Rapidity gap survival in central exclusive diffraction: Dynamical mechanisms and uncertainties[*]

Mark Strikman[a], Christian Weiss[b]
[a] Department of Physics, Pennsylvania State University, University Park, PA 16802, USA
[b] Theory Center, Jefferson Lab, Newport News, VA 23606, USA

Abstract

We summarize our understanding of the dynamical mechanisms governing rapidity gap survival in central exclusive diffraction, $pp \rightarrow p + H + p$ (H = high–mass system), and discuss the uncertainties in present estimates of the survival probability. The main suppression of diffractive scattering is due to inelastic soft spectator interactions at small pp impact parameters and can be described in a mean–field approximation (independent hard and soft interactions). Moderate extra suppression results from fluctuations of the partonic configurations of the colliding protons. At LHC energies absorptive interactions of hard spectator partons associated with the $gg \rightarrow H$ process reach the black–disk regime and cause substantial additional suppression, pushing the survival probability below 0.01.

1   Strong interaction dynamics in rapidity gap survival

Calculation of the cross section of central exclusive diffraction, $pp \rightarrow p + H + p$ (H = dijet, heavy quarkonium, Higgs boson, *etc.*) presents a major challenge for strong interaction physics. It involves treating the hard dynamics in the elementary $gg \rightarrow H$ subprocess, and calculating the probability that no other interactions leading to hadron production occur during the pp collision. The latter determines the suppression of diffractive relative to non-diffractive events with the same hard process, referred to as the rapidity gap survival (RGS) probability. In this article we summarize our understanding of the dynamical mechanisms determining the RGS probability, their phenomenological description, and the uncertainties in present numerical predictions.

RGS in central exclusive diffraction has extensively been discussed in an approach where soft interactions are modeled by eikonalized pomeron exchange; see Ref. [1] for a summary. More recently a partonic description was proposed, which allows for a model–independent formulation of the interplay of hard and soft interactions and reveals the essential role of the "transverse geometry" of the pp collision [2]. In the mean–field approximation, where hard and soft interactions are considered as independent aside from their common dependence on the impact parameter, we derived a simple "factorized" expression for the RGS probability, using closure of the partonic states to take into account inelastic diffractive intermediate states. The resulting RGS probability is smaller than in the models of Refs. [1, 3] without inelastic diffraction,

but comparable to the some of the versions of those models with multichannel diffraction. Our partonic description also permits us to go beyond the mean–field approximation and incorporate various types of correlations between the hard scattering process and spectator interactions. Here we discuss two such effects: (a) quantum fluctuations of the partonic configurations of the colliding protons, which somewhat reduce the survival probabilities at RHIC and Tevatron energies; (b) absorptive interactions of high-virtuality spectator partons ($k^2 \sim$ few GeV^2) associated with the hard scattering process, related to the onset of the black–disk regime (BDR) in hard interactions at LHC energies; this new effect substantially reduces the RGS probability compared to previously published estimates.

2 Soft spectator interactions in the mean–field approximation

A simple picture of RGS is obtained in the impact parameter representation. On one hand, to produce the heavy system H two hard gluons from each of the two protons need to collide in the same space–time point (actually, an area of transverse size $\sim 1/\langle k_T^2 \rangle$ in the hard process); because such gluons are concentrated around the transverse centers of the protons this is most likely when the protons collide at a small impact parameters, $b \lesssim 1\,\mathrm{fm}$. On the other hand, soft inelastic spectator interactions are strongest at small b and would favor collisions at $b \gg 1\,\mathrm{fm}$ for diffractive scattering. These different preferences limit diffraction to an intermediate range of impact parameters and ensure that its cross section is substantially suppressed compared to non–diffractive scattering. More precisely, the RGS probability is given by [2]

$$S^2 \;\; = \;\; \int d^2b \; P_{\mathrm{hard}}(b) \; |1 - \Gamma(b)|^2, \qquad\qquad b \equiv |\boldsymbol{b}|. \qquad (1)$$

Here $P_{\mathrm{hard}}(b)$ is the probability for two gluons to collide at the same transverse point as a function of the pp impact parameter, given by the convolution of the transverse spatial distributions of the gluons in the colliding protons, normalized to $\int d^2b \; P_{\mathrm{hard}}(b) = 1$ (see Fig. 1a). The factor $|1 - \Gamma(b)|^2$ is the probability for the two protons not to interact inelastically in a collision at the given impact parameter, calculable in terms of the profile function of the pp elastic amplitude, $\Gamma(b)$. Figure 1b shows the b–dependence of the two factors as well as their product, illustrating the interplay described above. While we have motivated Eq. (1) by probabilistic arguments, it actually can be derived (as well as the expression for the differential cross section) in the partonic description of Ref. [2] within the mean–field approximation, where one assumes no correlation between the presence of the gluons involved in the hard interaction (with the particular x) and the strength of the soft spectator interactions. In this approximation one can use closure to sum over the different diffractive intermediate states, and thus effectively include the contribution of inelastic diffraction.[1] The numerical values of the RGS probability obtained from Eq. (1) are of the order $S^2 \sim 0.03$ for $M_H = 100\,\mathrm{GeV}$ and $\sqrt{s} = 14\,\mathrm{TeV}$; see Ref. [2] for details.

It is worthwhile to discuss the uncertainty in the numerical predictions for S^2 in the mean–field approximation, Eq. (1), resulting from our imperfect knowledge of the functions in the

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[1]In principle there is also a contribution from excitation of a diffractive state by soft spectator interactions and subsequent transition back to the proton via the nondiagonal gluon GPD; however, it is strongly suppressed because the typical excitation masses in hard and soft diffraction are very different in the kinematics of Higgs production at the LHC ($10^{-8} \leq x_{I\!\!P} \leq 0.1$ for generic pp diffraction and $10^{-2} \leq x_{I\!\!P} \leq 0.1$ for the GPD); see Section IV C of Ref. [2].
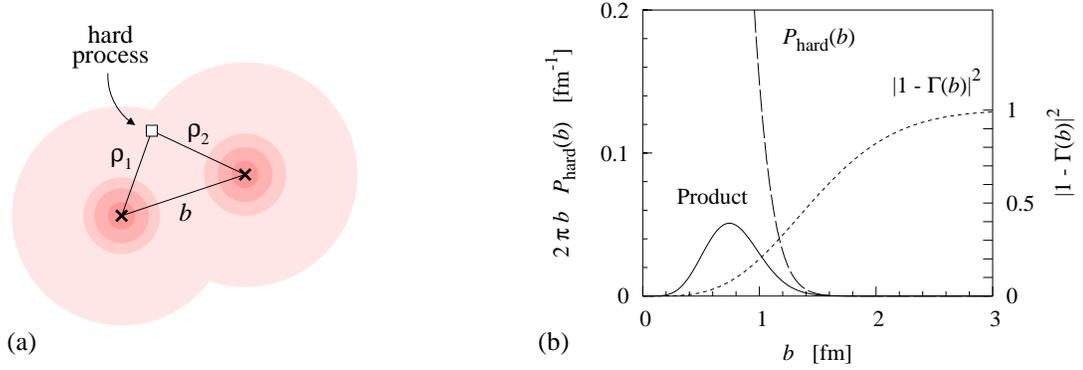
Fig. 1: (a) Transverse geometry of hard diffractive pp scattering. (b) RGS probability in the impact parameter representation *cf.* Eq. (1), for $\sqrt{s} = 14\,\text{TeV}$, $M_H \sim 100\,\text{GeV}$ [2]. Dashed line: Probability for hard scattering process $P_{\text{hard}}(b)$ (left vertical axis). Dotted line: Probability for no inelastic interactions between the protons, $|1 - \Gamma(b)|^2$ (right vertical axis). Solid line: Product $P_{\text{hard}}(b)|1 - \Gamma(b)|^2$ (left vertical axis). The RGS probability Eq. (1) is given by the area under this curve.
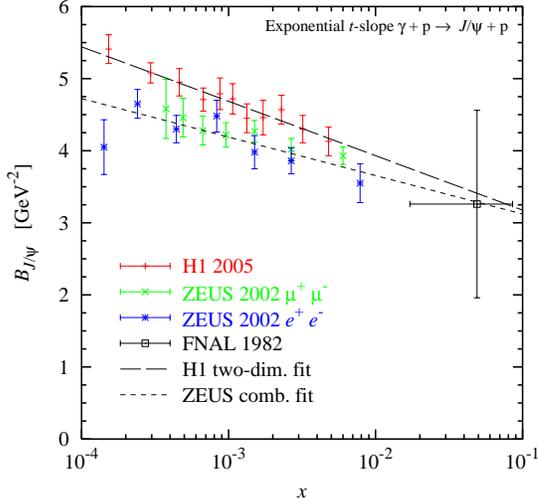

Fig. 2: The exponential t–slope, $B_{J/\psi}$, of exclusive J/ψ photoproduction, extracted from fits to the FNAL E401/E458 [4], HERA H1 [5], and ZEUS [6] data. The long–dashed and short–dashed lines represent fits to the x–dependence of the H1 and ZEUS t–slopes [5, 6], *cf.* Eq. (2). The t–slope of the gluon GPD, B_g, is obtained from $B_{J/\psi}$ after applying a small correction for the finite size of the J/ψ [7].

integrand. We first consider the transverse spatial distribution of gluons entering in $P_{\text{hard}}(b)$. The latter is obtained as the Fourier transform of the t–dependence (more precisely, transverse momentum dependence) of the gluon generalized parton distribution (GPD) measured in hard exclusive vector meson production. Extensive studies at HERA have shown that exclusive J/ψ photoproduction, $\gamma p \to J/\psi + p$, provides an effective means for probing the t–dependence of the gluon GPD at small and intermediate x (a small correction for the finite transverse size of the J/ψ is applied) [7]. Figure 2 summarizes the results for the exponential t–slope of this process, $B_{J/\psi}$, from HERA H1 [5] and ZEUS [6] and the FNAL E401/E458 experiment [4], as well as fits to the x–dependence of the H1 and ZEUS results of the form (here $x = M_{c\bar{c}}^2/W^2$)

$$B_{J/\psi}(x) \;=\; B_{J/\psi}(x_0) \;+\; 2\alpha'_{J/\psi}\,\ln(x_0/x). \tag{2}$$

There is a systematic difference between the H1 and ZEUS results due to different analysis methods [5, 6]; however, the fits to both sets agree well with the FNAL point when extrapolated to larger x. In diffractive production of a system with $M_H = 100\,\text{GeV}$ at $\sqrt{s} = 14\,\text{TeV}$ at zero rapidity the gluons coupling to the heavy system H have momentum fractions $x_{1,2} = M_H/\sqrt{s} = 0.007$. Assuming exponential t–dependence of the gluon GPD, we can estimate the uncertainty in the transverse spatial distribution of gluons at such x by evaluating the fits to the HERA data within the error bands quoted for $B_{J/\psi}(x_0)$ and $\alpha'_{J/\psi}$ [5, 6]. We find a 15-20% uncertainty of $B_{J/\psi}$ at $x = 0.007$ in this way, translating into a 20–30% uncertainty in the mean–field RGS probability, Eq. (1). We note that there is at least a comparable uncertainty in S^2 from the uncertainty of the shape of the t–dependence; this is seen from Fig. 10 of Ref. [2], where the exponential is compared with a theoretically motivated dipole form which also describes the FNAL data. Altogether, we estimate that our imperfect knowledge of the spatial distribution of gluons results in an uncertainty of the mean–field result for S^2 by a factor ~ 2. Dedicated analysis of the remaining HERA exclusive data, and particularly precision measurements with a future electron–ion collider (EIC), could substantially improve our knowledge of the transverse spatial distribution of gluons.

We now turn to the uncertainty in S^2 arising from the pp elastic amplitude, $\Gamma(b)$. Most phenomenological analyses of pp elastic and total cross section data find that for TeV energies $|1 - \Gamma(b)| \leq 0.05$ at $b = 0$, corresponding to near–unit probability of inelastic interactions at small impact parameters (BDR). This is supported by theoretical studies in the QCD dipole model, which show that the large–x partons with virtualities of up to several GeV^2 experience "black" interactions with the small–x gluon field in the other proton when passing through the other proton at transverse distances $\rho \leq 0.5\text{fm}$, and receive transverse momenta $k_T \geq 1\,\text{GeV}$ (see Ref. [7] for a summary). At pp impact parameter $b = 0$ the chance that none of the leading partons in the protons receive such a kick is extremely small, implying that $|1 - \Gamma(b)| \sim 0$ [8]. For the RGS probability in the mean–field approximation, Eq. (1), the fact that $|1 - \Gamma(b)|^2$ is small at $b = 0$ is essential, as this eliminates the contribution from small b in the integral (see Fig. 1b) and stabilizes the numerical predictions. However, present theoretical arguments and data analysis cannot exclude a small non-zero value of $|1 - \Gamma(b)|$ at $b = 0$; a recent analysis finds $|1 - \Gamma(b)| \sim 0.1$ [9]. To investigate the potential implications for the RGS probability, we evaluate Eq. (1) with the Gaussian parametrization of $\Gamma(b)$ of Ref. [2], Eq. (12), but with $\Gamma(b = 0) = 1 - \epsilon$. We find that a value of $\epsilon = 0.1$, corresponding to $|1 - \Gamma(b)|^2 = 0.01$, increases

the mean–field result for S^2 by a factor ~ 1.8, indicating significant uncertainty of the mean–field result. However, as explained in Sec. 4 below, hard spectator interactions associated with the $gg \to H$ process lead to an additional suppression of diffraction at small b (not contained in the soft RGS probability), which mitigates the impact of this uncertainty on the overall diffractive cross section.

3 Fluctuations of parton densities and soft–interaction strength

Corrections to the mean–field picture of RGS arise from fluctuations of the interacting configurations in the colliding protons. This concept is known well in soft diffraction, where fluctuations of the strength of interaction between the colliding hadrons give rise to inelastic diffraction. In hard diffraction, one expects that also the gluon density fluctuates; *e.g.* because the color fields are screened in configurations of small size [10]. In fact, the variance of the gluon density fluctuations can be directly related to the ratio of inelastic and elastic diffraction in processes such as $\gamma_L^* + p \to$ "vector meson" $+ X$,

$$\omega_g \; \equiv \; \frac{\langle G^2 \rangle - \langle G \rangle^2}{\langle G \rangle^2} \;\; = \;\; \frac{d\sigma_{\mathrm{inel}}}{dt} \bigg/ \frac{d\sigma_{\mathrm{el}}}{dt} \bigg|_{t=0}. \tag{3}$$

The HERA data are consistent with the dynamical model estimate of $\omega_g \sim 0.15 - 0.2$ for $Q^2 = 3\,\mathrm{GeV}^2$ and $x \sim 10^{-4} - 10^{-3}$ [10]; unfortunately, the limited Q^2 range and the lack of dedicated studies do not allow for a more precise extraction of this fundamental quantity.

In central exclusive diffraction, correlated fluctuations of the soft–interaction strength and the gluon density lower the RGS probability, because small-size configurations which experience little absorption have a lower gluon density. This effect can be modeled by a generalization of the mean–field expression (1), in which both the gluon GPDs in P_{hard} and the profile function fluctuate as a function of an external parameter controlling the overall size of the configurations [10]. Numerical studies find a reduction of the RGS probability by a factor $\sim 0.82\,(0.74)$ for a system with mass $M_H = 100\,\mathrm{GeV}$ produced at zero rapidity at $\sqrt{s} = 2\,(14)\,\mathrm{TeV}$. The dynamical model used in this estimate does not include fluctuations of the gluon density at larger $x(\sim 0.05 - 0.1)$, which could increase the suppression.

We emphasize again that inelastic diffraction *per se* is included in the partonic approach of Ref. [2] through the closure of partonic states. The effect discussed in this section is specifically related to correlations between the fluctuations of the parton densities and the soft–interaction strength; in the limit of zero correlations (independent fluctuations) we recover the mean–field result described above [10].

4 Black–disk regime in hard spectator interactions

Substantial changes in the mechanism of diffractive scattering are brought about by the onset of the BDR in hard interactions at LHC energies, where even highly virtual partons ($k^2 \sim$ few GeV^2) with $x \gtrsim 10^{-2}$ experience "black" interactions with the small–x gluons in the other proton. This new effect modifies the amplitude of central exclusive diffraction in several ways: (a) absorption of the "parent" partons of the gluons attached to the high–mass system; (b) absorption of the hard gluons attached to the high–mass system; (c) absorption due to local interactions
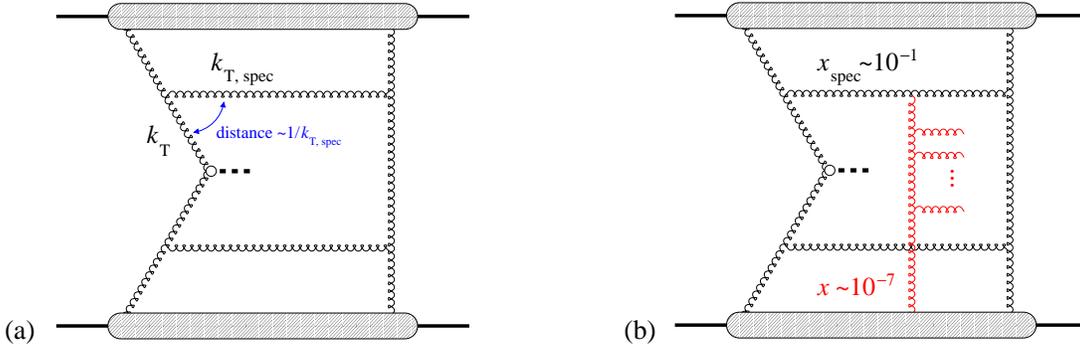
Fig. 3: (a) QCD evolution–induced correlation between hard partons. The transverse distance between the active parton and the spectator is $\sim 1/k_{T,\,\mathrm{spec}}$. (b) Absorptive interaction of the hard spectator with small–x gluons in the other proton.

within the partonic ladder. Such absorptive hard interactions cause additional suppression of diffractive scattering, not included in the traditional soft–interaction RGS probability [2]. Because of the generic nature of "black" interactions, we can estimate this effect by a certain modification of the mean–field picture in the impact parameter representation. Here we focus on mechanism (a) and show that it causes substantial suppression; the other mechanisms may result in further suppression.

According to Ref. [11] (and references therein) the dominant contribution to the hard amplitude of Higgs production at the LHC ($M_H = 100\,\mathrm{GeV}$, $x_{1,2} \sim 10^{-2}$) originates from gluons with transverse momenta of the order $k_T \sim 2\,\mathrm{GeV}$. Such gluons are typically generated by DGLAP evolution starting from the initial scale, Q_0^2, in which spectator partons, mostly gluons, are emitted (see Fig.3a). In the leading–log approximation $Q_0 \ll k_{T,\,\mathrm{spec}} \ll k_T$, and thus the transverse distance between the active and spectator parton is $\sim 1/k_{T,\,\mathrm{spec}} \ll R_{\mathrm{proton}}$, amounting to short–range correlations between partons. If the interactions of the spectator parton with the small–x gluons in the other proton become significant (see Fig.3b), the basic assumption of the mean–field approximation — that the spectator interactions are independent of the hard process — is violated, and the interactions of that parton need to be treated separately. Indeed, studies within the QCD dipole model show that at the LHC energy spectator gluons with $k_{T,\,\mathrm{spec}} \sim 1\,\mathrm{GeV}$ and $x_{\mathrm{spec}} \sim 10^{-1}$ "see" gluons with momentum fractions $x \sim 10^{-7}$ in the other proton, and are absorbed with near–unit probability if their impact parameters with the other proton are less than $\sim 1\,\mathrm{fm}$ [2].[2] For pp impact parameters $b < 1\,\mathrm{fm}$ about 90% of the strength in $P_{\mathrm{hard}}(b)$ comes from parton–proton impact parameters $\rho_{1,2} < 1\,\mathrm{fm}$ (*cf.* Fig. 1a), so that this effect practically eliminates diffraction at $b < 1\,\mathrm{fm}$. Since $b < 1\,\mathrm{fm}$ accounts for 2/3 of the cross section (see Fig. 1b), and the remaining contributions at $b > 1\,\mathrm{fm}$ are also reduced by absorption, we estimate that absorptive interactions of hard spectators in the BDR reduce the RGS probability at LHC to about 20% of its mean–field value. Much less suppression is expected

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[2]The cross section of "gluonic" (88) dipoles is larger than that of the quark–antiquark ($\bar{3}3$) dipoles in $\gamma^* p$ scattering [12] by a factor $9/4$. A summary plot of the profile function for gluon–proton scattering is given in Fig. 13 of Ref. [7] (right y–axis). Note that $\Gamma^{\mathrm{gluon-proton}} = 0.5$ already corresponds to a significant absorption probability of $1 - |1 - \Gamma^{\mathrm{gluon-proton}}|^2 = 0.75$.

at the Tevatron energy, where hard spectator interactions only marginally reach the BDR.

In the above argument one must also allow for the possibility of trajectories with no gluon emission, which correspond to the Sudakov form factor–suppressed $\delta(1 - x)$–term in the evolution kernel. While such trajectories are not affected by absorption, their contributions are small both because of the Sudakov suppression, and because they effectively probe the gluon density at a low scale, $Q_0^2 \sim 1\,\mathrm{GeV}^2$, where evolution–induced correlations between partons can be neglected. We estimate that the contribution of such trajectories to the cross section is suppressed compared to those with emissions by a factor $R = \left[S_G^2\, G(x, Q^2)/G(x, Q_0^2) \right]^2 \sim 1/10$, where $S_G^2 = \exp[-(3\alpha_s/\pi)\ln^2(Q^2/Q_0^2)]$ is the square of the Sudakov form factor, and $Q^2 \sim 4\,\mathrm{GeV}^2$. Their net contribution is thus comparable to that of the trajectories with emissions, because the latter are strongly suppressed by the absorption effect described above. Combining the two, we obtain an overall suppression by a factor of the order ~ 0.3. More accurate estimates would need to take into account fluctuations in the number of emissions; in particular, trajectories on which only one of the partons did not emit gluons are suppressed only by $\sqrt{R}$ and may make significant contributions.

The absorptive hard spectator interactions described here "push" diffractive pp scattering to even larger impact parameters than would be allowed by the soft spectator interactions included in the mean–field RGS probability, Eq. (1) (except for the Sudakov–suppressed contribution). One interesting consequence of this is that it makes the uncertainty in the mean–field prediction arising from $\Gamma(0) \neq 1$ (see Sec. 2) largely irrelevant, as the region of small impact parameters is now practically eliminated by the hard spectator interactions. Another consequence is that the final–state proton transverse momentum distribution is shifted to to smaller values; this could in principle be observed in p_T–dependent measurements of diffraction. We note that the estimates of hard spectator interactions reported here are based on the assumption that DGLAP evolution reasonably well describes the gluon density down to $x \sim 10^{-6}$; the details (but not the basic picture) may change if small–x resummation corrections were to significantly modify the gluon density at such values of x (see Ref. [13] and references therein).

5   Summary

The approach to the BDR in the interaction of hard spectator partons, caused by the increase of the gluon density at small x, has profound implications for central exclusive diffraction at LHC: *No saturation without disintegration!* The RGS probability is likely to be much smaller (by a factor of $\sim 1/3$ or less) than predicted by the mean–field approximation or corresponding models which neglect correlations of partons in the transverse plane. Diffractive scattering is relegated either to very large impact parameters ($b > 1\,\mathrm{fm}$) or to Sudakov–suppressed trajectories without gluon radiation. We estimate that the overall RGS probability at LHC is $S^2 < 0.01$. Extrapolation of the Tevatron results may be misleading because interactions of hard spectators are generally far from "black" at that energy. The new effects described here call for detailed MC–based studies of possible histories of the hard scattering process and their associated spectator interactions.

References

[1] Khoze, V. A. and Martin, A. D. and Ryskin, M. G., Eur. Phys. J. **C55**, 363 (2008).

[2] Frankfurt, L. and Hyde, C. E. and Strikman, M. and Weiss, C., Phys. Rev. **D75**, 054009 (2007).

[3] Gotsman, E. and Levin, E. and Maor, U. and Miller, J. S. (2008).

[4] Binkley, M. and others, Phys. Rev. Lett. **48**, 73 (1982).

[5] Aktas, A. and others, Eur. Phys. J. **C46**, 585 (2006).

[6] Chekanov, S. and others, Nucl. Phys. **B695**, 3 (2004).

[7] Frankfurt, L. and Strikman, M. and Weiss, C., Ann. Rev. Nucl. Part. Sci. **55**, 403 (2005).

[8] Frankfurt, L. and Strikman, M. and Zhalov, M., Phys. Lett. **B616**, 59 (2005).

[9] Luna, E. G. S. and Khoze, V. A. and Martin, A. D. and Ryskin, M. G. (2008).

[10] Frankfurt, L. and Strikman, M. and Treleani, D. and Weiss, C., Phys. Rev. Lett. **101**, 202003 (2008).

[11] Khoze, Valery A. and Martin, Alan D. and Ryskin, M. G., Eur. Phys. J. **C14**, 525 (2000).

[12] Rogers, T. and Guzey, V. and Strikman, M. and Zu, X., Phys. Rev. **D69**, 074011 (2004).

[13] Ciafaloni, M. and Colferai, D. and Salam, G. P. and Stasto, A. M., JHEP **08**, 046 (2007).