strong coupling constant \( \alpha_s \) at scale \( \Lambda_{\text{QCD}} \). Moreover, we match the resummed \( g_1^{h, \text{res}} \) to its NLO value, i.e. we subtract the \( \mathcal{O}(\alpha_s) \) expansion from the resummed expression and add the full NLO result:

\[ g_1^{h, \text{match}} \equiv g_1^{h, \text{res}} - g_1^{h, \text{res}}(\mathcal{O}(\alpha_s)) + g_1^{h, \text{NLO}}. \]  

The final resummed and matched expression for the spin asymmetry \( A_1^{h, \text{res}} \) is then given by

\[ A_1^{h, \text{res}}(x, z, Q^2) \equiv \frac{g_1^{h, \text{match}}(x, z, Q^2)}{F_1^{h, \text{match}}(x, z, Q^2)}. \]  

Similar considerations can be made for inclusive DIS, where again the resummation for \( g_1 \) proceeds identically to that of \( F_1 \) in moment space. Only single Mellin moments of the structure function have to be taken:

\[ \tilde{g}_1(N, Q^2) \equiv \int_0^1 dx x^{N-1} g_1(x, Q^2). \]  

The threshold resummed coefficient function is the same as in the spin-averaged case and is discussed for example in Ref. [13]. We note that the outgoing quark in the process \( \gamma^* q \to q \) remains “unobserved” in inclusive DIS. At higher orders this is known to generate Sudakov suppression effects that can be described in a similar way to SIDIS, where the outgoing quark fragments and hence is “observed”, so that both the initial and the final quark contribute to Sudakov enhancement. As a result, resummation effects are generally larger in SIDIS than in DIS, for given kinematics.

III. PHENOMENOLOGICAL RESULTS

We now analyze numerically the impact of threshold resummation on the semi-inclusive and inclusive DIS asymmetries \( A_1^{h, \text{res}} \) and \( A_1 \). Given that the resummed exponents are identical for the spin-averaged and spin-dependent structure functions, we expect the resummation effects to be generally very modest. On the other hand, it is also clear that the effects will not cancel identically in the spin asymmetries: Even though the resummed exponents for \( g_1 \) and \( F_1 \) are identical in Mellin-moment space, they are convoluted with different parton distributions and hence no longer give identical results after Mellin inversion. Moreover, the matching procedure also introduces differences since the NLO coefficient functions are somewhat different for \( g_1 \) and \( F_1 \). It is therefore still relevant to investigate the impact of resummation on the spin asymmetries. We will compare our results to data sets from HERMES [2] and COMPASS [3, 5]. In addition, we present some results relevant for measurements at the Jefferson Laboratory [6, 7], in particular those to be carried out in the near future after the CEBAF upgrade to 12 GeV [12].

For our calculations we use the NLO polarized parton distribution functions of Ref. [9] and the unpolarized ones of Ref. [25]. Our choice of the latter is motivated by the fact that this set was also adopted as the baseline unpolarized set in Ref. [9], so that the two sets are consistent in the sense that the same strong coupling constant is used. Additionally, in the case of SIDIS we choose the “de Florian-Sassot-Stratmann” [29] NLO set of fragmentation functions. In this work, we choose to focus only on pions in the final state. Resummation effects for other hadrons will be very similar. The factorization and renormalization scales are set to \( Q \).

Figures 1 and 2 present comparisons of our resummed calculations with HERMES data [2] for semi-inclusive (\( \pi^- \)) and inclusive DIS, respectively, both off a proton target at \( \sqrt{s} \approx 7.25 \) GeV. The error bars show the sta-
FIG. 2: Spin asymmetry for inclusive polarized DIS off a proton target. The data points are from [4] and show statistical errors only. The $\langle x \rangle$ and $\langle Q^2 \rangle$ values were taken accordingly to the HERMES measurements.

Fig. 3: Same as Fig. 1 but comparing to the COMPASS measurements [3].
FIG. 4: Same as Fig. 2 but comparing to the COMPASS measurements [5].

$g_{1,p}, g_{1,n}, F_{1,p}, F_{1,n}$ in [25]. Up to certain refinements required by the fact that measurements of the ratios $g_{1,p}/F_{1,p}$ and $g_{1,n}/F_{1,n}$ are more readily available than those of the individual structure functions, this is essentially the approach used by the Hall-A Collaboration (alternatively, one may also use the corresponding spin asymmetry for the deuteron instead of the neutron one [7]). In the following we explore the typical size of the corrections to the ratios due to higher orders. Figure 6 shows first of all the structure function ratios on the right-hand side of (25), computed at NLO using as before the polarized and unpolarized parton distribution functions of [9] and [28], respectively (solid lines). We have again chosen $Q^2 = x \times 8$ GeV$^2$. Using (25), these ratios would correspond to the “direct experimental determinations” of $R_u$ and $R_d$. The dashed lines in the figure show the actual ratios $(\Delta u + \Delta \bar{u})/(u + \bar{u})$ and $(\Delta d + \Delta \bar{d})/(d + \bar{d})$ as given by the sets of parton distribution functions that we use. Any difference between the solid and dashed lines is, therefore, a measure of the significance of effects related to strange quarks and antiquarks, and to NLO corrections. As one can see, these have relatively modest size. Finally, we estimate the potential effect of resummation on $R_u, R_d$. Following [30, 31], we define ‘resummed’ quark (and antiquark) distributions by demanding that their contributions to the structure functions $g_1, F_1$ match those of the corresponding NLO distributions, which is ensured by setting

$$\tilde{q}^{N,\text{res}}(Q^2) \equiv \frac{\tilde{C}_q^{\text{NLO}}(N, \alpha_s(Q^2))}{\tilde{C}_q^{\text{NLO}}(N, \alpha_s(Q^2))} \tilde{q}^{N,\text{NLO}}(Q^2)$$

in Mellin-moment space. Here, $\tilde{C}_q^{\text{NLO}}$ and $\tilde{C}_q^{\text{res}}$ are the NLO and resummed quark coefficient functions for the inclusive structure function $F_1$, respectively. We match the resummed coefficient function to the NLO one by subtracting out its NLO contribution and adding the full NLO one, in analogy with (22). Equation (26) can be straightforwardly extended to the spin-dependent case. The ratios $R_u, R_d$ for these ‘resummed’ parton distributions are shown by the dotted lines in Fig. 6. As one can see, they are quite close to the other results, indicating that resummation is not likely to induce very large changes in the parton polarizations extracted from future high-precision data. For illustration, we also show the Hall-A [6] and CLAS [7] data in the figure, which have been obtained using parton-model relations for the inclusive structure functions, similar to (25). One can see that the error bars of the data are presently still larger than the differences between our various theoretical results. This situation is expected to be improved with the advent of the Jefferson Lab 12-GeV upgrade [12] or an Electron Ion Collider [32]. As is well-known, SIDIS measurements provide additional information on $R_u, R_d$, albeit so far primarily at lower $x$ [2].

IV. CONCLUSIONS

We have investigated the size of threshold resummation effects on double-longitudinal spin asymmetries for inclusive and semi-inclusive deep inelastic scattering in fixed-target experiments. Overall, the asymmetries are...
The effects are included in their extraction, using Eq. (26). The dotted lines represent the NLO sets of [9] and [28]. The dashed lines show the actual parton distribution ratios as structure functions on the right-hand sides of Eq. (25), while the solid lines show the ratios of parton distributions. This work was supported in part by the German Bundesministerium für Bildung und Forschung (BMBF), grant no. 05P12VTCTG.

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