NLO Corrections to Heavy Quark Production with Polarized Beams

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Abstract

We present a calculation of the NLO QCD corrections to heavy flavor photoproduction with longitudinally polarized beams. We apply our results to study the spin asymmetry for total charm quark production which will be used for a first direct determination of $\Delta g$ by the COMPASS experiment. We also briefly discuss the main theoretical uncertainties inherent in this calculation.
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1 Introduction

Despite significant progress in the field of spin-dependent DIS, the polarized gluon density $\Delta g$ remains almost completely unconstrained by presently available DIS data. An important role plays here the lack of any direct constraint on $\Delta g$ from other processes. Upcoming spin experiments will thus put a special emphasis on exclusive measurements to provide further invaluable information for a more restrictive analysis of polarized parton densities in the future. In this context heavy quark ($Q = c, b$) photoproduction is considered to be one of the best options to pin down $\Delta g$ because in LO only the photon-gluon fusion (PGF) process (an arrow denotes a longitudinally polarized particle)

$$\vec{\gamma} \vec{g} \rightarrow Q \bar{Q}$$

contributes. Since LO estimates of (1) are rather unreliable, and since we already know from the unpolarized NLO calculation that the corrections are sizeable and that the ‘clean picture’ of (1) is obscured by new, light quark induced NLO subprocesses, the knowledge of the polarized NLO corrections is mandatory for a meaningful extraction of $\Delta g$. In what follows we will briefly highlight the main steps and results of our calculation.

2 Technical Framework

The NLO QCD corrections to the PGF mechanism in (1) consist of three parts: (i) the one-loop virtual corrections, (ii) the real corrections $\vec{\gamma} \vec{g} \rightarrow Q \bar{Q} q$, and (iii) a new generic NLO production mechanism $\vec{\gamma} \vec{q} \rightarrow Q \bar{Q} q$ ($\vec{q}$). In the calculation of (i)-(iii) one encounters UV, IR, and mass (M) singularities which are removed by renormalization, in the sum of (i) and (ii), and by factorization, respectively. To make all these singularities manifest we choose to work in the framework of $n$-dimensional regularization.
The required polarized squared matrix elements for (1) and (i)-(iii)

\[ \Delta |M|^2 = \frac{1}{2} \left[ |M|^2 (++) - |M|^2 (+-) \right], \]  

where the ± denote the helicities of the incoming particles, are obtained by projecting onto the helicity states of the bosons (photons or gluons) and quarks using the \( \epsilon_{\mu\nu\rho\sigma} \) tensor and the \( \gamma_5 \) matrix, respectively. By taking the sum instead of the difference in (2) we fully agree with the known unpolarized results and the abelian, \( \gamma^5 \rightarrow Q \bar{Q} \), part of our results agrees as well.

The presence of \( \gamma_5 \) and \( \epsilon_{\mu\nu\rho\sigma} \) in the polarized calculation introduces some complications since these objects have no unique continuation to \( n \neq 4 \). In the HVBM prescription, which we use, the usual \( n \)-dim. scalar products \( k \cdot p \) are accompanied by their \((n-4)\)-dim. subspace counterparts \( \hat{k} \cdot \hat{p} \) (‘hat momenta’). These terms deserve special attention when performing the \( 2 \rightarrow 3 \) phase space integrations. For single-inclusive heavy quark production, which we consider here, only a single hat momenta combination \( \hat{p}^2 \) shows up in \( \Delta |M|^2 \) and the appropriately modified \( 2 \rightarrow 3 \) phase space formula schematically reads

\[
dPS_3 = dPS_{3,\text{unp}}(\theta_1, \theta_2) \times \frac{1}{B \left( \frac{1}{2}, \frac{n-4}{2} \right)} \int_0^1 dx \frac{x^{(n-6)/2}}{\sqrt{1-x}}, \]  

where \( x \equiv 4(s_4 + m^2)\hat{p}^2/(s_4^2 \sin^2 \theta_1 \sin^2 \theta_2) \), \( B \) is the Beta function, \( m \) denotes the heavy quark mass, and \( s_4 \equiv s + \ell_1 + u_1 \). \( \theta_{1,2} \) are introduced to parametrize the momenta of the two not observed partons. However, due to the appearance of \( m \) in \( x \) it turns out that all contributions due to \( \hat{p}^2 \) are at least of \( O(n-4) \) and hence drop out when the limit \( n \rightarrow 4 \) is taken. The remaining phase space integration then proceeds as in the unpolarized case.

Finally, it should be recalled that in the factorization procedure for (iii) one has to introduce the parton content of the polarized photon which is experimentally completely unknown so far. A scheme independent result in \( O(\alpha_s^2\alpha) \) can thus only be obtained for the sum of the ‘direct’ and ‘resolved’ photon contributions. The NLO corrections for the latter are unknown and, for the time being, have to be estimated in LO.

\[ ^a \text{This differs from a calculation involving only massless particles.} \]
3 Numerical Results and Phenomenological Aspects

The total photon-parton cross section in NLO can be expressed in terms of scaling functions ($i = g, q, \bar{q}$)

$$\Delta \hat{\sigma}_{i\gamma}(s, m^2, \mu_f) = \frac{\alpha_s}{m^2} \left[ \Delta f^{(0)}_{i\gamma}(\eta)^2 + 4\pi\alpha_s \left\{ \Delta f^{(1)}_{i\gamma}(\eta) + \Delta \tilde{f}^{(1)}_{i\gamma}(\eta) \ln \frac{\mu_f^2}{m^2} \right\} \right]$$

where $\Delta f^{(0)}_{i\gamma}$ and $\Delta f^{(1)}_{i\gamma}$, $\Delta \tilde{f}^{(1)}_{i\gamma}$ stand for the LO and NLO corrections, respectively, $\mu_f$ denotes the factorization scale (for simplicity we choose $\mu_r = \mu_f$), and $\eta \equiv s/4m^2 - 1$. The scaling functions can be further decomposed depending on the electric charge of the heavy and light quarks, $e_Q$ and $e_q$, respectively:

$$\Delta f_{g\gamma}(\eta) = e_Q^2 \Delta c_{g\gamma}(\eta), \quad \Delta f_{q\gamma}(\eta) = e_Q^2 \Delta c_{q\gamma}(\eta) + e_q^2 \Delta d_{q\gamma}(\eta)$$

with corresponding expressions for the $\Delta \tilde{f}_{i\gamma}$.

In Fig. 1 we present $\Delta c^{(0)}_{g\gamma}$, $\Delta c^{(1)}_{g\gamma}$, and $\Delta c^{(1)}_{q\gamma}$ as a function of $\eta$ in the MS scheme. For the discussion below it is important to notice that $\Delta c_{g\gamma}^{(0)}$ (solid line) changes sign at $\eta \approx 3$. Upon adding the NLO terms, multiplied by a factor $4\pi\alpha_s$ (see Fig. 5 in [4]), the zero is shifted towards $\eta \approx 1$. We also note that for $\eta \lesssim 0.1$ the $\mathcal{O}(\alpha_s)$ corrections dominate over the LO result when we include that factor. Comparing our polarized results with the unpolarized ones (see Fig. 5 in [4]), one observes that for $\eta \to 0$ $\Delta c_{g\gamma} \to c_{g\gamma}$, implying that $|M_{g\gamma}|^2(\pm) \to 0$. On the contrary, for $\eta \to \infty$ the unpolarized NLO coefficients approach a plateau value dominating over the LO result while all polarized coefficients tend to zero here, implying that $|M_{q\gamma}|^2(\mp) \to |M_{q\gamma}|^2(\mp)$. The numerically less important quark coefficients can be found in [7].
Using (4) it is easy to calculate the experimentally relevant spin asymmetry for the total hadronic heavy flavor photoproduction cross section

\[ A_{Q \gamma p}(S_{\gamma p}, m^2, \mu_f) = \frac{\Delta \sigma_{Q \gamma p}(S_{\gamma p}, m^2, \mu_f)}{\sigma_{Q \gamma p}(S_{\gamma p}, m^2, \mu_f)} \]  

(6)
as a function of the photon-proton c.m.s. energy \( S_{\gamma p} \) and where

\[ \Delta \sigma_{Q \gamma p}(S_{\gamma p}, m^2, \mu_f) = \sum_{f=q, q\bar{q}, g} \int \frac{d x \Delta \hat{\sigma}_{f \gamma}(x S_{\gamma p}, m^2, \mu_f) \Delta f_p(x, \mu_f^2)}{x S_{\gamma p}} \]  

(7)
(with a similar expression for the unpolarized cross section \( \sigma_{Q \gamma p}^Q \)).

In Fig. 2 we show the charm asymmetry \( A_{c \gamma p} \) in LO and NLO, using \( m = 1.5 \) GeV, \( \mu_f = 2m \), and three sets of polarized densities \( ^1,^2 \) in the energy region relevant for COMPASS \( ^{12} \) (they will operate at \( \approx 10 \) GeV). The NLO corrections are large, depend strongly on \( \sqrt{S_{\gamma p}} \) and do not cancel in the ratio (6) as one may naively expect. However, their origin is readily explained. For GRSV \( ^1 \) and GS (A) \( ^2 \) and \( \sqrt{S_{\gamma p}} \geq 12 \) GeV the corrections stem from the shift of the zero in the gluonic coefficient function. \( A_{c \gamma p} \) changes sign at some \( \sqrt{S_{\gamma p}} > 20 \) GeV and large NLO corrections in the vicinity of a zero are natural. For \( \sqrt{S_{\gamma p}} \lesssim 12 \) GeV, where one probes \( \Delta g \) at \( x \gtrsim 0.1 \), the corrections are, on the other hand, entirely due to the badly constrained \( \Delta g \), more precisely, due to the large differences between the LO and NLO \( \Delta g \) in the two sets \( ^1,^2 \). This is illustrated by the dotted curve where we use the NLO GRSV gluon \( ^1 \) to calculate the LO \( A_{c \gamma p} \). In this energy region the observed large corrections

\[
\begin{figure}
\centering
\includegraphics[width=\textwidth]{figure2.png}
\caption{\( A_{c \gamma p} \) in LO and NLO for three sets of polarized parton densities \( ^1,^2 \). \( \sigma_{c \gamma p} \) in (6) was calculated using the GRV \( ^{11} \) densities. The dotted line shows \( A_{c \gamma p} \) in LO using a NLO \( \Delta g \) (see text). The bar shows the uncertainty for such a measurement at COMPASS \( ^{12} \).}
\end{figure}
\]
thus should not be taken too literally. For the GS (C) set the situation is more involved since here $\Delta g$ oscillates as well.

Finally we briefly discuss the importance of the main theoretical uncertainties. Fig. 3 shows the dependence of $\Delta \sigma_{cp}$ on the choice of $\mu_f$. The improved scale dependence in NLO clearly underlines the importance of our NLO calculation. Light quark induced subprocesses contribute about 5% at $\sqrt{S_{cp}} \simeq 10$ GeV, but more for GS (C), and should be subtracted before extracting $\Delta g$. The ‘resolved’ photon contribution was shown to be small in $\text{LO}$. More important is the unknown value of $m$, leading to shifts of about 30% when $m$ is varied by 0.2 GeV around 1.5 GeV used in our calculations. Finally, for a reliable extraction of $\Delta g$ from $A_{cp}$ our knowledge of the unpolarized gluon has to be improved as well since at large $x$ the uncertainty in $g(x, \mu^2)$ is rather sizeable.

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