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Heavy quark pair correlations in QCD

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Abstract. The azimuthal correlations of heavy quarks produced in high energy pp ($p\overline{p}$) collisions are calculated in the perturbative QCD without usual assumptions of the parton model. The virtual nature of the interacting gluons as well as their transverse motion and different polarizations are taken into account within the framework of the semihard processes theory describing the parton distributions in the region of the small Bjorken variable x. We give some predictions for the azimuthal correlations of charm and beauty hadrons produced at Tevatron-collider and LHC. Our approach can be of interest also for HERA energy region, since it shows a difference with the conventional parton model in the small x domain

1 Introduction

The investigation of the production of heavy quarks in high energy hadron processes provides a method for studying the internal structure of hadrons. Realistic estimates of the cross section of the heavy quark production as well as their correlations are necessary in order to plan experiments on existing and future accelerators. These predictions are usually obtained in the parton model framework and depend significantly on the quark and gluon structure functions. The last ones are more or less known experimentally from the data of HERA, but unknown at very small values of Bjorken variable $x < 10^{-4}$. However it is just the region that dominates in the heavy quark production at high energies.

The Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) evolution equation is usually applied to calculate the structure functions. It sums up in leading logarithm approximation (LLA) all the QCD diagram contributions proportional to $(\alpha_s \ln q^2)^n$ but it does not take into account the terms proportional to $(\alpha_s \ln 1/x)^n$, so this approximation does not give the correct asymptotic behaviour of the structure function in the small x region. For the correct description of these phenomena not only the terms of the form $(\alpha_s \ln q^2)^n$ have to be collected in the Feynman diagrams but also the terms $(\alpha_s \ln 1/x)^n$ and $(\alpha_s \ln q^2 \cdot \ln 1/x)^n$.

Another problem that appears at $x \sim 0$ is that of the absorption (screening) corrections which must stop the increase of the cross section when $x \to 0$ in accordance with the unitarity condition. It can be interpreted as saturation of the parton density. For relatively small virtuality $q^2 \leq q_0^2(x)$ the gluon structure function behaves as $xG(x,q^2) \sim q^2R^2$, so the cross section for the interaction of a point-like parton with a target, $\sigma \sim (1/q^2)xG(x,q^2) \sim R^2$, obeys the unitarity condition. The quantity $q_0^2(x)$ can be treated as a new infrared-cutoff parameter which plays the role of a typical transverse momentum of partons in the parton cascade of the hadron in semihard processes. The behaviour of $q_0(x)$ was discussed in [1]. It increases with $\log(1/x)$ and at $x = 0.01 - 0.001$ the values of $q_0(x)$ are about 2-4 GeV.

The predictions for the cross section value of heavy quark pair production are based usually on the parton model calculations [2,3]. In this model all particles involved are assumed to be on mass shell, having only longitudinal component of the momentum (so called collinear approximation) and the cross section is averaged over two transverse polarizations of the gluons. The virtualities q^2 of the initial partons are taken into account only through their densities. The latter are calculated in LLA through DGLAP evolution equation. The probabilistic picture of noninteracting partons underlies this way of proceeding. In the region where the transverse mass m_T of the produced heavy quark Q is close to $q_0(x)$ the dependence of the amplitude of the most important at high energy subprocess $gg \to \overline{Q}Q$ on the virtualities and polarizations of the gluons should be taken into account, i.e., the matrix elements of these subprocesses should be calculated more accurately than is usually done in the parton model.[1] The matrix elements of the QCD subprocesses accounting for the virtualities and polarizations of the gluons are very complicated. In our previous paper [9] we presented the results for main and simplest subprocess, $gg \to \overline{Q}Q$ ($\sim \alpha_s^2$) for hadroproduction and $\gamma g \to \overline{Q}Q$ ($\sim \alpha_s$) for photo- and electroproduction. The contributions of high-order subprocesses can be essential but our aim is to discuss the qualitative difference between our results and the parton

[1] The results accounting for the virtualities of incident gluons in a different approach based on k-factorization formulae can be found, for example, in [4–8]

model predictions that can be done on the level of low-order diagrams.

In this paper we calculate the azimuthal correlations [10] of heavy quarks produced in hadron-hadron collisions in the lowest ($\sim \alpha_s^2$) order.

2 Cross sections of heavy flavour production in QCD

The cross section of heavy quarks hadroproduction is given schematically by the graphs in Fig. 1. The main contribution to the cross section at small x is known to come from gluons. The lower and upper ladder blocks present the two-dimensional gluon distribution $\varphi(x,q_1^2)$ and $\varphi(x,q_2^2)$, which are the functions of the fraction (x and y) of the longitudinal momentum of the initial hadron and the gluon virtuality. Their distribution over x and transverse momenta q_T in hadron is given in semihard theory [1] by function $\varphi(x,q^2)$. It differs from the usual function $G(x,q^2)$:

$$xG(x,q^2) \;=\; \frac{1}{4\sqrt{2}\,\pi^3}\int_0^{q^2} \varphi(x,q_1^2)\,dq_1^2. \tag{1}$$

Such definition of $\varphi(x,q^2)$ makes possible to treat correctly the effects arising from the gluon virtualities. The exact expression for this function can be obtained as a solution of the evolution equation which, contrary to the parton model case, is nonlinear due to interactions between the partons in small x region.

In what follows we use Sudakov decomposition for quark momenta $p_{1,2}$ through the momenta of colliding hadrons p_A and p_B ($p_A^2 = p_B^2 \simeq 0$) and transverse ones $p_{1,2T}$:

$$p_{1,2} = x_{1,2}p_B + y_{1,2}p_A + p_{1,2T}. \tag{2}$$

The differential cross sections of heavy quarks hadroproduction have the form:[2]

$$\begin{aligned}\frac{d\sigma_{pp}}{dy_1^* dy_2^* d^2p_{1T} d^2p_{2T}} \;&=\; \frac{1}{(2\pi)^8}\frac{1}{(s)^2}\\ &\times \int d^2q_{1T}d^2q_{2T}\delta(q_{1T}+q_{2T}-p_{1T}-p_{2T})\\ &\times \frac{\alpha_s(q_1^2)}{q_1^2}\frac{\alpha_s(q_2^2)}{q_2^2}\varphi(q_1^2,y)\varphi(q_2^2,x)|M_{QQ}|^2.\end{aligned} \tag{3}$$

Here $s = 2p_Ap_B$, $q_{1,2T}$ are the gluons' transverse momenta and $y_{1,2}^*$ are the quarks' rapidities in the hadron-hadron c.m.s. frame,

$$\begin{aligned} x_1 &= \tfrac{m_{1T}}{\sqrt{s}}\,e^{-y_1^*},\; x_2 = \tfrac{m_{2T}}{\sqrt{s}}\,e^{-y_2^*},\; x = x_1 + x_2\\ y_1 &= \tfrac{m_{1T}}{\sqrt{s}}\,e^{y_1^*}, \quad\; y_2 = \tfrac{m_{2T}}{\sqrt{s}}\,e^{y_2^*},\; y = y_1 + y_2. \end{aligned} \tag{4}$$

$|M_{pp}|^2$ is the square of the matrix element for the heavy quark pair hadroproduction.

[2] We put the argument of α_S to be equal to gluon virtuality, which is very close to the BLM scheme [11]; (see also [12])

In LLA kinematic

$$q_1 \simeq yp_A + q_{1T},\; q_2 \simeq xp_B + q_{2T}. \tag{5}$$

so

$$q_1^2 \simeq -q_{1T}^2,\; q_2^2 \simeq -q_{2T}^2. \tag{6}$$

(The more accurate relations are $q_1^2 = -\frac{q_{1T}^2}{1-y}$, $q_2^2 = -\frac{q_{2T}^2}{1-x}$ but we are working in the kinematics where $x, y \sim 0$).

The matrix element M is calculated in the Born order of QCD without standart simplifications of the parton model since in the small x domain there are no grounds for neglecting the transverse momenta of the gluons q_{1T} and q_{2T} in comparision with the quark mass and the parameter $q_0(x)$. In the axial gauge $p_B^\mu A_\mu = 0$ the gluon propagator takes the form $D_{\mu\nu}(q) = d_{\mu\nu}(q)/q^2$,

$$d_{\mu\nu}(q) \;=\; \delta_{\mu\nu} -\; (q^\mu p_B^\nu \;+\; q^\nu p_B^\mu)/(p_Bq). \tag{7}$$

For the gluons in $t-$channel the main contribution comes from the so called 'nonsense' polarization $g_{\mu\nu}^n$, which can be picked out by decomposing the numerator into longitudinal and transverse parts:

$$\begin{aligned}\delta_{\mu\nu}(q) \;&=\; 2(p_B^\mu p_A^\nu \;+\; p_A^\mu p_B^\nu)/s \;+\; \delta_{\mu\nu}^T \;\approx\; 2p_B^\mu p_A^\nu/s\\ &\equiv\; g_{\mu\nu}^n.\end{aligned} \tag{8}$$

The other contributions are suppressed by the powers of s. Since the sum of the diagrams in Fig. 1a-1c is gauge invariant in the LLA, the transversality condition for the ends of gluon line enables one to replace p_A^μ by $-q_{1T}^\mu/x$ in the expression for $g_{\mu\nu}^n$. Thus we get

$$d_{\mu\nu}(q) \;\approx\; -2\,\frac{p_B^\mu q_T^\nu}{x\,s}, \tag{9}$$

or

$$d_{\mu\nu}(q) \;\approx\; 2\,\frac{q_T^\mu q_T^\nu}{xys}, \tag{10}$$

if we do such a trick for the vector p_B too. Both these equations for $d_{\mu\nu}$ can be used but for the form (9) one has to modify the gluon vertex slightly (to account for several ways of gluon emission – see [9]):

$$\begin{aligned}\Gamma_{eff}^\nu = \frac{2}{xys}\,&[(xys - q_{1T}^2)\,q_{1T}^\nu - q_{1T}^2 q_{2T}^\nu\\ &+2x\,(q_{1T}q_{2T})\,p_B^\nu].\end{aligned} \tag{11}$$

As a result the colliding gluons can be treated as aligned ones and their polarization vectors are directed along the transverse momenta. Ultimately, the nontrivial azimuthal correlations must arise between the transverse momenta p_{1T} and p_{2T} of the heavy quarks.

From the formal point of view there is a danger to loose the gauge invariance in dealing with the off mass shell gluons. Say, in the covariant Feynman gauge the new graphs (similar to the 'bremsstruhlung' from the initial or final quark line, as it is shown in Fig. 1d) may contribute in the central plato rapidity region. However this is not the fact. Within the “semihard” accuracy, when the

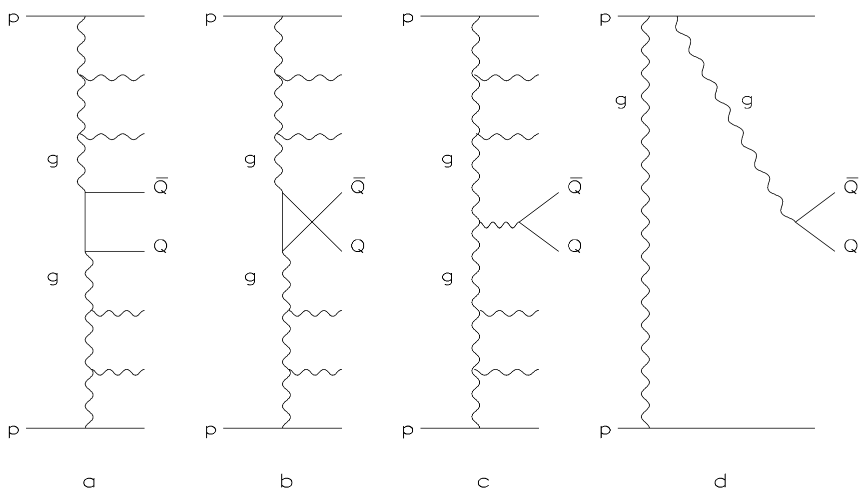

Fig. 1. Low order QCD diagrams for heavy quark production in pp ($p\bar{p}$) collisions via gluon-gluon fusion **a**–**c** and the diagram **d** formally violating the gauge invariance, that is restored within logarithmic accuracy

function $\phi(x, q^2)$ collects the terms of the form $\alpha_s^k(\ln q^2)^n \times(\ln(1/x))^m$ with $n+m \geq k$, the triple gluon vertex (11) includes effectively all the leading logarithmic contributions of the Fig. 1d type [13,14]. For instance, the upper part of the graph shown in Fig. 1d corresponds in terms of the BFKL equation to the t-channel gluon reggeization. Thus the final expression (17) is gauge invariant (except a small, non-logarithmic, $O(\alpha_s)$ corrections).

Although the situation considered here seems to be quite opposite to the parton model there is a certain limit, in which our formulae can be transformed into parton model ones. Let us consider the pp case and assume now that the characteristic values of quark momenta p_{1T} and p_{2T} are much larger than the values of gluon momenta q_{1T}, q_{2T}

$$< p_{1T} > \gg < q_{1T} > \,, \quad < p_{2T} > \gg < q_{2T} > \tag{12}$$

and one can keep only lowest powers of q_{1T}, q_{2T}. It means that we can put $q_{1T} = q_{2T} = 0$ everywhere in the matrix element M except the vertices. Introducing the polar coordinates

$$d^2 q_{1T} \;=\; \frac{1}{2}\, dq_{1T}^2 d\theta_1 \tag{13}$$

(and the same for q_{2T}) and performing angular integration with the help of the formula

$$\int_0^{2\pi} d\theta_1 q_{1T}^{\mu} q_{1T}^{\nu} \;=\; \pi\, q_{1T}^2 \delta_T^{\mu\nu} \tag{14}$$

we obtain

$$\int_0^{2\pi} d\theta_1 \frac{q_{1T}^{\mu}}{y}\frac{q_{1T}^{\nu}}{y} \int_0^{2\pi} d\theta_2 \frac{q_{2T}^{\lambda}}{x}\frac{q_{2T}^{\sigma}}{x} M_{\mu\nu}\overline{M}_{\lambda\sigma} = 2\pi^2 \frac{q_{1T}^2 q_{2T}^2}{(x\,y)^2}\,|M_{part}|^2. \tag{15}$$

Here M_{part} is just the matrix element in the parton model, since the result is the same as that calculated for the real (mass shell) gluons and averaged over transverse polarizations. Then we obtain the cross section (3) in the form [2, 3]:

$$\frac{d\,\sigma}{dy_1^* dy_2^* d^2 p_{1T}} = |M_{part}|^2 \frac{1}{(\hat{s})^2} \times \int \frac{\alpha_s(q_1^2)\varphi(y, q_{1T}^2)}{4\sqrt{2}\,\pi^3}\frac{\alpha_s(q_2^2)\varphi(x, q_{2T}^2)}{4\sqrt{2}\,\pi^3}\, dq_{1T}^2\, dq_{2T}^2 = \frac{\alpha_s^2(q^2)}{(\hat{s})^2}\,|M_{part}|^2\, xG(x, q_2^2)\, yG(y, q_1^2) \tag{16}$$

where $\hat{s} = xys$ is the mass square of $\overline{Q}Q$ pair.

However the assumption (12) is not fulfilled in a more or less realistic case. The transverse momenta of produced quarks as well as the gluon virtualities (QCD scale values) should be of the order of heavy quark masses.

3 Azimuthal correlations

Consider the distributions over the azimuthal angle ϕ, which is defined as an opening angle between two produced heavy quarks projected onto a plane perpendicular to the beam and defined as xy-plane. In the conventional LO parton model the sum of the produced heavy quarks momenta projected onto this plane is exactly zero and the angle between them is always 180^o. In the case of NLO parton model the sum of three momenta (two quark's and one gluon's) in the xy-plane should be equal to zero, therefore a non-trivial distribution over ϕ angle appears.

The theoretical as well as experimental investigation of such distributions are very important for the control of the level of our understanding of the considered processes. The problem is that in the case of one-particle inclusive distributions for heavy quark production in hadron collisions the sum of LO and NLO contributions of the parton model practically coinsides [15] with the LO contribution multiplied by so-called K-factor that is of the order of $1.5 \div 2$. So in the case of too small or too large NLO contribution the agreement with the experimental data can be achieved by fitting of one parameter, which can work as a normalization factor (say, QCD scale). In the case of azimuthal correlations all difference from the trivial $\delta(\phi-\pi)$

distribution comes from NLO contribution to the parton model.

The experimental data on these correlations are claimed (see [16] and references therein) to be in disagreement with the conventional parton model predictions for the cases of charm pair hadro- and photoproduction at fixed target energies. At the same time the experimental data on charm azimuthal correlations in different processes are in reasonable agreement with each other.

It was shown in [17], that the conventional approach can describe the experimental data on azimuthal correlations assuming the comparatively large intrinsic transverse momentum of incoming partons (k_T kick) equal to 1 GeV/c. However the large intrinsic transverse momentum significantly changes one-dimentional p_T distributions of heavy flavour hadrons which were earlier in good agreement with the data. It is presented in [17] that the agreement with the data can be re-established using the fragmentation function for quark-to-hadron transition process.

We shall try to consider another approach, when the azimuthal correlations of produced heavy quarks in the small-x region result from the diffusion of transverse momenta in the gluon evolution. This diffusion is described by the function $\phi(x, q^2)$, (1), determined in the semihard theory [1] by the derivative of the gluon distributions. The matrix element accounting for the gluon virtualities and polarizations is much more complicate than the parton model one. That is why we shall consider only LO contribution of the subprocess $gg \to Q\bar{Q}$. Note that our mechanism responsible for the appearance of the azimuthal correlations differs from the conventional parton model one, where the ϕ-distribution is related to the hard gluon emission. The NLO azimuthal correlations in our approach can be estimated as a kind of "convolution" of ϕ-distributions of our LO results and the NLO parton model distribution, however there exists a danger of double counting related to the problem of redefenition of NLO structure functions [15].

4 Results of calculations

Equation (3) enables to calculate straightforwardly the distributions over the azimuthal angle ϕ introduced in the previous section. Let the first heavy quark flights, for the definitnes, in the x direction of the xy-plane, i.e. $p_{1Tx} = p_{1T}$, $p_{1Ty} = 0$. In this case $\cos\phi = p_{2Tx}/p_{2T}$ and the distribution over ϕ can be easely evaluated during the integration of (3) using, say, VEGAS code.

Since the functions $\varphi(x, q_2^2)$ and $\varphi(y, q_1^2)$ are unknown at the small values of q_2^2 and q_1^2 we rewrite the integrals in (3) as

$$
\begin{aligned}
&\int d^2q_{1T} d^2q_{2T} \delta(q_{1T} + q_{2T} - p_{1T} - p_{2T}) \frac{\alpha_s(q_1^2)}{q_1^2} \frac{\alpha_s(q_2^2)}{q_2^2} \\
&\quad \times \varphi(q_1^2, y)\varphi(q_2^2, x)|M_{QQ}|^2 \\
&= 4\sqrt{2}\,\pi^3 \alpha_s(Q_0^2))^2\, xG(x, Q_0^2)\, yG(y, Q_0^2) \left(\frac{|M_{QQ}|^2}{q_1^2 q_2^2}\right)_{q_{1,2}\to 0} \\
&+ 4\sqrt{2}\,\pi^3 \alpha_s(Q_0^2) xG(x, Q_0^2) \int_{Q_0^2}^{\infty} dq_{1T}^2\, \delta(q_{1T} - p_{1T} - p_{2T}) \\
&\quad \times \frac{\alpha_s(q_1^2)}{q_1^2} \varphi(q_1^2, y) \left(\frac{|M_{QQ}|^2}{q_2^2}\right)_{q_2\to 0} \\
&+ 4\sqrt{2}\,\pi^3 \alpha_s(Q_0^2)\, yG(y, Q_0^2) \int_{Q_0^2}^{\infty} dq_{2T}^2\, \delta(q_{2T} - p_{1T} - p_{2T}) \\
&\quad \times \frac{\alpha_s(q_2^2)}{q_2^2} \varphi(q_2^2, x) \left(\frac{|M_{QQ}|^2}{q_1^2}\right)_{q_1\to 0} \\
&+ \int_{Q_0^2}^{\infty} d^2q_{1T} \int_{Q_0^2}^{\infty} d^2q_{2T}\, \delta(q_{1T} + q_{2T} - p_{1T} - p_{2T}) \\
&\quad \times \frac{\alpha_s(q_1^2)}{q_1^2} \frac{\alpha_s(q_2^2)}{q_2^2} \varphi(q_1^2, y)\varphi(q_2^2, x)|M_{QQ}|^2 \,, \qquad (17)
\end{aligned}
$$

where (1) is used[3]. Thus we obtain the sum of three different contributions: the first term in the r.h.s. of (17), $w_1(\phi)$; the sum of the second and third terms, $w_2(\phi)$; and, finally, the fourth term, $w_3(\phi)$. The first contribution, $w_1(\phi)$, is very similar to the conventional LO parton model expression, in which the sum of the produced heavy quarks momenta is exactly zero and the angle between them is always 180^o. However the angle between two heavy hadrons can be sligtly different from this value due to a hadronization process. To take it into account we assume that in the first contribution, where quarks are produced back-to-back, the probability to find a quark pair with azimuthal angle ϕ is determined by the expression

$$
w_1(\phi) = \frac{p_h}{\sqrt{p_h^2 + p_T^2}} \,, \qquad (18)
$$

where p_h is a transverse momentum in the azimuthal plane of a hadron coming from hadronization process. The two last contributions, $w_2(\phi)$ and $w_3(\phi)$, result in a more or less broad distribution over the angle between the produced quarks so we neglect here their small modification in a hadronization. The total probability is normalized to unity,

$$
\int_0^{\pi} d\phi(w_1(\phi) + w_2(\phi) + w_3(\phi)) = 1 \,. \qquad (19)
$$

It is necessary to repeat, that our approach is justified only at small x region, that is at the high enough initial energies. Unfortunately, the highest energy, where the experimental data on charm azimuthal correlations exist, is only $\sqrt{s} = 39$ GeV [16]. The values of q_1^2 and q_2^2 in (3) giving an essential contributions to the charm production cross section are not large enough at this energy, so at not very small Q_0^2 value the first contribution in (17) dominates and our predictions coinside practically with the results of LO parton model.

At the higher energies the second and third contributions become larger and here we get some difference with the conventional parton model. The predictions of the charm pair azimuthal correlations with GRV HO [18] parton distributions at $\sqrt{s} = 1800$ GeV and $Q_0^2 = 4$ GeV^2

[3] The value of Q_0^2 should not be mixed here with the function $q_0^2(x)$ discussed in the Introduction

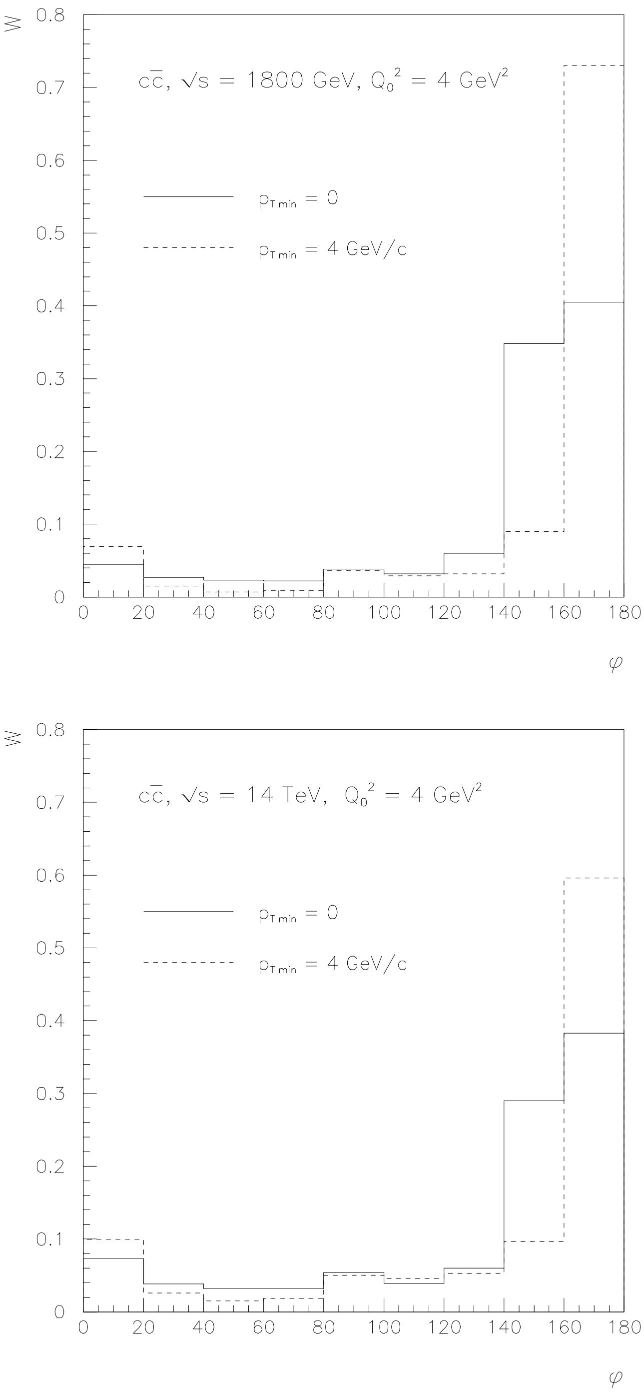

Fig. 2. The calculated azimuthal correlations of charm pair production in pp ($\overline{p}p$) collisions at $\sqrt{s}$ = 1800 GeV **a** and $\sqrt{s}$ = 14 TeV **b** for all events (*solid histograms*) and for the events, where both charm particle have $p_T > p_{Tmin}$ (*dashed hystograms*)

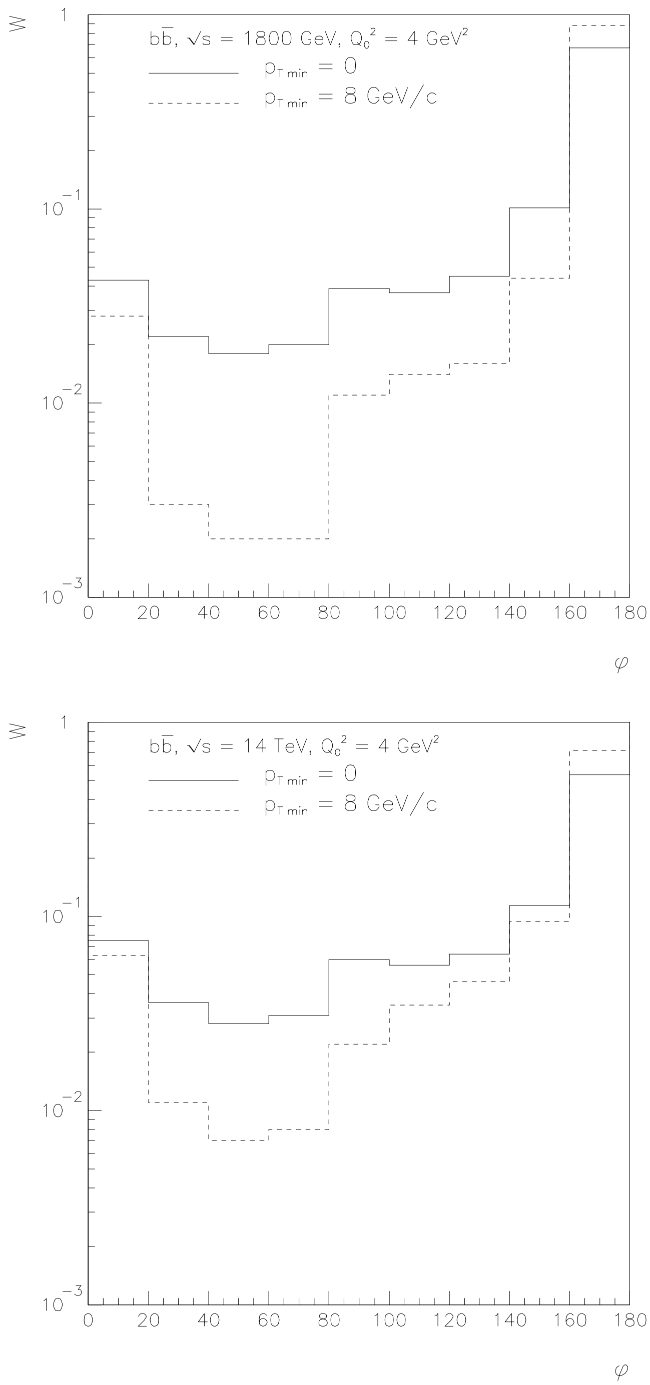

Fig. 3. The calculated azimuthal correlations of beauty pair production in pp ($\overline{p}p$) collisions at $\sqrt{s}$ = 1800 GeV **a** and 14 TeV **b** for all events (*solid histograms*) and for the events, where both beauty particle have $p_T > p_{Tmin}$ (*dashed histograms*)

are presented in Fig. 2a. The solid hystogram shows the results for all produced charm particles. However in this case some hadronization mechanism can contribute to the azimuthal correlations of charmed hadrons with small relative momenta. To decrease this uncertainity we present by dashed histogram the same results for the events, where both charmed hadrons have the transverse momenta $p_T > p_{Tmin}$, p_{Tmin} = 4 GeV/c.

The results at the energy $\sqrt{s}$ = 14 TeV are shown in Fig. 2b. Note that the predicted azimuthal correlations are energy dependent, they become more broad with the increase of the initial energy, although this energy dependence is weak, especially in the case of $p_{Tmin} = 0$. Our

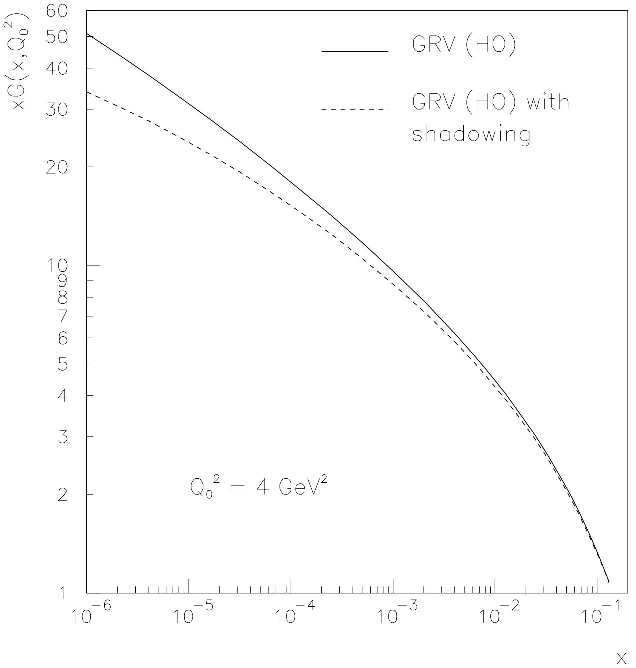

Fig. 4. Gluon structure function of the nucleon [18] without (*solid curve*) and with (*dashed curve*) shadow correction, (19)

ϕ-distributions seem to be more broad than the results of the parton model calculations.

The similar predictions for the beauty pair azimuthal correlations at the energies $\sqrt{s} = 1800$ GeV and 14 TeV are presented in Fig. 3 for two values of p_{Tmin}: 0 and 8 GeV/c. At the both values of p_{Tmin} the distributions become more broad, when the energy increases.

Contrary to the parton model results we predict two peaks for all cases: the standard one at $\phi = 180^o$ and the second peak with smaller altitude, at $\phi = 0^o$ originated from the contribution of the diagram Fig. 1c.

The HERA experimental data on $F_2(x, Q^2)$ at small x, actually $10^{-4} < x < 10^{-2}$, show a singular x-behaviour at moderate Q^2 (say, $Q^2 \sim 10^1$ GeV^2). Both ZEUS [19] as well as H1 [20] Coll. data can be parametrized as $x^{-\delta}$ with $\delta = 0.1 \div 0.25$. Such a behaviour at $x \to 0$ is in evident contradiction with the unitarity and has to be stopped by a shadow mechanism [1,21,22].

To see the possible influence of shadow effects on the azimuthal correlations we make a simplest assumption, that the shadowing modifies the gluon distribution in such a way, that the real distribution can be written as

$$xG(x, Q^2) = \frac{xG_0(x, Q^2)}{1 + \epsilon xG_0(x, Q^2)} \,, \qquad (20)$$

where $\epsilon \ll 1$ and $xG_0(x, Q^2)$ is the bare GRV (HO) gluon distribution [18]. Both $xG(x, Q^2)$ and $xG_0(x, Q^2)$ distributions are shown in Fig. 4 for the value $\epsilon = 0.01$. One can see that the difference between them at the smallest x, where the data exist ($x \sim 10^{-4}$), is about 10 %.

The calculated results for azimuthal correlations of heavy flavour pairs with "shadowed" gluon distributions,

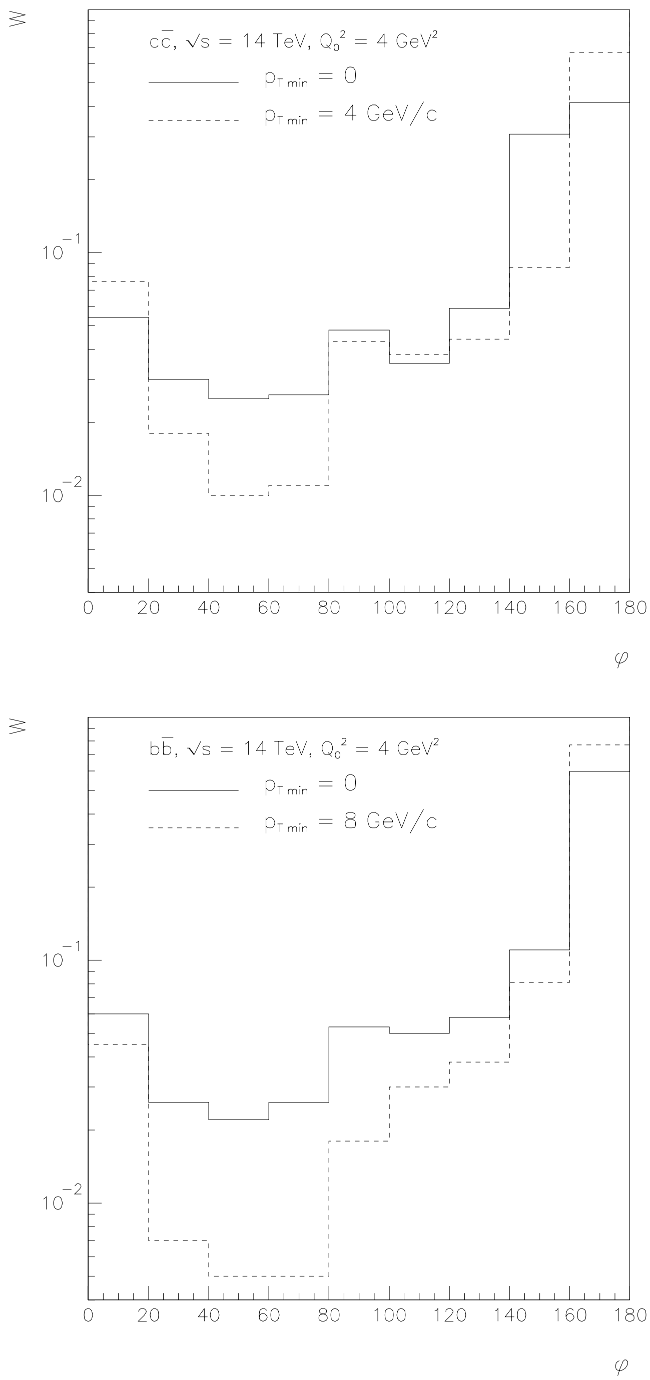

Fig. 5. The calculated azimuthal correlations of charm **a** and beauty **b** pair production in pp ($\overline{p}p$) collisions at $\sqrt{s} = 14$ TeV for all events (*solid histograms*) and for the events, where both heavy flavour particle have $p_T > p_{Tmin}$ (*dashed histograms*)

(20), are presented in Fig. 5 and show that the shadowing does not significantly affect our results.

5 Conclusion

The above discussion shows that the accounting for the virtual nature of the interacting gluons as well as their transverse motion and different polarizations result in a qualitative differences with the parton model predictions. The reason is that the NLO parton model allows the discussed distributions to differ from δ-functions only due to a possibility of one hard gluon emission. In contrast, our approach effectively incorporates the emission of all evolution gluons via the phenomenological gluon distribution. It gives a possibility to distinct experimentally between these two approaches. In the parton model practically all events with ϕ significantly smaller than π should be accompanied by hard gluon jet, whereas in our approach we can expect such jet only with some (not very large) probability.

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