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QCD EFFECTS IN SEMI-INCLUSIVE DEEP INELASTIC SCATTERING

FROM A POLARIZED TARGET

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A B S T R A C T

The next-to-leading logarithmic corrections to cross-sections for semi-inclusive deep inelastic scattering from a polarized target are worked out. Predictions for pion asymmetries and pion charge ratios are given.

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1. - INTRODUCTION

Deep inelastic scattering provides a good testing ground for the dynamics of strong interactions. The quantitative agreement between scaling violations¹⁾ predicted by perturbative QCD and the data accumulated so far²⁾ can be considered as a striking success of the theory. Nevertheless spin-dependent observables which are dominated by short-distance physics³⁾ have to be confronted with the predictions of perturbative QCD in order to get a complete description of strong interactions. One such process is muoproduction using a polarized target since the EMC collaboration⁴⁾ will soon allow us to compare their data with theoretical predictions.

Recently a detailed study⁵⁾ of scaling violations for spin-dependent structure functions has been worked out starting from partonic structure functions, constrained by the two Bjorken sum rules⁶⁾, in agreement with SLAC-Yale data⁷⁾. For semi-inclusive deep inelastic scattering (hereafter denoted SIDI), QCD predicts at a next-to-leading level a factorization breaking⁸⁾ which has been seen in leptonproduction processes⁹⁾. Calculations for SIDI processes involving a polarized target have been done¹⁰⁾ only at the partonic level, predicting appreciable spin dependent effects. Next-to-leading corrections to helicity cross-sections have been calculated for Drell-Yan processes¹¹⁾ and found to cancel when considering the asymmetry. Our purpose is to refine the predictions of Ref. 10) by computing QCD corrections at the next-to-leading level and by using a more realistic set of spin-dependent structure functions and also partonic fragmentation functions derived from experiment. Such refinements are expected to mainly affect the small x region which is precisely the kinematical domain accurately investigated by the EMC collaboration.

The paper is organized as follows. We give in Section 2 the expressions for the helicity-dependent SIDI cross-sections. We deal in Section 3 with numerical results to be directly confronted to the NA2 experiment. We conclude in Section 4.

2. - HELICITY-DEPENDENT SIDI CROSS-SECTIONS

The usual kinematical variables relevant for the reaction $\ell H \rightarrow \ell' H' + X$ are:

$$x_H = Q^2 / 2 P_H \cdot q, \quad y = P_H \cdot q / P_H \cdot \ell, \quad z_{H'} = P_{H'} \cdot P_H / P_H \cdot q$$

where $Q^2 = -q^2 = (\ell - \ell')^2$. In the quark parton model the differential cross-section for the unpolarized case has the schematic form:

$$\frac{d\sigma}{dx_H dy dz_{H'}} = \frac{2\pi\alpha^2}{Q^2} \sum_i e_i^2 \int_{x_H}^1 \frac{dx}{x} \int_{z_{H'}}^1 \frac{dz}{z} \tilde{\sigma}(x, z) q_i^H\left(\frac{x_H}{x}\right) D_i^{H'}\left(\frac{z_{H'}}{z}\right) \quad (1)$$

where $q_i^H(x)$ is the quark of type i distribution in the hadron H and $D_i^{H'}(z)$ the fragmentation function of the parton of type i in the hadron H' . For the polarized case:

$$\frac{d\sigma_{\lambda h}}{dx_H dy dz_{H'}} = \frac{2\pi\alpha^2}{Q^2} \sum_i e_i^2 \int_{x_H}^1 \frac{dx}{x} \int_{z_{H'}}^1 \frac{dz}{z} \tilde{\sigma}_{\lambda h}(x, z) q_i^{Hh}\left(\frac{x_H}{x}\right) D_i^{H'}\left(\frac{z_{H'}}{z}\right) \quad (2)$$

where $\lambda(h)$ refers to the incoming lepton (parton) helicity, $q_i^{Hh}(x)$ is the distribution inside the hadron H of the quark i with helicity $h = + (-)$ parallel (antiparallel) to the hadron H nuclear helicity. The quantity $\tilde{\sigma}(x, z)$ is the subprocess cross-section. In the quark parton model:

$$\tilde{\sigma}(x, z) = \frac{1 + (1-y)^2}{y} \delta(1-x) \delta(1-z) \quad (3)$$

and

$$\tilde{\sigma}_{\lambda h}(x, z) = \frac{\delta(1-x) \delta(1-z)}{y} [1 + \lambda h + (1-y)^2 (1 - \lambda h)] \quad (4)$$

The observable we will consider is the double spin asymmetry

$$\begin{aligned} A_1^{H'}(x_H, y, z_{H'}) &= \frac{\frac{d\sigma_{++}}{dx_H dy dz_{H'}} - \frac{d\sigma_{+-}}{dx_H dy dz_{H'}} + \frac{d\sigma_{--}}{dx_H dy dz_{H'}} - \frac{d\sigma_{-+}}{dx_H dy dz_{H'}}}{\frac{d\sigma}{dx_H dy dz_{H'}}} \\ &= \frac{d\Delta\sigma/dx_H dy dz_{H'}}{d\sigma/dx_H dy dz_{H'}} \end{aligned} \quad (5)$$

Its expression reads:

$$A_1^{H'}(x_H, y, z_{H'}) = \frac{\sum_i e_i^2 \Delta q_i^H(x_H) D_i^{H'}(z_{H'}) [1 - (1-y)^2]}{\sum_i e_i^2 q_i^H(x_H) D_i^{H'}(z_{H'}) [1 + (1-y)^2]} \quad (6)$$

where $\Delta q_i^H(x) = q_i^{H+}(x) - q_i^{H-}(x)$.

At the leading level one has simply to replace, in Eq. (6), the structure and fragmentation functions by the evolved ones, the evolution scale usually chosen being Q^2 . To first order in α_s three subprocesses contribute: the first one involves an initial quark and an outgoing quark, the second an initial quark and an outgoing gluon and the third one an initial gluon and an outgoing quark. As it is well known, the first subprocess is infra-red divergent. The regularization method we adopt consists of giving a small mass m_g to gluons. This singularity, which corresponds to a $\ln^2(Q^2/m_g^2)$ term cancels when one adds virtual gluon corrections to real emission. For the three subprocesses we are left with mass singularities (i.e., terms like $\ln(Q^2/m_g^2)$) which are absorbed into evolved structure functions. As previously mentioned the first-order corrections have already been computed for the unpolarized case. Their expression can be found for example in Ref. 13) and will not be recalled here.

Corrections to the quantity $d\Delta\sigma/dx_H dy dz_H$, defined in Eq. (5) for the first subprocess read (after adding real and virtual contributions):

$$\begin{aligned} d\Delta\tilde{\sigma}(x, y, z) = & \frac{2\pi\alpha_s^2}{Q^2} e_i^2 \frac{1 - (1-y)^2}{y} \frac{2\alpha_s}{3\pi} \left\{ -\delta(1-x) \delta(1-z) \right. \\ & \left(\frac{7}{2} + \frac{2\pi^2}{3} \right) + \delta(1-x) \left[\left(\frac{1+z^2}{1-z} \right)_+ \log\left(\frac{Q^2}{m_g^2}\right) + \frac{1+z^2}{(1-z)_+} \log(1-z) - (1-z) \right] \\ & + \delta(1-z) \left[\left(\frac{1+x^2}{1-x} \right)_+ \log\left(\frac{Q^2}{m_g^2}\right) + \frac{1+x^2}{(1-x)_+} \log\left(\frac{1-x}{x^2}\right) - (1-x) \right] \\ & \left. + 2(x+z) + \frac{x^2+z^2}{(1-x)_+(1-z)_+} \right\} \quad (7) \end{aligned}$$

For the second subprocess one gets:

$$\begin{aligned} d\Delta\tilde{\sigma}(x, y, z) = & \frac{2\pi\alpha_s^2}{Q^2} e_i^2 \frac{1 - (1-y)^2}{y} \frac{2\alpha_s}{3\pi} \left\{ \delta(1-x) \left[\frac{1 + (1-z)^2}{z} \right. \right. \\ & \left. \log\left(\frac{Q^2 z}{m_g^2}\right) - z \right] + 2(x + (1-z)) + \frac{x^2 + (1-z)^2}{(1-x)_+ z} \left. \right\} \quad (8) \end{aligned}$$

One similarly obtains for the subprocess involving a polarized gluon:

$$d\Delta\tilde{\sigma}(x, y, z) = \frac{2\pi\alpha_s^2}{Q^2} \frac{\alpha_s}{2\pi} \frac{1-(1-y)^2}{y} \left\{ \delta(1-z) \frac{1}{2} (2x-1) \right. \\ \left. \left(\log \frac{Q^2}{m_g^2 x^2} - 1 \right) - x + \frac{2x-1}{2z} + \frac{1-2z(1-x)}{2(1-z)_+} \right\} \quad (9)$$

In order to cancel the mass singularities we have to compute the first-order corrections to the polarized structure functions. For DIS on polarized quarks next-to-leading corrections take the form:

$$\delta\Delta q_i^H(x_H, Q^2) = e_i^2 \int_{x_H}^1 \frac{dx}{x} \Delta P_i(x, Q^2) \Delta q_i^H\left(\frac{x_H}{x}\right) \quad (10)$$

where

$$\Delta P_i(x, Q^2) = \frac{2\alpha_s}{3\pi} \left\{ 3x - \frac{3}{2(1-x)_+} - \delta(1-x) \left(\frac{9}{4} + \frac{2\pi^2}{3} \right) \right. \\ \left. + \left(\frac{1+x^2}{1-x} \right)_+ \log\left(\frac{Q^2}{m_g^2}\right) + \frac{1+x^2}{(1-x)_+} \log\left(\frac{1-x}{x^2}\right) \right\} \quad (11)$$

For DIS scattering on polarized gluons we get

$$\delta\Delta G(x_H, Q^2) = \int_{x_H}^1 \frac{dx}{x} \Delta P_g(x, Q^2) \Delta G\left(\frac{x_H}{x}\right) \quad (12)$$

where

$$\Delta P_g(x, Q^2) = \frac{\alpha_s}{4\pi} (2x-1) \left[\log\left(\frac{Q^2}{m_g^2 x^2}\right) - 2 \right] \quad (13)$$

These results are in agreement with those of Ref. 11) using a dimensional regularization for infra-red divergences. The next-to-leading corrections to the fragmentation functions read:

$$\delta D_i^{H'}(z_{H'}, Q^2) = \int_{z_{H'}}^1 \frac{dz}{z} F_i(z, Q^2) D_i^{H'}\left(\frac{z_{H'}}{z}\right) \quad (14)$$

with

$$F_i(z, Q^2) = \frac{2\alpha_s}{3\pi} \left[\delta(1-z) \left(\frac{\pi^2}{3} - \frac{9}{4} \right) + \left(\frac{1+z^2}{1-z} \right)_+ \log\left(\frac{Q^2}{m_g^2}\right) \right. \\ \left. + \frac{1+z^2}{(1-z)_+} \log z(1-z) - \frac{3}{2(1-z)_+} + \frac{z+1}{2} \right] \quad (15)$$

for the quark fragmentation and

$$\mathcal{S} D_g^{H'}(z_{H'}, Q^2) = \int_{z_{H'}}^1 \frac{dz}{z} F_g(z, Q^2) D_g^{H'}\left(\frac{z_{H'}}{z}\right) \quad (16)$$

with

$$F_g(z, Q^2) = \frac{2\alpha_s}{3\pi} \left[\frac{1+(1-z)^2}{z} \log\left(\frac{Q^2}{m_g^2}\right) + \frac{1+(1-z)^2}{z} \log z^2 - 2z \right] \quad (17)$$

for the gluon fragmentation.

As a check of the validity of our results we recover the electroproduction sum rules by integration over the z variable and for the polarized quark structure functions the $(1-(\alpha_s/\pi))$ first-order correction to the Bjorken sum rule. One finally gets:

$$\begin{aligned} \frac{d\Delta\sigma^q}{dx_H dy dz_{H'}} &= \frac{2\pi\alpha^2}{Q^2} \left(\frac{1-(1-y)^2}{y} \right) \sum_i e_i^2 \left\{ \Delta q_i^H(x_H, Q^2) D_i^{H'}(z_{H'}, Q^2) \right. \\ &\quad \left[1 + \frac{2\alpha_s}{3\pi} \left(\frac{5}{2} - \frac{\pi^2}{3} \right) \right] + \frac{2\alpha_s}{3\pi} \left\{ \Delta q_i^H(x_H, Q^2) \int_{z_{H'}}^1 \frac{dz}{z} D_i^{H'}\left(\frac{z_{H'}}{z}, Q^2\right) \right. \\ &\quad \left[\frac{1+z^2}{(1-z)_+} \log \frac{1}{z} + \frac{3}{2(1-z)_+} + \frac{(z-3)}{2} \right] + D_i^{H'}(z_{H'}, Q^2) \int_{x_H}^1 \frac{dx}{x} \\ &\quad \Delta q_i^H\left(\frac{x_H}{x}, Q^2\right) \left(\frac{3}{2(1-x)_+} - 2x - 1 \right) + \int_{z_{H'}}^1 \frac{dz}{z} D_i^{H'}\left(\frac{z_{H'}}{z}, Q^2\right) \int_{x_H}^1 \frac{dx}{x} \\ &\quad \Delta q_i^H\left(\frac{x_H}{x}, Q^2\right) \left(2x + 2z + \frac{x^2+z^2}{(1-x)_+(1-z)_+} \right) + \Delta q_i^H(x_H, Q^2) \int_{z_{H'}}^1 \frac{dz}{z} \\ &\quad D_g^{H'}\left(\frac{z_{H'}}{z}, Q^2\right) \left(\frac{1+(1-z)^2}{z} \log \frac{1}{z} + z \right) + \int_{x_H}^1 \frac{dx}{x} \Delta q_i^H\left(\frac{x_H}{x}, Q^2\right) \int_{z_{H'}}^1 \frac{dz}{z} \\ &\quad \left. D_g^{H'}\left(\frac{z_{H'}}{z}, Q^2\right) \left[2x + 2(1-z) + \frac{x^2+(1-z)^2}{(1-x)_+ z} \right] \right\} \quad (18) \end{aligned}$$

for initial polarized quark and

$$\begin{aligned} \frac{d\Delta\sigma^g}{dx_H dy dz_{H'}} &= \frac{2\pi\alpha^2}{Q^2} \frac{1-(1-y)^2}{y} \sum_i e_i^2 \left\{ \Delta G(x_H, Q^2) D_i^{H'}(z_{H'}, Q^2) \right. \\ &\quad \left. + \frac{\alpha_s}{2\pi} D_i^{H'}(z_{H'}, Q^2) \int_{x_H}^1 \frac{dx}{x} \Delta G\left(\frac{x_H}{x}, Q^2\right) \left(\frac{2x-1}{2} \right) \right\} \quad (19) \end{aligned}$$

$$+ \left\{ \int_{z_H'}^1 \frac{dz}{z} D_i^{H'}\left(\frac{z_H'}{z}, Q^2\right) \int_{x_H}^1 \frac{dx}{x} \Delta G\left(\frac{x_H}{x}, Q^2\right) \left(\frac{2x-1}{2}\right) \left(\frac{2z^2-2z+1}{z(1-z)}\right) \right\} \quad (19)$$

cont.

A comment is in order here. In Eq. (18) we recognize the infra-red singular factor

$$\frac{2\alpha_s}{3\pi} \frac{x^2 + z^2}{(1-x)_+ (1-z)_+} \quad (20)$$

When convoluted with the quark and fragmentation functions it leads to the term:

$$\frac{4\alpha_s}{3\pi} \ln(1-x_H) \ln(1-z_{H'}) \sum_i e_i^2 \Delta q_i^H(x_H, Q^2) D_i^{H'}(z_{H'}, Q^2) \quad (21)$$

Its exponentiation can be done as in Ref. 12) for lepton pair production. One has simply to replace the term

$$\left[1 + \frac{4\alpha_s}{3\pi} \ln(1-x_H) \ln(1-z_{H'}) \right] \sum_i e_i^2 \Delta q_i^H(x_H, Q^2) D_i^{H'}(z_{H'}, Q^2) \quad (22)$$

by

$$\begin{aligned} & \exp - \frac{12\pi}{(33-2N_f)} \frac{4}{3} \left[\ln\left(\frac{Q^2(1-x_H)}{\Lambda^2}\right) \ln\left(\frac{\alpha_s(Q^2)}{\alpha_s(Q^2(1-x_H))}\right) \right. \\ & + \ln\left(\frac{Q^2(1-z_{H'})}{\Lambda^2}\right) \ln\left(\frac{\alpha_s(Q^2)}{\alpha_s(Q^2(1-z_{H'}))}\right) - \ln\left(\frac{Q^2(1-x_H)(1-z_{H'})}{\Lambda^2}\right) \ln\left(\frac{\alpha_s(Q^2)}{\alpha_s(Q^2(1-x_H)(1-z_{H'}))}\right) \\ & \left. \cdot \sum_i e_i^2 \Delta q_i^H(x_H, Q^2) D_i^{H'}(z_{H'}, Q^2) \right] \quad (23) \end{aligned}$$

This exponentiation will give deviations from the first order only in the kinematical domain x_H and $z_{H'} \rightarrow 1$ which is out of present experimental access. Moreover this exponentiated term is the same as for the unpolarized case¹³⁾ so that the effects due to soft gluon resummation will cancel when considering the asymmetry. Therefore they will be neglected throughout this paper.

3. - Numerical results

We now discuss the implications of QCD corrections for polarized SIDI. We will restrict ourselves to the production of charged pions from interactions of muons of momentum $p_{lab} = 200$ GeV/c on a proton or neutron target. We need an

input of the structure and fragmentation functions at a given Q^2 . We have taken for spin dependent structure functions the two sets given in Ref. 5). Let us recall that set I (II) corresponds to a positive (negative) sea polarization. The pion fragmentation functions are parametrized at $Q^2 = 50 \text{ GeV}^2$ as¹⁴⁾:

$$\begin{aligned} z D_U^{\pi^+}(z) &= 0.7 (1-z)^{1.75} \\ z D_U^{\pi^-}(z) &= \left(\frac{1-z}{1+z} \right) z D_U^{\pi^+}(z) \end{aligned} \quad (24)$$

We have checked that our results are not dependent on another choice of fragmentation functions like¹⁵⁾ for $z_H, < 0.6$. We have plotted in Fig. 1 the K factor:

$$K(x_H, y, z_H) = \frac{d\Delta\sigma/dx_H dy dz_H' |_{\text{NEXT TO LEADING}}}{d\Delta\sigma/dx_H dy dz_H' |_{\text{LEADING}}} \quad (25)$$

as a function of x_H for two z_H , values (0.1 and 0.55) in the case of positive pions as an illustration. For small z_H , values K is slowly decreasing around 1 indicating that next-to-leading corrections are very small. For large z_H , values deviations from the leading order are more pronounced since $K(x_H, Q^2, z_H)$ increases from 0.8 to 1.1 at large x_H . This factorization breaking is very similar to the unpolarized case⁸⁾.

We give in Fig. 2 the asymmetries $A_1^{\pi^\pm}$ [see Eq. (5)] for positive and negative pions from a proton target at $Q^2 = 20 \text{ GeV}^2$ as a function of x_H for several z_H , values (0.1, 0.3 and 0.55). Predictions are given for the set II of polarized structure functions corresponding to a negative polarization of the sea. $A_1^{\pi^+}$ has a sizeable magnitude at small x_H reaching 30% for $z_H = 0.55$. $A_1^{\pi^-}$ is smaller reaching roughly 20% for small z_H . Both asymmetries decrease with x_H . This effect is mainly due to the factor $K_{\text{cin}} = (1-(1-y)^2)/(1+(1-y)^2)$ since for $x_H \rightarrow 1$, $y \rightarrow Q^2/s$ which is very small in the kinematical range explored. The z_H behaviour is different. Since at fixed x_H and neglecting the sea $A_1^{\pi^+}|_{z_H=0} \sim K_{\text{cin}} (4\Delta U_V + \Delta d_V/4U_V + d_V)$ for z_H , and $A_1^{\pi^+}|_{z_H=1} \sim K_{\text{cin}} (\Delta U_V/U_V)$ for $z_H = 1$ $A_1^{\pi^+}|_{z_H=1}$ is larger than $A_1^{\pi^+}|_{z_H=0}$ (since Δd_V is negative) and therefore $A_1^{\pi^+}$ increases with z_H . $A_1^{\pi^-}$ decreases with z_H , since $A_1^{\pi^-}|_{z_H=0} \sim K_{\text{cin}} (4\Delta U_V + \Delta d_V/4U_V + d_V)$ and $A_1^{\pi^-}|_{z_H=1} \sim K_{\text{cin}} (\Delta d_V/d_V)$. As shown by the shaded area, we are sensitive to the set of polarized structure functions in the low x_H regime: $A_1^{\pi^+}$ increases by roughly 15%

at $x_H = 0.05$ if one uses set I. This result is directly related to the inclusive spin asymmetry A_1^P since a negative polarized sea (set II) led to smallest values for $x_H < 0.4$ ⁵⁾. We have plotted in Fig. 3 $A_1^{\pi^+}$ at $Q^2 = 50 \text{ GeV}^2$. The asymmetry is larger in magnitude and still decreasing with x_H . These effects are due to the kinematical factor K_{cin} .

We give now in Fig. 4 the asymmetries $A_1^{\pi^\pm}$ from a neutron target. They are smaller in magnitude compared to the proton case. They increase with x_H at fixed z_H , since for small x_H we are sensitive to $(\Delta d_V + \Delta \bar{q}) / (d_V + \bar{q})$ whereas for large x_H we are sensitive to $\Delta U_V / U_V$. $A_1^{\pi^+}$ is positive except for small z_H , and x_H . $A_1^{\pi^-}$ is always negative. For higher Q^2 , the magnitude of $A_1^{\pi^+}$ increases and $A_1^{\pi^-}$ is still more negative for large z_H values.

We will finally consider the charge ratios:

$$R^P = \frac{d\Delta\sigma / dx_H dy dz_{H'} \big|_{H'=\pi^+}}{d\Delta\sigma / dx_H dy dz_{H'} \big|_{H'=\pi^-}} \quad (26)$$

for the polarized case and

$$R^{NP} = \frac{d\sigma / dx_H dy dz_{H'} \big|_{H'=\pi^+}}{d\sigma / dx_H dy dz_{H'} \big|_{H'=\pi^-}} \quad (27)$$

for the unpolarized case. R^P and R^{NP} obtained from a proton target are plotted in Fig. 5 at $Q^2 = 20 \text{ GeV}^2$ as a function of z_H , for two x_H values (0.15 and 0.5). R^P is always larger than R^{NP} . For small x_H we are sensitive to the choice of polarized structure functions: as an illustration, the charged ratio R^P for $z_H = 0.5$ predicted from a negative polarized sea (set II) is roughly 15% higher than the ratio obtained from set I (dashed area). For larger x_H R^P decreases whereas R^{NP} increases since for $x_H \rightarrow 1$ we expect $R^P \sim R^{NP} \sim D_U^{\pi^+} / D_U^{\pi^-}$. We give in Fig. 6 the charge ratio from a neutron target for $x_H = 0.4$. For an unpolarized target R^{NP} is quite independent of x_H and z_H , and roughly equal to one. Since $A_1^{\pi^-}$ changes sign we predict a huge variation of R^P around the z_H value whose $A_1^{\pi^-}$ is zero.

4. - CONCLUSIONS

We have shown that sizeable spin-dependent effects can be found in semi-inclusive deep inelastic scattering. The largest asymmetry from a proton target is obtained for positive pions (reaching 30% for small x_H values at $Q^2 = 20 \text{ GeV}^2$). For a neutron target the asymmetry for negative pions changes sign at a given x_H value which increases with $z_{H'}$. The polarized charge pion ratio R^P is enhanced from the unpolarized case. R^P from a polarized neutron target exhibits strong variations in the kinematical range where the polarized asymmetry for negative pions is zero.

As for the unpolarized case, next-to-leading corrections break the factorization. This effect is small in magnitude and therefore will be difficult to establish. The polarized observables we have considered are also sensitive at small x_H to the sign of the polarized sea. Here also the effect is tiny and should require accurate measurements in order to be conclusive.

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FIGURE CAPTIONS

- Fig. 1 x_H dependence of the K factor at $Q^2 = 20 \text{ GeV}^2$. Full curve: $z_H = 0.55$; dashed curve: $z_H = 0.1$.
- Fig. 2 x_H dependence of the pion asymmetry at $Q^2 = 20 \text{ GeV}^2$ from a proton target for several z_H values. Full curves: leading predictions for negative pions. Dashed curves: next-to-leading corrections. The shaded area gives the uncertainties coming from the parton densities.
- Fig. 3 x_H dependence of the pion asymmetry at $Q^2 = 50 \text{ GeV}^2$ from a proton target for two z_H values. Full curves: next-to-leading predictions for positive pions using set II of polarized structure functions. Dashed curves: next-to-leading predictions for positive pions using set I. Dotted-dashed curves: next-to-leading corrections for negative pions using set I.
- Fig. 4 The same as Fig. 2 from a neutron target.
- Fig. 5 z_H dependence of the pion charge ratio from a proton target at $Q^2 = 20 \text{ GeV}^2$ for two x_H values. Full curves: leading predictions for R^P using set II. Dashed curve: leading prediction for R^P using set I. Dotted-dashed curves: leading prediction for R^{NP} .
- Fig. 6 z_H dependence of the pion charge ratio from a neutron target at $Q^2 = 20 \text{ GeV}^2$ for $x_H = 0.4$. Full curves: leading prediction for R^P using set II. Dashed curves: next-to-leading corrections for R^P . Dotted-dashed curve: R^{NP} . The shaded area gives the uncertainties coming from the parton densities.

Fig. 1

Fig. 2

Fig. 3

Fig. 4

Fig. 5

Fig. 6