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Exclusive photoproduction of J/ψ and Υ at one-loop accuracy in collinear factorisation with GPD evolution and high-energy resummation

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We report on the first complete study at one-loop accuracy of J/ψ and Υ exclusive photoproduction off protons in Collinear Factorisation (CF) including GPD evolution. We present our solution to the NLO perturbative instability of the cross section at high photon-proton centre of mass energies $W_{\gamma p}$: we have matched high-energy factorisation (HEF) which resums higher-order QCD corrections enhanced by a logarithm of the parton energies in the Doubly-Logarithmic Approximation (DLA) to CF. This is particularly relevant at high $W_{\gamma p}$ UPCs at the LHC.

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1 Introduction

Like other hard exclusive reactions (DVCS, DVMP, etc.) the exclusive photoproduction of J/ψ and Υ (denoted $\mathcal{Q}$) probe the 3D hadron structure through Generalised Parton Distributions (GPDs) [1–3]. Indeed the process $\gamma(q,\lambda)p_s(p) \to \mathcal{Q}(p_\mathcal{Q},\lambda')p'_{s'}(p')$ is expected to admit at small $t=(p'-p)^2$, even for $q^2 \to 0$, a collinear factorisation [4] at the level of the amplitude, $\mathcal{A}^{\lambda\lambda'}_{ss'}$, as a convolution of a process-dependent perturbative hard-scattering coefficient function, C_i and a universal, non-perturbative quark and/or gluon GPD correlator $F_{i,ss'}$ $(i=q,g)$ as

$$\mathcal{A}^{\lambda\lambda'}_{ss'} = -\varepsilon^{\gamma}_{\lambda}.\varepsilon^{*\mathcal{Q}}_{\lambda'} \sum_{i=q,g} \int_{-1}^{1} \frac{dx}{x^{1+\delta_{ig}}} C_i\left(x,\xi;\mu_F,\mu_R\right) F_{i,ss'}(x,\xi,t;\mu_F), \tag{1}$$

where $\varepsilon^{\gamma}_{\lambda}$ is the photon-polarisation vector with helicity λ in the gauge $\varepsilon^{\gamma}_{\lambda}.p=0$ and $\varepsilon^{\mathcal{Q}}_{\lambda'}$ that of the quarkonium, $\xi=-\Delta^+/2P^+$ $(P=(p+p')/2)$ is the fraction of the longitudinal proton momentum transfer, or skewness, with $\Delta=(p'-p)$ being the momentum transfer. The light-cone components are here defined as $k^{\pm} \equiv (k^0 \pm k^3)/\sqrt{2}$ and $\mathbf{k}_\perp \equiv (k^1,k^2)$ in the γp centre-of-momentum frame, with the z axis defined along the proton direction. $\mu_{R,F}$ are respectively the renormalisation (factorisation) scales.

Such a factorisation is not proven to all orders in α_s but explicit one-loop computations in the non-relativistic static limit have shown that it holds up to NLO in α_s [4]. In such a case, the hard scale is provided by the heavy-quark mass m_Q. Note that factorisation also holds at NLO for the case of finite q^2 [5]. One interesting feature of the CF amplitude of Eq. (1) is that the t dependence uniquely follows from that of GPD F_i. As usual, the $\gamma p \to \mathcal{Q}p$ two-body t-differential cross section is readily obtained from the amplitude Eq. (1) via

$$\frac{d\sigma}{dt} = \frac{1}{16\pi(W^2_{\gamma p}-m_p^2)^2} \overline{\sum_{\lambda,\lambda'=\pm}} \overline{\sum_{s,s'}} |\mathcal{A}^{\lambda\lambda'}_{ss'}|^2, \tag{2}$$

where the sums include the initial-state averaging over the incoming proton spins (s,s') and the photon transverse helicities $(\pm)$. Gluon GPDs which give the largest contribution to exclusive quarkonium photoproduction are poorly known, in particular their t dependence which is unconstrained. For this reason, we have decided to focus on the small t region close to $t=t_{\min}$. In practice it is well approximated by taking $t=0$ at large $W_{\gamma p}$ values which are within the reach of existing HERA and LHC UPC data as well as future one at the EIC. In fact, HERA data [6] are so precise that even for $t \simeq 0$ and for various $W_{\gamma p}$, useful data-theory comparisons can be made. In the case of UPC, t-differential data are currently not yet available in pPb collisions for direct comparison with $\gamma p \to \mathcal{Q}p$. A discussion of the t dependence and of the t-integrated cross sections is thus certainly needed but is beyond the scope of our current discussion.

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The presence of x, ξ and t in F illustrate that collinear factorisation is in terms of 3D GPDs and not in terms of the usual 1D PDFs. In addition, we note the presence of μ_F both in the hard coefficient and the GPDs which are subject to evolution equations.

In the light-cone gauge and at leading twist, F_q and F_g are respectively Fourier transforms of matrix elements of chiral-even operators constructed from quark fields ψ^q or gluon-field-strength tensors $F^{\mu\nu}$ as follows (with $y^+ = 0$ and $\mathbf{y}_\perp = \mathbf{0}$):

$$
\begin{aligned}
F_{q,ss'} &= \frac{1}{2}\int \frac{\mathrm{d}y^-}{2\pi} e^{ixP^+y^-} \langle p', s'|\bar{\psi}^q\left(\tfrac{-y}{2}\right)\gamma^+\psi^q\left(\tfrac{y}{2}\right)|p,s\rangle, \\
F_{g,ss'} &= \frac{1}{P^+}\int \frac{\mathrm{d}y^-}{2\pi} e^{ixP^+y^-} \langle p', s'|F^{+\mu}\left(\tfrac{-y}{2}\right)F_\mu{}^+\left(\tfrac{y}{2}\right)|p,s\rangle.
\end{aligned}
\tag{3}
$$

We parametrise them through known Lorentz structures and twist-2 parton-helicity conserving GPDs $H^{q,g}(x,\xi,t;\mu_F)$ and $E^{q,g}(x,\xi,t;\mu_F)$, as follows ($j = q, g$)

$$
F_{j,ss'} = \frac{1}{2P^+}\left[\bar{u}_{s'}(p')\left(H_j\gamma^+ + E_j\frac{i\sigma^{+\Delta}}{2m_p}\right)u_s(p)\right]. \tag{4}
$$

While the LO cross section for exclusive $\mathcal{Q}$ photoproduction uniquely depends on gluon GPDs, a class of NLO corrections is sensitive to quark GPDs. Explicit μ_R and μ_F scale dependences also arise in the hard coefficient function and are supposed to reduce the scale dependence from respectively the running of $\alpha_s(\mu_R)$ and the evolution of the GPD $H^g(x,\xi,t;\mu_F)$.

Such a decrease of the μ_F dependence was nevertheless not observed in the existing phenomenological studies for $W_{\gamma p} \gg m_Q$ [4, 7, 8] as the NLO cross sections get dramatically sensitive to μ_F. This is in fact reminiscent of several inclusive quarkonium observables [9–11] where negative p_T-integrated cross section were observed along with a very strong μ_F dependence. The solution found to solve this issue in the inclusive case was either to fix μ_F to reduce these anomalously large NLO corrections [12, 13] or to resum the high-energy logarithms generating these perturbative instabilities with High-Energy Factorisation (HEF) [14, 15]. In [16], our aim was to show that such a resummation can also cure these instabilities in exclusive reactions. Thereby, we could go beyond the scale-fixing criterion advocated in [8]. We could also perform the first complete phenomenological study of this exclusive process with GPD evolution and compare it to precise HERA data.

2 Reminder on GPDs and their evolution

In the forward limit $\xi \to 0$ and $t \to 0$, the GPDs should coincide with the corresponding PDFs, namely $H_g(\pm x, 0; \mu_F) = xg(x;\mu_F)$ and $H_q(x,0;\mu_F) = q(x;\mu_F)$, $H_q(-x,0;\mu_F) = -\bar{q}(x;\mu_F)$ for $x > 0$. In addition, Lorentz invariance (see e.g. Ref. [17]) imposes that x moments of GPDs are specific polynomials in ξ. This is usually referred to as the polynomiality of GPDs. Any proper modelling of GPDs should satisfy both these properties. It is important to note that scale evolution preserves such properties.

This is why one usually resorts to a construction of GPDs at the initial scale by performing a Radon transform of objects called double distributions [1], $f_i(\beta,\alpha)$, where $i = q, g$:

$$
H_i(x,\xi;\mu_0) = \int_{-1}^{1} d\beta \int_{-1+|\beta|}^{1-|\beta|} d\alpha\, \delta(\beta + \xi\alpha - x) f_i(\beta,\alpha;\mu_0). \tag{5}
$$

Provided that $f_i(\beta,\alpha)$ are even functions of α, GPD polynomiality is automatically satisfied.

The forward limit constraint can nicely be enforced by a further factorisation Ansatz [18, 19]:

$$
f_i(\beta,\alpha;\mu_0) = h_i(\beta,\alpha) \times \left\{
\begin{array}{cc}
|\beta| g(|\beta|;\mu_0) & \text{for } i = g, \\
\theta(\beta) q_{\text{val}}(|\beta|;\mu_0) & \text{for valence } q, \\
\text{sgn}(\beta) q_{\text{sea}}(|\beta|;\mu_0) & \text{for sea } q,
\end{array}\right. \tag{6}
$$

where $g(x)$ is the gluon PDF, $q_{\text{val}}(x)$ and $q_{\text{sea}}(x)$ are the valence and sea components of quark PDFs where $h_i(\beta,\alpha)$ is the so-called profile function satisfying

$$
\int_{-1+|\beta|}^{1-|\beta|} d\alpha\, h_i(\beta,\alpha) = 1, \tag{7}
$$

as originally proposed in [20, 21].

The choice of $h_i(\beta,\alpha)$ is where the ξ dependence can be modelled. Lacking further constraints, one usually takes [22]:

$$
h_i(\beta,\alpha) \equiv \frac{\Gamma(2n_i+2)}{2^{2n_i+1}\Gamma^2(n_i+1)} \frac{\left((1-|\beta|)^2 - \alpha^2\right)^{n_i}}{(1-|\beta|)^{2n_i+1}}. \tag{8}
$$

The values of n_i, which are flavour specific, control the width of the profile function and the resulting ξ dependence. An extreme case is when $n \to \infty$, $h_q(\beta,\alpha)|_{n\to\infty} = \delta(\alpha)$, hence $H_q(x,\xi;\mu_0)|_{n\to\infty} = q(x;\mu_0)$. The ξ dependence thus disappears. We stress that a different parametrisation of h_i than that of Eq. (8) can be used, especially since such a choice lacks flexibility and certainly does not span all the allowed interplay between x and ξ.

There are no constraints on the t dependence apart from the limit $t \to 0$. It can be easily included for instance by extending $h_i(\beta, \alpha)$ to $h_i(\beta, \alpha, t)$, with $\lim_{t\to 0} h_i(\beta, \alpha, t) = h_i(\beta, \alpha)$.

We have used the Goloskokov-Kroll (GK) DD model [23], i.e. a simple exponential dependence in t : $h_i(\beta, \alpha, t) \equiv e^{bt} h_i(\beta, \alpha)$ with $n_g = n_q^{\rm sea} = 2$ and $n_q^{\rm val} = 1$. We have used the central set of the CTEQ6M PDF set [24] as in the original GK GPD as well as of CT18NLO [25] and other PDFs for more up-to-date predictions. We have also focused on the differential cross section at the minimum value of $|t|$, $t_{\rm min} = -4\xi^2 m_p^2/(1-\xi^2)$. In the limit of large $W_{\gamma p}$ that we focused on, $\xi \simeq M_{\mathcal{Q}}^2/(2W_{\gamma p}^2)$ and $t_{\rm min} \simeq -m_p^2 M_{\mathcal{Q}}^4/W_{\gamma p}^4$. Such a value is much smaller than the the other hadronic scales and we can skip modelling the t dependence of GPDs.

We have performed full leading-logarithmic (LL) GPD evolution with kernels at one-loop [26]. We have used a DD at $\mu_0 = 2$ GeV as our initial condition for the GPD evolution. This was achieved with the `PARTONS` framework [27] interfaced to the `APFEL++` code [28].

Besides its lack of flexibility, modelling GPDs with DDs does not allow one to include contributions with the largest ξ power usually modelled by the so-called D term [29]. In the limit of large $W_{\gamma p}$, or small ξ, these are, however, not expected to be important.

3 NLO instabilities of cross sections in Collinear Factorisation

In the collinearly-factorised expression of the amplitude, Eq. (1), the integration over $x \in [-1, 1]$ spans two physically distinct regions: the so-called DGLAP region (with $\xi < |x| < 1$) and the ERBL region (with $|x| < \xi$). In these, the GPDs obey two different evolution equations [26].

Let us start with the LO results in α_s (and in the squared relative velocity of the heavy quarks, v^2), where only the gluon coefficient function is nonzero: $C_g^{(0)}(x, \xi; \mu_R) = x^2 c(x + \xi - i\varepsilon)^{-1}(x - \xi + i\varepsilon)^{-1}$, where $c = (4\pi\alpha_s(\mu_R) e e_Q R_{\mathcal{Q}}(0))/(m_Q^{3/2}\sqrt{2\pi N_c})$ with $R_{\mathcal{Q}}(0)$ being the value at the origin of the spatial radial wave function of the quarkonium $\mathcal{Q}$, e_Q is the electric charge of the heavy quark Q in units of the positron charge and e is the positron charge in natural units.

The imaginary part of the LO coefficient function, $\mathrm{Im}\, C_g^{\rm LO}(\xi/x) = -\pi c/2\,[\delta(\xi/x - 1) + \delta(\xi/x + 1)]$ leads to the imaginary part of the amplitude at LO in α_s, proportional to $H_g(\xi, \xi)/\xi$ upon integration over x with the GPD. We took the heavy-quark mass m_Q to be $M_{\mathcal{Q}}/2$. Inspired by potential-model values [30, 31], we chose $R_{J/\psi}(0) = 1$ GeV$^{3/2}$ and $R_\Upsilon(0) = 3$ GeV$^{3/2}$. With regard to α_s, we evolve at two-loop with `CRunDec` [32] using $\alpha_s(M_Z^2) = 0.118$. The resulting LO cross section is shown in grey on Fig. 1. The μ_R uncertainty at LO is a trivial normalisation shift and is not shown. One observes a good agreement with the data within large theoretical uncertainties from μ_F which we discuss now. First it grows with energy following the singular $1/x$ behaviour of P_{gg}. Second, one observes a strange behaviour of the lower bound which exhibits a slight decrease vs. energy above 30 GeV. Such a behaviour should in fact be attributed to that of the gluon PDF at very low scale which can, for some PDFs like CTEQ6M used here, develop a local minimum for x on the order of 0.001 (see [12]). In both cases, this is not connected to the GPD itself.

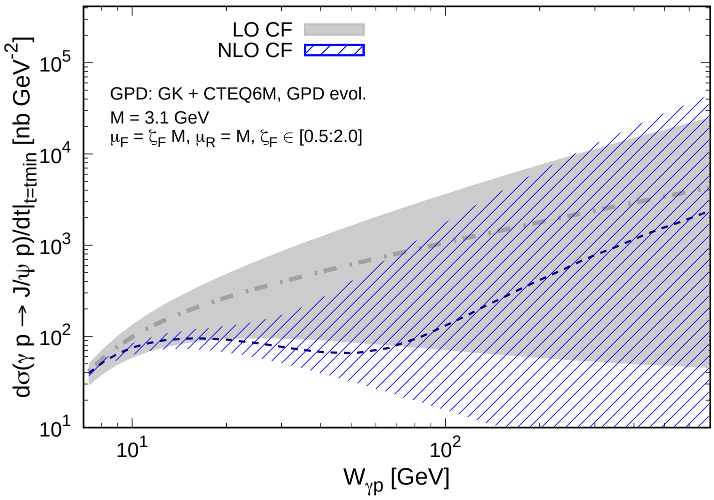


Figure 1: LO and NLO t-differential cross section for exclusive J/ψ photoproduction.

Let us now move on to the discussion of the NLO cross section at large $W_{\gamma p}$, i.e. small ξ. In this limit, the NLO coefficient functions $C_{g,q}^{(1)}(x, \xi; \mu_F, \mu_R)$ scale like

$$-\frac{i\pi c|x|}{2\xi}\frac{\alpha_s(\mu_R)}{\pi}\ln\left(\frac{m_Q^2}{\mu_F^2}\right)\left\{C_A, 2C_F\right\}, \qquad \xi < |x| < 1\,, \tag{9}$$

and we refer to this limiting scaling as $C^{(1,\text{ asy.})}_{\{g,q\}}(x,\xi;\mu_F,\mu_R)$.

Under these conditions, we can trivially deduce the effect of the x integration once one notes that the x dependence of the gluon GPD $H_g(x,\xi)$ is close to a constant for $\mu_F \sim 2-3$ GeV which is precisely the range used for theoretical computations for J/ψ. It follows that the corresponding NLO contribution to the amplitude of Eq. (1) is enhanced by $\ln(1/\xi)$ at small ξ. A similar observation can be made for the (NLO) quark contributions with the flavour-singlet GPD scaling like $1/x$ once one notes the different power in the prefactor $1/x$. Such observations were already made in Ref. [4]. This leads to the catastrophically large μ_F dependence of the NLO CF (blue) J/ψ cross section in Fig. 1 compared to the LO (grey) one at large $W_{\gamma p}$ accompanied by an unphysical oscillation due to the change of sign of the imaginary part of the amplitude. In corresponding inclusive computations [9–13], the cross section simply gets negative for increasing colliding energies and the unphysical character of the energy dependence is even more explicit.

One solution to cure this issue is simply to nullify $C^{(1,\text{ asy.})}_{q,g}$ by setting $\mu_F = m_Q$. Better results were indeed observed for low μ_F values in [4] and such a scale choice (denoted $\hat{\mu}_F$ in what follows) was advocated [8] as an approximate way to resum a series of corrections to the amplitude which are proportional to $(\alpha_s \ln(\mu_F/m_Q)\ln(1/\xi))^n$. It was then used as the unique μ_F choice in several NLO phenomenological analyses [33–35]. A similar scale-fixing prescription was advocated for P_T-integrated cross sections of *inclusive* hadro- [12] and photoproduction [13] of heavy quarkonia and it was identified in [14, 15] that such a choice however misses all corrections $\propto (\alpha_s \ln(1/\xi))^n$ without $\ln(\mu_F)$. As we discuss later, this is also the case in exclusive photoproduction.

4 Resummation via High-Energy Factorisation

To ease the discussion, let us introduce the variable $\rho = \xi/x$ and consider the limit $|\rho| \ll 1$ relevant for high energies. As discussed in [4, 36], $C_i(x,\xi;\mu_F,\mu_R)$ receives a series of corrections proportional to $(\alpha_s^n \ln^{n-1}(1/|\rho|))/|\rho|$. This series can be resummed using *High-Energy Factorisation* [37–40] which was developed for *inclusive* processes. The application of HEF to *exclusive* processes is possible as explained in [16] and the resummation of the above mentioned series corresponds to the Leading-Logarithmic Approximation (LLA) which is within theory reach.

At leading power in $|\rho| \ll 1$, the resummed coefficient function $C_i^{(\text{HEF})}$ replaces the CF coefficient C_i in Eq. (1) which factorises in the following way [7, 36]:

$$C_i^{(\text{HEF})}(\rho;\mu_F,\mu_R) = \frac{-i}{|\rho|}\int_0^\infty d\mathbf{k}_\perp^2\ \mathcal{C}_{gi}(|\rho|,\mathbf{k}_\perp^2;\mu_F,\mu_R)h(\mathbf{k}_\perp^2). \tag{10}$$

Here, $\mathcal{C}_{gi}(\rho,\mathbf{k}_\perp^2;\mu_F,\mu_R)$ is a process independent HEF resummation factor and $h(\mathbf{k}_\perp^2)$ is the process dependent HEF coefficient function at LO in α_s which is well known (see e.g. [36, 41]) and reads $h(\mathbf{k}_\perp^2) = \pi c m_Q^2/(2(m_Q^2+\mathbf{k}_\perp^2))$. We refer to [16] for derivation of the explicit results.

As we noted in [14], $\mathcal{C}_{gi}$ in the (complete) LLA includes [37–40] some μ_F-dependent terms which will not be compensated by the usual fixed-order evolution of GPDs. To cope with this limitation, we advocate a truncation of LLA down to the *doubly-logarithmic approximation (DLA)* in order to stay consistent with GPD evolution. Doing so, one only resums corrections scaling like $\alpha_s^n(\mu_R)\ln^{n-1}(1/\rho)\ln^{n-1}(\mu_F^2/\mathbf{k}_\perp^2)$ in $\mathcal{C}_{gi}$. Introducing the Mellin conjugate variable conjugate to ρ, N, one gets a compact expression for the gluon channel [39, 42]:

$$\mathcal{C}_{gg}^{(\text{DL})}(\rho,\mathbf{k}_\perp^2) = \int \frac{dN}{2\pi i}\,\rho^{-N}\frac{\gamma_N}{\mathbf{k}_\perp^2}\left(\frac{\mathbf{k}_\perp^2}{\mu_F^2}\right)^{\gamma_N}, \tag{11}$$

where $\gamma_N = \hat{\alpha}_s(\mu_R)/N$ and $\hat{\alpha}_s = \alpha_s C_A/\pi$.

Injecting Eq. (11) into Eq. (10), the resummed coefficient function in the DLA HEF in Mellin space $C_g^{(\text{DLA HEF})}(N)$ has a closed-form expression:

$$C_g^{(\text{DLA HEF})}(N) = \frac{-i\pi c}{2}\left(\frac{m_Q^2}{\mu_F^2}\right)^{\gamma_N}\frac{\pi\gamma_N}{\sin(\pi\gamma_N)}. \tag{12}$$

It can be expanded in α_s and transformed to ρ space to interpret the results. Up to NNLO, we have [16] (with an overall factor $-i\pi c/2$): $\delta(|\rho|-1) + \frac{\hat{\alpha}_s}{|\rho|}\ln\left(\frac{m_Q^2}{\mu_F^2}\right) + \frac{\hat{\alpha}_s^2}{|\rho|}\ln\frac{1}{|\rho|}\left[\frac{\pi^2}{6}+\frac{1}{2}\ln^2\left(\frac{m_Q^2}{\mu_F^2}\right)\right]$. The first two terms coincide (as they should) with the imaginary part of $C_g^{(0)}$ and $C_g^{(1,\text{ asy.})}$. The same applies for the quark contribution. Such an expansion can also give us a prediction of the NNLO term of the coefficient function in this asymptotic limit as our DLA indeed coincides with the complete LLA up to NNLO in α_s.

We guide the reader to [16] for the complete DLA expression in ρ space which can be used for phenomenological applications. Let us however note that, in the limit $|\rho| \ll 1$, one has[b] $\check{C}_g^{(\text{HEF})} \propto |\rho|^{-\hat{\alpha}_s - 1}$. This signals the presence of a *hard-Pomeron* behaviour of the coefficient function unlike the case of the scale-fixing solution. Let us add that, in a complete LLA treatment [43–45], the value of the exponent to $|\rho|$ would exhibit $4\hat{\alpha}_s\ln 2$ instead $\hat{\alpha}_s$ with DLA.

[b] The ˇ mark on C indicates the subtraction of the α_s^0 term.

5 Matching of the NLO and resummed CF results

The NLO CF of Sec. 3. and the DLA HEF of Sec. 4 are two approximations for the (all order) CF coefficient function which are most accurate in different ρ regions. We would like to use both and it is important to combine them while keeping in mind that they can contain the same contributions. To proceed, we employ a simple *subtractive-matching* prescription:

$$C_{g,q}^{(\text{match.})}(x,\xi) = C_{g,q}^{(0)}(x,\xi) + C_{g,q}^{(1)}(x,\xi) + [\check{C}_{g,q}^{(\text{HEF})}\,(\xi/|x|) - C_{g,q}^{(1,\text{ asy.})}(x,\xi)]\theta(|x|-\xi). \tag{13}$$

In what follows, we will refer to results obtained with such a matching as "NLO CF $\oplus$ DLA HEF" results. $C_{g,q}^{(\text{HEF})}$ is applicable only for $|x| > \xi$, namely the DGLAP region, and this is why we added $\theta(|x|-\xi)$ in Eq. (13). For $\rho \to 0$ or equivalently $\xi \ll |x| < 1$, we obtain, by definition, $C_{g,q}^{(1)}(x,\xi) \simeq C_{g,q}^{(1,\text{ asy.})}(x,\xi)$. For $\rho \to 1$ or equivalently to $\xi \lesssim x < 1$, $\ln^n(1/|\rho|)$ tend to 0 and $\check{C}_{g,q}^{(\text{HEF})}(\xi/|x|) \simeq C_{g,q}^{(1,\text{ asy.})}(x,\xi)$. It follows that, in both limits, the matched result corresponds to the most accurate approximation, i.e. $C_{g,q}^{(\text{HEF})}$ from DLA HEF for $\rho \to 0$ and $C_{g,q}^{(0)} + C_{g,q}^{(1)}$ from NLO CF for $\rho \to 1$. Even though the limiting cases are exactly what we want, we note that such a subtraction affects the matched contributions in the entire DGLAP region. A matching procedure using *inverse-error-weighting* (InEW) [14, 15, 46] would avoid such a shortcoming and is left for future investigation.

6 Resummed results

The scale dependence of the NLO CF $\oplus$ DLA HEF results has been thoroughly discussed in [16]. We briefly summarise here the most relevant features. With regard to the μ_F sensitivity,

- for the J/ψ case, it remains roughly constant and below a factor of 3 for the NLO CF $\oplus$ DLA HEF, unlike the NLO (blue) ones which blows up;
- for the Υ case, it only slowly increases to reach a factor of 2 at 2 TeV, unlike the NLO one which shows a prominent maximum close to a factor of 5 around 1 TeV which corresponds to the point where the imaginary part changes sign;

In both cases, NLO CF $\oplus$ DLA HEF exhibits a clear smaller sensitivity than the LO, as expected from higher-order computations. Regarding the μ_R sensitivity,

- at LO, it trivial comes from the normalisation shift from $\alpha_s(4m_Q)$ to $\alpha_s(m_Q)$;
- the NLO CF $\oplus$ DLA HEF for the J/ψ case shows a slow increase with energy which is compatible with the μ_R dependence of the hard-pomeron contribution with an asymptotic behaviour as $|\rho|^{-\hat{\alpha}_s(\mu_R)-1}$, unlike the NLO (blue) one in Fig. 1 which is completely out of control reaching a factor of 100;
- for the Υ case, it is at most 20% in the energy region that we consider, unlike the NLO one which reaches a factor of 20. The hard-pomeron behaviour seems to set in even above the energies at a possible FCC-eh.

In short, we can claim that the perturbative instabilities of the NLO results are cured by the HEF resummation. The only caveat is that the μ_R dependence at very high energy will only be mitigated with a complete NLLA computation. It will also be instructive to investigate a different matching which may be less sensitive to the completely anomalous behaviour of the NLO contributions at small ρ.

Now that we have summarised how the HEF resummation improves the stability of the results, we can proceed with comparisons with experimental data. We guide the reader to [16] regarding the (small) corrections we brought in to make comparisons at $t = t_{\min}$. To assess the theoretical uncertainties, lacking a better procedure, we consider a 9-point scale variation[c] about $M_{\mathcal{Q}}$. Although it is standard, we recognise that it remains arbitrary and is not meant to be considered as a true 1σ uncertainty.

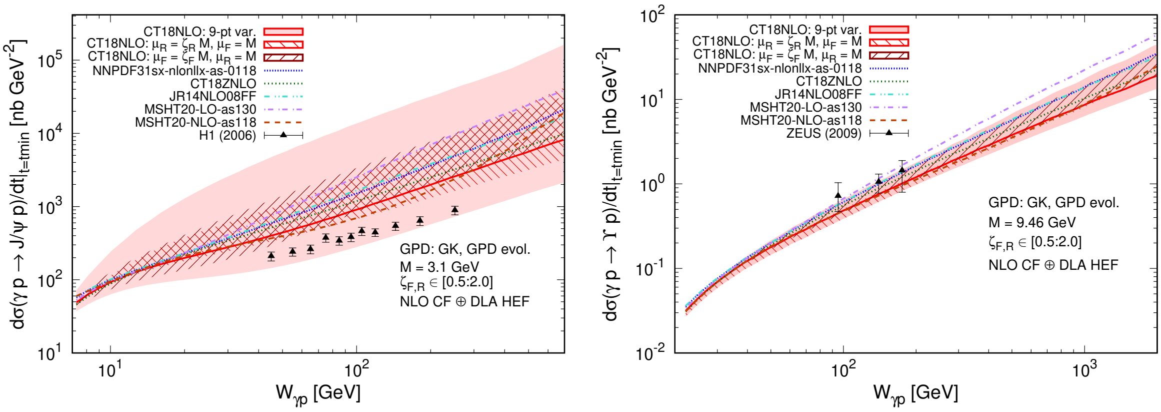
Regarding the input PDF choices, we have considered the central eigenvector of several PDF sets: CT18NLO [25], CT18ZNLO [25], an alternate CT18 fit with a different gluon PDF, the LO set MSHT20-LO-as130 [47], another conventional NLO set MSHT20-NLO-as118 [47], and two sets known to be steeper at low x and low scales, JR14NLO08FF [48] and NNPDF31sx-nlonllx-as-0118 [49]. Using all the eigenvectors of a set would have been significantly more resource-consuming in the GPD grid generation and has been left for a future work.

As shown in Fig. 2a, our J/ψ results feature large uncertainties but are in agreement with the experimental data which tend to lie on the lower side. The cross-section predictions obtained using the central values of different PDF sets differ by a factor of $\mathcal{O}(2-3)$, similar to the size of the μ_F and μ_R uncertainties taken separately. For Υ, our band shown in Fig. 2b also agrees with the experimental data. We note that in this case, both types of uncertainties are of similar size hinting at a clear possibility to better constrain PDFs using Υ high-energy data.

7 Conclusions

We have reviewed the main results of our recent study [16] of exclusive J/ψ and Υ production in CF at NLO supplemented by the resummation of a series of higher-order corrections in CF which are enhanced by high-energy

[c] μ_R and μ_F are independently set to $\mu_F = \mu_0 \times \zeta_F$ and $\mu_R = \mu_0 \times \zeta_R$, for $\zeta_{F,R} \in \{1/2, 1, 2\}$ and $\mu_0 = M_{\mathcal{Q}}$. The filled band is the envelope of the corresponding cross sections. As a complement, we show the envelope corresponding to the sole variation of μ_F and μ_R with hatched bands.



(a) (b)

Figure 2: (a) NLO CF $\oplus$ DLA HEF t-differential cross section (see text for details) compared to data extrapolated down to $t = t_{\text{min}}$ from the H1 collaboration [50]. (b) Same as (a) but for Υ production and data from the ZEUS collaboration [51].

logarithms. We have found that such a resummation, performed via HEF, is crucial to resolve the perturbative instability of the NLO CF computation at $W_{\gamma p} \gg M_{\mathcal{Q}}$. This approach offers an accurate description of this observable, over a wider energy range than other approaches such as e.g. CGC [52], small-x [53] or k_T-factorisation [41] approaches, which are applicable only at high energies. We believe our work is an important addition to other recent related studies [54–57].

Our results for the J/ψ case are characterised by large scale uncertainties, which is regrettably so since it features the largest experimental dataset. The uncertainties are significantly smaller for the Υ case and we are clearly eager for more data, in particular from LHC pPb UPCs [58], especially if they can be measured differentially in t. In addition, the scale uncertainty will be reduced once we develop this computation beyond the DLA (see e.g. [59]).

Let us add that our study confirmed that exclusive J/ψ and Υ photoproduction is nearly entirely gluon-induced. Quark-induced contributions, which only arise at NLO from radiative corrections, are at most at the level of 20% and essentially via interference with gluon-induced contributions for $W_{\gamma p} > 100$ GeV. It is also clear that the theoretical features, like the scale and flavour dependencies, encoded in NLO CF and resummation results cannot be reproduced if one simply approximates GPDs to a square of PDFs. In general, more work is needed to extract reliable constraints on PDFs from a global GPD fit with PDFs as forward constraints. We are currently working on including relativistic-v^2 corrections in a similar HEF computation matched to CF. Preliminary results indicate a sizable negative correction to the amplitude which would result in a smaller cross section, in better agreement with data.

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