Polarized Electroproduction of Charm

J. Smith[a*]

[a]Institute for Theoretical Physics, SUNY at Stony Brook ,
Stony Brook, NY 11794-3840, USA

We discuss the NLO corrections to polarized electroproduction of charmed quarks.

Deep inelastic electroproduction of charm quarks has yielded important information about the gluon density inside the proton (see [1]). Until recently almost all attention was paid to unpolarized charm quark electroproduction. The complete next-to-leading order (NLO) cross-section for this reaction is known [2]. Computer programs are available for the structure functions $F_{2,c}(x,Q^2,m^2)$ and $F_{L,c}(x,Q^2,m^2)$ where x denotes the Bjorken scaling variable, Q^2 ($Q^2 = -q^2 \geq 0$) the virtuality of the photon exchanged between the lepton and the hadron and m the mass of the heavy quark. Analytic expressions for the heavy quark coefficient functions in the large Q^2 limit are also available in [3]. These calculations give an excellent description of the HERA data on $D^{*\pm}$ production.

Now there is a great deal of interest in charm production in polarized deep inelastic lepton-hadron scattering (see articles in these proceedings). One is especially interested in extracting the polarized gluon density since it plays an important role in the description of the longitudinal spin structure function $g_1(x,Q^2)$. The heavy flavor contribution $g_1(x,Q^2,m^2)$ can be expressed in terms of convolution integrals over the <u>partonic</u> scaling variable $z = Q^2/(s+Q^2)$ where s is the square of the photon-parton centre-of-mass energy ($s \geq 4m^2$). This yields the following result

$$g_1(x,Q^2,m^2) = \frac{1}{2n_f}\sum_{k=1}^{n_f} e_k^2 \int_x^{z_{\rm max}} \frac{dz}{z}$$
$$\times\Big[\Sigma\Big(\frac{x}{z},\mu^2\Big) L_q^{\rm S}\Big(z,\frac{Q^2}{m^2},\frac{m^2}{\mu^2}\Big)$$
$$+G\Big(\frac{x}{z},\mu^2\Big) L_g^{\rm S}\Big(z,\frac{Q^2}{m^2},\frac{m^2}{\mu^2}\Big)$$
$$+n_f\Delta_k\Big(\frac{x}{z},\mu^2\Big) L_q^{\rm NS}\Big(z,\frac{Q^2}{m^2},\frac{m^2}{\mu^2}\Big)\Big]$$
$$+\frac{1}{2}e_H^2 \int_x^{z_{\rm max}} \frac{dz}{z}\Big\{\Sigma\Big(\frac{x}{z},\mu^2\Big)$$
$$\times H_q^{\rm PS}\Big(z,\frac{Q^2}{m^2},\frac{m^2}{\mu^2}\Big)$$
$$+G\Big(\frac{x}{z},\mu^2\Big) H_g^{\rm S}\Big(z,\frac{Q^2}{m^2},\frac{m^2}{\mu^2}\Big)\Big\}\,, \quad (1)$$

where $z_{\rm max} = Q^2/(4m^2+Q^2)$. The function $G(z,\mu^2)$ represents the <u>polarized</u> gluon density, while $\Sigma(z,\mu^2)$ and $\Delta_k(z,\mu^2)$ represent the singlet and non-singlet <u>polarized</u> quark densities. In the above expression the charges of the light quarks and the heavy quark are denoted by e_i and e_H respectively, in units of e. Furthermore n_f stands for the number of light quarks and μ denotes the mass factorization scale μ_F, which is chosen equal to the renormalization scale μ_R. Both polarized coefficient functions L_k and H_k, which are the differences of the amplitudes for polarized photon scattering from left and right handed quarks and/or left or right handed gluons, can be expanded in a power series in $\alpha_s/4\pi$ ($\alpha_s = g^2/4\pi$) as follows

$$L_k = \sum_{n=2}^{\infty} \frac{\alpha_s}{4\pi} L_k^{(n)}\,, H_k = \sum_{n=1}^{\infty} \frac{\alpha_s}{4\pi} H_k^{(n)}\,. \quad (2)$$

There are several processes which contribute to the spin dependent heavy quark coefficient functions up to order $(\alpha_s/4\pi)^2$. The lowest order (LO) coefficient function $H_g^{\rm S,(1)}$ is given by the polarized Born photon-gluon fusion process

*Work supported in part by NSF-9722101

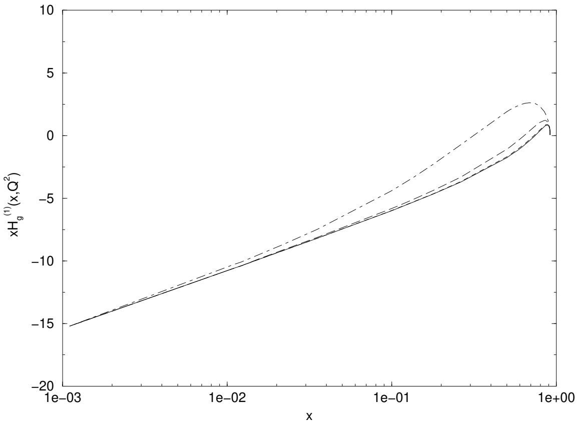

Figure 1. $xH_g^{(1)}(x,Q^2)$ versus x for fixed Q^2.

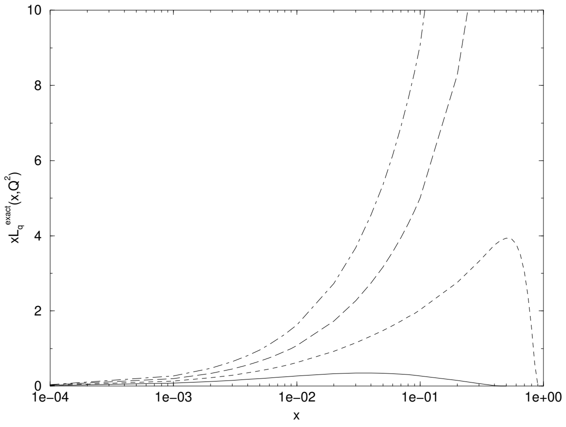

Figure 2. $xL_q^{(2)}(x,Q^2)$ versus x for fixed Q^2.

(see [4] and Fig.1. The curves are for $Q^2 = 1$ $(\mathrm{GeV/c})^2$ (solid), $Q^2 = 10$ $(\mathrm{GeV/c})^2$ (short-dashed), $Q^2 = 100$ $(\mathrm{GeV/c})^2$ (long-dashed), and $Q^2 = 1000$ $(\mathrm{GeV/c})^2$ (dot-dashed), while $m = 4.5$ GeV/c^2). The first moment of this function vanishes because it has both positive and negative parts. This explains the oscillatory behaviour of the LO expression for $g_1(x,Q^2,m^2)$. This same sum rule holds for the higher order coefficient $H_g^{\mathrm{S},(2)}$ and its mass factorization part $\bar{H}_g^{\mathrm{S},(2)}$ (the coefficient of the $\ln(m^2/\mu^2)$ term). However it does not hold for $H_q^{\mathrm{PS},(2)}$ or its mass factorization part $\bar{H}_q^{\mathrm{PS},(2)}$.

Recently [5] an analytic result was derived for one of the photon-quark coefficient functions corresponding to the reaction $\gamma^* + q \to c + \bar{c} + q$ (Compton process, $L_q^{\mathrm{NS},(2)}$, finite in four dimensions and shown in Fig.2 for the same Q^2 and m^2 values as in the previous figure), which shows up for the first time in NLO. Expressions for the other NLO heavy quark coefficient functions are available in the asymptotic limit $Q^2 \gg m^2$ in [5]. Contrary to the unpolarized electroproduction of charm quarks, which is studied at large values of Q^2 (e.g. HERA), most polarized deep inelastic experiments are carried out at small Q^2. Therefore it is necessary to know the NLO heavy quark coefficient functions for general x and Q^2 values.

The calculation of the pQCD corrections reveals the presence of UV, IR and C divergences which have to be regularized. The projection operators needed to calculate $H_g^{\mathrm{S},(2)}$ involve Levi-Civita tensors while those for $H_q^{\mathrm{PS},(2)}$ involve the γ_5 spin projector. We have evaluated the graphs using n-dimensional regularization [6]. In the case of the Levi-Civita tensor we followed the prescription advocated in [7] and contracted the Lorentz indices in n dimensions. This can be achieved because product of these tensors can be written as a determinant containing products of the metric tensors $g_{\mu\nu}$. This method was used in [8] to obtain the two-loop massless partonic coefficient functions corresponding to the structure functions $g_1(x,Q^2)$. We also have to adopt a prescription for the γ_5-matrix appearing in spin projector. We chose the HVBM prescription as presented in [6], [9]. Since only one γ_5-matrix occurs in our analysis the prescription above is equivalent to the replacement [7], [10]

$$\gamma_5 \gamma_\mu k_1^\mu = -\frac{i}{6} \epsilon_{\mu\rho\sigma\nu} k_1^\mu \gamma^\rho \gamma^\sigma \gamma^\nu \,. \tag{3}$$

The γ-matrix algebra can now be carried out in n-dimensions. Note that the NLO processes under consideration here contain massive quarks so that we avoid one collinear pole and thereby problems associated with additional C singularities in the massless limit, which introduce anomalies (see [12]). Since the NLO contributions oscillate we have evaluated the two mass factorization terms

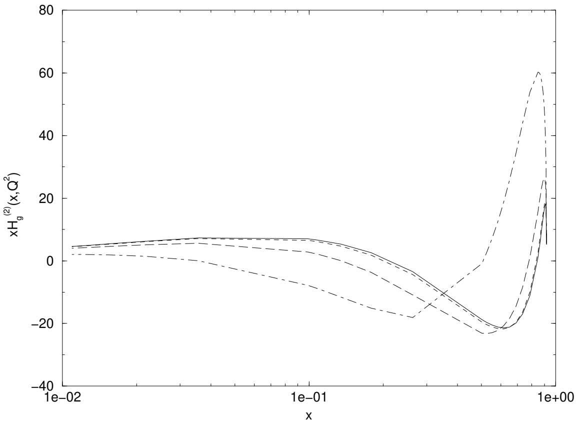

Figure 3. $x\bar{H}_g^{(2)}(x,Q^2)$ versus x for fixed Q^2.

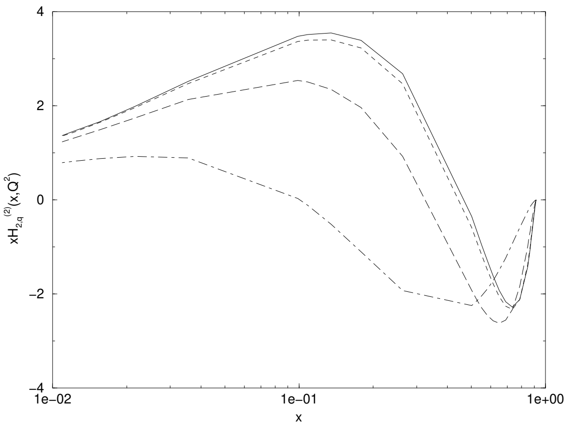

Figure 4. $x\bar{H}_q^{(2)}(x,Q^2)$ versus x for fixed Q^2.

$\bar{H}_g^{\mathrm{PS},(2)}$ and $\bar{H}_q^{\mathrm{S},(2)}$ analytically [13] and they are shown in Figs. 3 and 4 respectively, with the same notation as in the previous figures. The other terms $H_g^{\mathrm{PS},(2)}$ and $H_q^{\mathrm{S},(2)}$ are available in the form of two dimensional integrals. The answers can be checked in the large Q^2 limit with the results in [5] and in the small Q^2 limit with the results in [12]. We are presently making these checks and will publish our results in due course. Due to the ocillatory behaviour of all the polarized coefficient functions we expect that the the final corrections to the longtitudinal structure function $g_1(x,Q^2,m^2)$ in Eq. (1) will be quite small.

REFERENCES

1. J.J. Aubert *et al.* (EMC Collaboration), Nucl. Phys. **B213** (1983) 31; C. Adloff, (H1 Collaboration), these proceedings; I. Redondo, (ZEUS collaboration), these proceedings.
2. E. Laenen, S. Riemersma, J. Smith and W. L. van Neerven, Nucl. Phys. **B392** (1993) 162, 229; S. Riemersma, J. Smith and W. L. van Neerven, Phys. Lett. **B347** (1995) 143; B. W. Harris and J. Smith, Nucl. Phys. **B452** (1995) 109, Phys.Rev. **D57** (1998) 2806.
3. M. Buza, Y. Matiounine, J. Smith, R. Migneron and W. L. van Neerven, Nucl. Phys. **B472** (1996) 611.
4. A. D. Watson, Z. Phys. **C12** (1982) 123; W. Vogelsang, Z. Phys. **C50** (1991) 275; M. Glück, E. Reya and W. Vogelsang. Nucl. Phys. **B351** (1991) 579.
5. M. Buza, Y. Matiounine, J. Smith and W. L. van Neerven, Nucl. Phys. **B485** (1997) 420.
6. G. 't Hooft and M. Veltman, Nucl. Phys. **B44** (1972) 189.
7. S. A. Larin, Phys. Lett. **B303** (1993) 113.
8. E. B. Zijlstra and W. L. van Neerven, Nucl. Phys. **B417** (1994) 61, Erratum ibid. **B426** (1994) 245.
9. P. Breitenlohner and B. Maison, Commun. Math. Phys. **52** (1977) 11; **52** (1977) 39; **52** (1977) 55.
10. D. Akyeampong and R. Delburgo, Nuovo Cimento **A17** (1973) 578; **A18** (1973) 94; **A19** (1974) 219.
11. Y. Matiounine, J. Smith and W. L. van Neerven, Phys. Rev. **D58** (1998) 076002.
12. I. Bojak and M. Stratmann, Phys.Lett.**B433** (1998) 411, hep-ph/9807405; A.P. Contogouris et al. Nucl. Phys. Proc. Suppl. **64** (1998) 117; W.L. van Neerven, Acta Phys. Polon. **B29** (1998) 1175.
13. J. Smith and W.L. van Neerven, to be published.