Electroweak boson production in double parton scattering

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We study the $W^+W^-$ and $ZZ$ electroweak boson production in double parton scattering using QCD evolution equations for double parton distributions. In particular, we analyse the role of splitting terms in the evolution equations. We show that they give dominant contribution to the cross sections due to high mass of electroweak bosons and the lack of hadron-like form factor which suppresses the contribution discussed so far in the literature.

Keywords: quantum chromodynamics, double parton scattering, double parton distributions evolution equations, electroweak bosons

I. INTRODUCTION

The double parton scattering (DPS) in high-energy hadron scattering is a process in which two hard interactions with large scales (much bigger that nucleon mass) take place in one scattering event. Such a process is usually interpreted in QCD as scattering of two pairs of partons (quarks or gluons) from incoming hadrons. The DPS is the simplest multiparton process with hard scales which allows to gain information on parton correlations inside hadrons. Thus, it has been studied for many years from both theoretical and phenomenological sides, as well as for better description of multiparton interactions needed for modeling the underlying event, see Refs. [22, 28]. It is, therefore, very important to use a rigorous approach to the DPS which is based on QCD.

The inclusive DPS cross section in the collinear approximation takes the form [12, 15]:

$$\sigma_{AB} = N \sum_{f_1,f_2,f'_1,f'_2} \int dx_1 dx_2 dz_1 dz_2 \frac{d^2q}{(2\pi)^2} \times D_{f_1 f_2}(x_1, x_2, Q_1, Q_2, q) \hat{\sigma}^A_{f_1 f'_1}(x_1, z_1, Q_1) \hat{\sigma}^B_{f_2 f'_2}(x_2, z_2, Q_2) D_{f'_1 f'_2}(z_1, z_2, Q_1, Q_2, -q)$$ (1)

where $A$ and $B$ denote the two hard processes, $N$ is a symmetry factor equal to 1 for $A = B$ and 2 otherwise and $D_{f_1 f_2}(x_1, x_2, Q_1, Q_2, q)$ are collinear double parton distribution functions (DPDFs) in a hadron. They depend on parton flavors $f_1, f_2$, parton momentum fractions $x_1, x_2$, two hard scales $Q_1, Q_2$ and transverse relative momentum $q$. The latter momentum is related to the momentum structure of four parton fields in the definition of the DPDFs, see Fig. 1 and [15] for more details. The longitudinal momentum fractions obey the condition

$$0 < x_1 + x_2 \leq 1,$$ (2)

which says that the sum of parton longitudinal momenta cannot exceed the total nucleon momentum. This is the basic parton correlation which has to be taken into account. For more advanced aspects of parton correlations see Refs. [15, 16]. One remark is in order at this point. We start from the cross section formula in which the DPDFs depends on exchange momentum $q$ rather than on the Fourier conjugate variable $b$, being interpreted as transverse distance between two partons taking part in hard scattering. The latter possibility as a starting point creates more problems than answers [10, 15], despite its apparent attractiveness for phenomenological modeling of the DPDFs dependence on this variable.

The DPDFs obey QCD evolution equations known at present in the leading logarithmic approximation [1, 2, 5, 6, 45]. These are the Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) type evolution equations with additional non-homogeneous terms which describe splitting of a single parton into two partons. The role of these terms for the DPS
cross section predictions is the main subject of this paper. We will focus on electroweak boson production which is one of the cleanest process for such an analysis since the electroweak bosons are color singlets, thus the collinear factorization formula is not endangered by final state soft gluon interactions which might break the factorization. We show that the commonly used way to estimate the DPS cross sections no longer works in this case because of the contribution from the splitting terms in the evolution equations.

The paper is organized as follows. In Sec. I we describe the collinear evolution equations in the leading logarithmic approximation. In Sec. III we derive the general solution to these equation in the Mellin moment space while in Sec. IV we present the model of Ref. [12] for the relative momentum dependence of the DPDFs. In Sec. V we apply this model to the electroweak boson production at the LHC, paying particular attention to the role of the splitting term contribution to the cross sections. The summary are presented in the last section.

II. EVOLUTION EQUATIONS FOR DPDFS

We start by presenting evolution equations for double parton distributions with equal hard scales: \( Q_1 = Q_2 \equiv Q \). They are known in the leading logarithmic approximation (LLA) in which large powers of \( (\alpha_s \ln(Q^2/\Lambda^2))^n \) are resummed to all orders in \( n \). They were derived in [1, 2, 5, 6] for the relative momentum \( q = 0 \). However, in [10, 15] it is argued that the evolution equations in the LLA can be written for arbitrary \( q \) in the following form

\[
\partial_t D_{f_1 f_2}(x_1, x_2, t, q) = \sum_{f'} \int_{x_1}^{1-x_2} \frac{du}{u} P_{f_1 f'} \left( \frac{x_1}{u} \right) D_{f' f_2}(u, x_2, t, q) + \sum_{f'} \int_{x_2}^{1-x_1} \frac{du}{u} P_{f_2 f'} \left( \frac{x_2}{u} \right) D_{f_1 f'}(x_1, u, t, q) + \frac{1}{x_1 + x_2} \sum_{f'} P_{f' \rightarrow f_1 f_2} \left( \frac{x_1}{x_1 + x_2} \right) D_{f'}(x_1 + x_2, t)
\]

(3)

where we absorbed the leading order strong coupling constant into the definition of the evolution parameter

\[
t = t(Q) \equiv \int_{Q_0^2}^{Q^2} \frac{\alpha_s(\mu^2) d\mu^2}{2\pi \mu^2} = \frac{1}{b} \ln \frac{\ln(Q^2/\Lambda_{QCD}^2)}{\ln(Q_0^2/\Lambda_{QCD}^2)}
\]

(4)

and \( \Lambda_{QCD} \) is the dimensional parameter of QCD, \( b = 6/(33 - 2n_f) \), \( n_f \) is the number of active quark flavors. We also introduced shorthand notation for the double parton distributions with equal hard scales

\[
D_{f_1 f_2}(x_1, x_2, t, q) \equiv D_{f_1 f_2}(x_1, x_2, Q, Q, q).
\]

(5)

Notice that \( t = 0 \) corresponds to an initial scale \( Q_0 \) at which the parton distributions need to be specified. Strictly speaking, Eqs. [3] were derived in [10, 15] for the quark-antiquark double parton distributions, \( D_{\bar{q}q} \). Nevertheless, we postulate that Eqs. [3] are also valid for all combinations of partons (including gluon). The first discussion of the next-to-leading corrections to these equations can be found in [3, 4].

The integral kernels in Eqs. [3] are the leading order Altarelli-Parisi splitting functions, with virtual corrections included, which describe the splitting of one of the two partons, while the remaining parton stays intact. This gives the upper integration limits resulting from condition [2].

FIG. 1: Parton momenta in the definition of the double parton distributions.
The third term on the r.h.s describes the real splitting of parton $f'$ into two given partons $f_1$ and $f_2$. It is why the single parton distribution, $D_f(x_1 + x_2)$, appears there. The functions $P_{f' f_1 f_2}$ are directly related to the real emission Altarelli-Parisi splitting functions in the LLA, $P_{f' f}^{(0)}$. In particular, we have

$$P_{q ightarrow qg}(z) = P_{qq}^{(0)}(z), \quad P_{q ightarrow gg}(z) = P_{qq}^{(0)}(z), \quad P_{g ightarrow qg}(z) = P_{gq}^{(0)}(z), \quad P_{g ightarrow gg}(z) = P_{gg}^{(0)}(z).$$

(6)

The splitting term provides additional dependence on $t$ in Eqs. (3), imposed by the Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) evolution equations for the SPDFs

$$\partial_t D_f(x, t) = \sum_{f'} \int_0^1 \frac{du}{u} P_{f' f} \left( \frac{x}{u} \right) D_{f'}(u, t).$$

(7)

Its significance for DPS cross sections is the main subject of our paper.

III. SOLUTION TO THE EVOLUTION EQUATIONS

The evolution equations (3) can be written in a simpler form using the double Mellin moments with respect to the momentum fractions, $x_1$ and $x_2$, which take into account condition (2)

$$\tilde{D}_{f_1 f_2}(n_1, n_2, t, q) = \int_0^1 dx_1 \int_0^1 dx_2 x_1^{n_1} x_2^{n_2} \Theta(1 - x_1 - x_2) D_{f_1 f_2}(x_1, x_2, t, q)$$

(8)

where $n_1, n_2$ are complex numbers and $\theta$ is the Heaviside step function. Introducing the matrix notation with respect to parton flavours (including gluon), $\hat{D} = (\hat{D}_{f_1 f_2})$ and $\hat{P} = (P_{f_1 f_2})$, we obtain from Eqs. (3) for the Mellin moments

$$\partial_t \hat{D}(n_1, n_2, t, q) = \gamma(n_1) \hat{D}(n_1, n_2, t, q) + \hat{D}(n_1, n_2, t, q) \gamma^T(n_2) + \hat{\gamma}(n_1, n_2) \hat{D}(n_1 + n_2, t)$$

(9)

where the known matrices of anomalous dimensions read

$$\gamma(n) = \int_0^1 dx \; x^n P(x), \quad \hat{\gamma}(n_1, n_2) = \int_0^1 dx \; x^{n_1}(1 - x)^{n_2} P(x).$$

(10)

Notice that in the non-homogeneous part of Eqs. (9) $\hat{D} = (\hat{D}_f)$ is vector of the Mellin moments of the SPDFs

$$\hat{D}_f(n, t) = \int_0^1 dx \; x^n D_f(x, t).$$

(11)

Eqs. (8) are linear non-homogeneous first order differential equations. The general solution for them is the sum of the general solution to the homogeneous equations (without the splitting term) and a particular solution to the non-homogeneous equations.

The homogeneous equation has the following general solution

$$\hat{D}(n_1, n_2, t, q) = e^{\gamma(n_1) t} A(n_1, n_2, q) e^{\gamma^T(n_2) t}.$$  

(12)

where the exponentials generate two collinear DGLAP evolutions since the DGLAP equations (7) in the Mellin representation,

$$\partial_t \hat{D}(n, t) = \gamma(n) \hat{D}(n, t),$$

(13)

have the solution

$$\hat{D}(n, t) = e^{\gamma(n) t} \hat{D}_0(n)$$

(14)

where $\hat{D}_0(n)$ is an initial condition. A particular solution to Eqs. (9) can now be found by making $A(n_1, n_2, q)$ time dependent. Substituting such an ansatz into Eqs. (9), we find the equation

$$\partial_t A(n_1, n_2, t, q) = e^{-\gamma(n_1) t} \hat{\gamma}(n_1, n_2) \hat{D}(n_1 + n_2, t) e^{-\gamma^T(n_2) t}$$

(15)
two partons originate from a point-like parton through its splitting, such a correlation (and the form factor) no longer reflects their correlation inside the nucleon, described by a non-perturbative form factor. On the other hand, if the matter of a physically motivated modeling. The basic idea is that for two partons originating from a nucleon, the dependence is not taken into account. Therefore, the first term in Eq. (17) has been postulated in [12] with the factorized q-dependence

\[ \tilde{D}^{(1)}(n_1, n_2, t, q) = e^{\gamma(n_1)t} \tilde{D}_0(n_1, n_2) e^{\gamma(n_2)t} F_{2g}(q) \]  

where \( F_{2g}(q) \) is the two-gluon nucleon form factor in the dipole form

\[ F_{2g}(q) = \frac{1}{(1 + q^2/m_g^2)^2} \]  

with \( m_g \) being an effective gluon mass. In principle, the form factor could depend on parton flavours, however, this dependence is not taken into account.
where for equal hard scales

\[ \tilde{D}^{(2)}(n_1, n_2, t, q) = \int_{t_0(q)}^{t} dt' e^{\gamma(n_1)(t-t')} \tilde{\gamma}(n_1, n_2) \tilde{D}(n_1 + n_2, t') e^{\gamma(n_2)(t-t')} \]  

(21)

where

\[ t_0(q) = \begin{cases} 
0 & \text{if } |q| \leq Q_0 \\
t(|q|) & \text{if } Q_0 < |q| \leq Q 
\end{cases} \]

(22)

and \( t(|q|) \) is given by Eq. 14. Thus, for \( |q| > Q_0 \), \(|q|\) is the scale from which the splitting starts, see Fig. 3. For \( |q| < Q_0 \), the relative loop momentum is small and may be neglected due to strong ordering in transverse parton momenta in the DGLAP approximation. Notice that \( |q| \leq Q \), which means that \( Q \) is the largest scale in the problem. Summarizing, after the transformation into the x-space, solution (17) is the sum of the two discussed contributions

\[ D(x_1, x_2, t, q) = D^{(1)}(x_1, x_2, t, q) + D^{(2)}(x_1, x_2, t, q). \]

(23)

The DPS cross section (11) can be written in terms of these contributions as the following sum

\[ \sigma_{AB} = \sigma_{AB}^{(11)} + \sigma_{AB}^{(12+21)} + \sigma_{AB}^{(22)} \]

(24)

where for equal hard scales

\[ \sigma_{AB}^{(ij)} = \frac{N}{2} \sum_{\text{flav}} \int dx_1 dx_2 dx_1' dx_2' \frac{d^2 q}{(2\pi)^2} \]

\[ \times D_{f_1 f_2}^{(ij)}(x_1, x_2, t, q) \tilde{\sigma}_{f_1 f_1}'(x_1, x_1', Q) \tilde{\sigma}_{f_2 f_2}'(x_2, x_2', Q) D_{f_1' f_2'}^{(ij)}(x_1', x_2', t, -q). \]

(25)

Each term has a clear interpretation, \( \sigma_{AB}^{(11)} \) is a contribution without parton splitting in the sense introduced by the analysis of the evolution equations, while \( \sigma_{AB}^{(22)} \) is a pure parton splitting contribution. There is also a mixed contribution, \( \sigma_{AB}^{(12+21)} \), with parton splitting from one hadron side only. Notice that \( \sigma_{AB}^{(11)} \) is proportional to

\[ \int \frac{d^2 q}{(2\pi)^2} F_{2g}^4(q) = \frac{m_g^2}{28\pi}, \]

where we assumed that \( Q \) which bounds \(|q|\) is much bigger than the mass \( m_g \). Similarly, \( \sigma_{AB}^{(12+21)} \) contains a similar factor with the second power of the form factor equal to \( m_g^2/12\pi \), while \( \sigma_{AB}^{(22)} \) is computed without the form factor. Therefore, we expect that the lack of strong suppression for large values of \(|q|\) makes the pure splitting contribution increasingly important with rising hard scale \( Q \). The importance of the splitting contribution has been also discussed in [10] in the context of cross sections which are not integrated over the transverse momenta of final state particles from hard interactions.

In the standard approach, the estimation of DPS cross sections is usually made with the cross section \( \sigma_{AB}^{(11)} \) in which the DPDFs are factorized into the product of two single PDFs. In such a case

\[ \sigma_{AB} = \frac{N}{2} \frac{\sigma_A \sigma_B}{\sigma_{\text{eff}}}, \]

(27)
where \( \sigma_A \) and \( \sigma_B \) are the single parton scattering cross sections and \( \sigma_{\text{eff}} \) is an effective cross section which is the inverse of the integral\(^{(28)}\)

\[
\frac{1}{\sigma_{\text{eff}}} = \int \frac{d^2 q}{(2\pi)^2} F_4^2(q). \tag{28}
\]

The value \( \sigma_{\text{eff}} \approx 15 \text{ mb} \), estimated by the CDF and D0 Collaborations from the DPS data\(^{38-40}\), gives gluon mass \( m_g \approx 1.5 \text{ GeV} \). In the next section we will show that the pure splitting contribution, \( \sigma_{AB}^{(22)} \), dominates the cross sections for the DPS electroweak boson production and formula\(^{(27)}\) can no longer be used for their estimation.

V. ELECTROWEAK BOSON PRODUCTION IN DPS

As an application of the presented formalism, we consider the DPS electroweak boson production \( W^+W^- \) and \( ZZ \) in the proton-proton scattering at the LHC center-of-mass energy \( \sqrt{s} = 14 \text{ TeV} \). The hard scale in such a case is boson mass, \( Q = M_{W,Z} \), see Fig. 4. In the collinear approach, the parton momentum fractions obey the condition

\[
x_{1,2} = \frac{Q}{\sqrt{s}} e^{y_{1,2}}, \quad z_{1,2} = \frac{Q}{\sqrt{s}} e^{-y_{1,2}}, \tag{29}
\]

where \( y_{1,2} \) are rapidities of the two produced bosons. The allowed values of rapidities, resulting from the relations: \( x_{1,2}, z_{1,2} \in [0,1] \) and \( x_1 + x_2 \), \( z_1 + z_2 \in [0,1] \), are shown in Fig. 5 (left) where the solid lines show constant values of \( x_1 + x_2 \) and \( z_1 + z_2 \) while the dashed lines correspond to constant ratios \( x_2/x_1 = z_1/z_2 \). Note that the point \( y_1 = y_2 = 0 \) in the middle corresponds to fully symmetric sharing of parton momentum fractions \( x \approx 0.5 \cdot 10^{-2} \).

In Fig. 5 (right) we show the DPS cross section

\[
\frac{d\sigma_{W^+W^-}}{dy_1 dy_2} = \frac{1}{\sigma_{\text{eff}}} \frac{d\sigma_{W^+}}{dy_1} \frac{d\sigma_{W^-}}{dy_2}, \tag{30}
\]
**FIG. 6**: The three contributions to $W^+W^-$ production cross section (24) in femtobarns. Contours of constant values are shown below.

Computation using the standard formula (27) with $\sigma_{\text{eff}} \approx 15$ mb and the single scattering cross sections $\sigma_{WW}^{\pm}$ (11), $\sigma_{WW}^{(12+21)}$ (12), $\sigma_{WW}^{(22)}$ (22) was used.

The three components of the cross section (24) are shown in Fig. 6. We see that the pure splitting contribution, $\sigma_{WW}^{(22)}$, is the dominant one. The standard contribution $\sigma_{WW}^{(11)}$ (which stays very close to the cross section from Fig. 5 (right)) is several times smaller. Thus, the factorized form (30) is no longer valid for the estimation of the DPS production of $W$ bosons since the pure splitting contribution dominates.
Integrand suppression at $y_1 = y_2 = 0$

$$d\sigma_{WW}/dy_1dy_2dq^2\ (fb/GeV^2)$$

FIG. 7: Form factor suppression in $d\sigma_{WW}/dy_1dy_2dq^2$ for the three indicated contributions. The upper limit for $q^2$ equals $M_W^2$.

This is illustrated in Fig. 7 where the cross sections $d\sigma_{WW}^{(ij)}/dy_1dy_2dq^2$ for the indicated components taken at $y_1 = y_2 = 0$ are shown. The strong form factor suppression for $q^2 > 10$ GeV$^2$ is clearly seen for the (11) and (12 + 21) components. Therefore, for such large hard scales like $Q^2 = M_{W,Z}^2 \gg 10$ GeV$^2$, the pure splitting contribution is much bigger than the standard one.

We repeat our analysis for the $ZZ$ boson production in DPS obtaining qualitatively the same results. In Fig. 8 we show the three components of the DPS cross section (24) in this case. The pure splitting contribution dominates again.

FIG. 8: The three contributions to $ZZ$ production cross section (24) in femtobarns. Contours of constant values are shown below.
The $W^+W^-$ and $ZZ$ splitting contributions are peaked around $y_1 \simeq y_2 \approx 0$. This is the configuration with parton momentum fractions, $x_1 \simeq x_2 \sim 10^{-2}$, for which the splitting contribution is determined by strongly rising single gluon distribution in the process $g \to q\bar{q}$.

VI. SUMMARY

In this paper, we have analyzed the DPS processes in the collinear approximation using evolution equations for the DPDFs. We have concentrated on the significance of the splitting terms in these equations for the DPS processes with large hard scale $Q \sim 100$ GeV. For the illustration, we have considered $W^+W^-$ and $ZZ$ boson production in DPS at the LHC energy 14 TeV. In order to compute the cross sections we have specified the dependence of of the DPDFs on the relative momentum $q$ following the model of Ryskin and Snigirev ([12]). In this model, the splitting contribution is not strongly suppressed at large values of $|q|$ like the standard contribution because it originates from the splitting of a point-like parton. We constructed a numerical program solving the evolution equations which allowed to compute the relevant cross sections. We have found that the splitting contribution dominates the electroweak boson production cross sections in DPS, making the standard way of their estimation no longer valid.

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VII. REFERENCES

[1] V. Shelest, A. Snigirev and G. Zinovev, Phys.Lett. B113, 325 (1982).
[2] G. Zinovev, A. Snigirev and V. Shelest, Theor.Math.Phys. 51, 523 (1982).
[3] R. K. Ellis, W. Furmanski and R. Petronzio, Nucl.Phys. B212, 29 (1983).
[4] A. Bukhvostov, G. Frolov, L. Lipatov and E. Kuraev, Nucl.Phys. B258, 601 (1985).
[5] A. M. Snigirev, Phys. Rev. D68, 114012 (2003), [hep-ph/0304172].
[6] V. L. Korotkich and A. M. Snigirev, Phys. Lett. B594, 171 (2004), [hep-ph/0404155].
[7] J. R. Gaunt and W. J. Stirling, JHEP 03, 005 (2010), [0910.4347].
[8] B. Blok, Y. Dokshitzer, L. Frankfurt and M. Strikman, Phys.Rev. D83, 071501 (2011), [1009.2714].
[9] F. A. Cecconopieri, Phys. Lett. B697, 482 (2011), [1011.6586].
[10] M. Diehl and A. Schafer, Phys. Lett. B698, 389 (2011), [1102.3081].
[11] J. R. Gaunt and W. J. Stirling, JHEP 1106, 048 (2011), [1103.1888].
[12] M. Ryskin and A. Snigirev, Phys.Rev. D83, 114047 (2011), [1103.3495].
[13] J. Bartels and M. G. Ryskin, 1105.1638.
[14] B. Blok, Y. Dokshitzer, L. Frankfurt and M. Strikman, Eur.Phys.J. C72, 1963 (2012), [1106.5533].
[15] M. Diehl, D. Ostermeier and A. Schafer, JHEP 1203, 089 (2012), [1111.0910].
[16] A. V. Manohar and W. J. Waalewijn, Phys.Rev. D85, 114009 (2012), [1202.3794].
[17] M. Ryskin and A. Snigirev, Phys.Rev. D86, 014018 (2012), [1203.2330].
[18] J. R. Gaunt, JHEP 1301, 042 (2013), [1207.0480].
[19] A. Snigirev, N. Snigireva and G. Zinovjev, 1403.6947.
[20] T. Sjostrand and M. van Zijl, Phys.Rev. D36, 2019 (1987).
[21] A. Del Favabro and D. Treleane, Phys. Rev. D61, 077502 (2000), [hep-ph/9911358].
[22] A. Kulesza and W. J. Stirling, Phys.Lett. B475, 168 (2000), [hep-ph/9912232].
[23] A. Del Fabbrabro and D. Treleane, Phys. Rev. D66, 074012 (2002), [hep-ph/0207311].
[24] E. Cattaruzza, A. Del Fabbro and D. Treleane, Phys. Rev. D72, 034022 (2005), [hep-ph/0507052].
[25] T. Sjostrand and P. Z. Skands, JHEP 0403, 053 (2004), [hep-ph/0402078].
[26] E. L. Berger, C. Jackson and G. Shaughnessy, Phys.Rev. D81, 014014 (2010), [0911.5348].
[27] J. R. Gaunt, C.-H. Kom, A. Kulesza and W. J. Stirling, Eur. Phys. J. C69, 53 (2010), [1003.3953].
[28] A. Snigirev, Phys.Rev. D81, 065014 (2010), [1001.0104].
[29] C. Kom, A. Kulesza and W. Stirling, Phys.Rev.Lett. 107, 082002 (2011), [1105.4186].
[30] E. L. Berger, C. Jackson, S. Quackenbush and G. Shaughnessy, Phys.Rev. D84, 074021 (2011), [1107.3150].
[31] C. Kom, A. Kulesza and W. Stirling, Eur.Phys.J. C71, 1802 (2011), [1109.0309].
[32] P. Bartalini et al., 1111.0469.
[33] M. Luszczak, R. Maciula and A. Szczurek, Phys.Rev. D85, 094034 (2012), [1111.3255].
[34] D. d’Enterria and A. M. Snigirev, Phys.Lett. B718, 1395 (2013), [1211.0197].
[35] D. d’Enterria and A. M. Snigirev, Phys.Lett. B727, 157 (2013), [1301.5845].
[36] R. Maciula and A. Szczurek, Phys.Rev. D87, 074039 (2013), [1301.4469].
[37] Axial Field Spectrometer Collaboration, T. Akesson et al., Z.Phys. C34, 163 (1987).
[38] CDF Collaboration, F. Abe et al., Phys.Rev.Lett. 79, 584 (1997).
[39] CDF Collaboration, F. Abe et al., Phys.Rev. D56, 3811 (1997).
[40] D0 Collaboration, V. Abazov et al., Phys.Rev. D81, 052012 (2010), [0912.5104].
[41] ATLAS Collaboration, G. Aad et al., New J.Phys. 15, 033038 (2013), [1301.6872].
[42] CMS Collaboration, S. Chatrchyan et al., JHEP 1403, 032 (2014), [1312.5729].
[43] ATLAS Collaboration, G. Aad et al., JHEP 1404, 172 (2014), [1401.2831].
[44] M. W. Krasny and W. Placzek, Acta Phys.Polon. B45, 71 (2014), [1305.1769].
[45] R. Kirschner, Phys.Lett. B84, 266 (1979).
[46] A. Snigirev, Phys.Atom.Nucl. 74, 158 (2011).
[47] A. Martin, W. Stirling, R. Thorne and G. Watt, Eur.Phys.J. C63, 189 (2009), [0901.0002].
[48] K. Golec-Biernat and E. Lewandowska, to be published.