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PERIPHERAL ULTRARELATIVISTIC HEAVY-ION
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**Deduction of the in-medium gluon distribution
from photon-gluon fusion processes
in peripheral ultrarelativistic heavy-ion collisions**

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ABSTRACT

The photon-gluon fusion process into a pair of heavy quarks is studied for peripheral Pb+Pb collisions at LHC energies. The double differential cross section with respect to the invariant mass and the rapidity of the produced quark pair at zero rapidity is directly proportional to the gluon distribution in the nuclear medium. Differential cross sections for the b -quark production lie well in the $\mu\text{barn}/\text{GeV}$ regime which will allow the deduction of the gluon distribution at low x . Rapidity cuts for the actual detectors are also considered.

1 Introduction

Very strong electromagnetic fields prevail in ultrarelativistic heavy-ion collisions such as Au+Au at RHIC or Pb+Pb at LHC. It has been pointed out previously [1] that these sources may be used to study two-photon production of heavy exotic mesons, glueballs and, eventually, supersymmetric particles, Higgs bosons and magnetic monopoles. Those processes occur in peripheral collisions and could be studied for "free" at the ultrarelativistic heavy-ion colliders, which will be mainly devoted to investigate a possible phase transition to a Quark-Gluon-Plasma in central collisions.

Besides purely electromagnetic processes also photon-hadron reactions occur in peripheral ultrarelativistic heavy-ion collisions. A photon has a rho meson component, which interacts with the other photon; however, this process is only of minor importance because the available photon energies are still too small [2]. On the other hand a photon from one of the two nuclei may penetrate directly into the other nucleus and interact with one of its quarks or gluons. In this publication we will concentrate on such a photon-gluon fusion into a pair of heavy quarks.

It has been suggested previously [3, 4] that the photon-gluon fusion mechanism into a heavy quark pair ($c\bar{c}$ or $b\bar{b}$) occurring in peripheral ultrarelativistic heavy-ion collisions represents a good opportunity to extract information about the gluon distribution in a nuclear medium. The kinematics is rather clean. The transverse momentum of the outgoing two quark jets contain only negligible contributions from the ingoing photon and gluon of the order 250 MeV. The photon has a maximal transverse momentum $|\vec{k}_\perp|_{\gamma,\max} \approx 1/R \approx 30$ MeV corresponding to the inverse nuclear radius (of Pb) whereas the gluon has a typical transverse momentum $|\vec{p}_\perp|_{g,\max} \approx 1/r \approx 200$ MeV corresponding to the inverse radius of a nucleon. The trigger would thus be two quark jets with opposite $\vec{p}_\perp$ and nothing else. – Information on the gluon content of a bound nucleon which could be gained from investigating the process $\gamma g \rightarrow q\bar{q}$ in peripheral ultrarelativistic heavy-ion collisions might represent an important supplement for a better understanding of the original EMC effect [5]. The measured in-medium quark structure functions suggest that there might be a gluon enhancement in bound nucleons. The gluon distribution could be quite different from what it is in free nucleons, which makes it in any case very worth while to measure it directly.

We have organized this publication as follows: For the theoretical formulation of the photon-gluon fusion process in peripheral heavy-ion collisions we adopt the picture of an equivalent photon-partonic gluon interaction; it turns out that the double differential cross-section $\frac{d^2\sigma}{dMdy}|_{y=0}$ with respect to invariant mass and rapidity of the produced quark pair at zero rapidity is directly proportional to the in-medium gluon distribution. Rapidity cuts as for the proposed ALICE detector at LHC [6] are also worked in. All of this is covered in section 2. In section 3 we study the sensitivity of the cross section on the photon absorption part and present methods how to deduce in-medium gluon distributions from various differential cross sections with and without rapidity cuts. Here

we have restricted our considerations to $b\bar{b}$ production in peripheral Pb+Pb collisions at LHC energies. In the conclusions (section 4) we sketch some further refinements, which will be worked out in the near future.

2 Theoretical formulation

We adopt a picture for the photon-gluon fusion process which is in close analogy to the equivalent photon method [7, 8, 9]. Two ultrarelativistic heavy-ions collide with an impact parameter $\vec{b}$. The transverse electromagnetic field of one nucleus can be simulated by an equivalent swarm of photons. One of these photons penetrates into the other nucleus and fuses with a gluon. This gluon belongs to a nucleon, which has a transverse distance $\vec{x}_\perp$ to the center of the nucleus; with respect to the center of the nucleon the gluon has the distance $\vec{x}'_\perp$ in the transverse plane (confer fig. 1).

The cross section for photon-gluon fusion into a pair of heavy quarks $q\bar{q}$ in a relativistic heavy-ion collision can be cast in the following form:

$$\sigma = \int_{2R}^{\infty} d^2b \int_0^{\infty} d\omega_1 \int_0^p d\omega_2 \left[\tilde{n}_{\parallel}(\omega_1, \omega_2; \vec{b}) \sigma_{\gamma g \rightarrow q\bar{q}}^{\parallel}(\omega_1, \omega_2) + \tilde{n}_{\perp}(\omega_1, \omega_2; \vec{b}) \sigma_{\gamma g \rightarrow q\bar{q}}^{\perp}(\omega_1, \omega_2) \right] . \quad (1)$$

The integration over the impact parameter is restricted to peripheral collisions, where b is larger than twice the nuclear radius R . The photon-gluon distribution functions $\tilde{n}_{\parallel}$ and $\tilde{n}_{\perp}$ depend on the photon frequency ω_1 and the gluon frequency $\omega_2 = xp$, where x represents the fraction of the nucleon momentum p . They read:

$$\tilde{n}_{\parallel}(\omega_1, \omega_2; \vec{b}) = \int d\vec{x}_\perp \int d\vec{x}'_\perp n(\omega_1, \vec{x}_\perp + \vec{x}'_\perp - \vec{b}) N(\vec{x}_\perp) g(\omega_2, \vec{x}'_\perp) \left(\frac{(\vec{x}_\perp + \vec{x}'_\perp - \vec{b}) \cdot \vec{x}'_\perp}{|\vec{x}_\perp + \vec{x}'_\perp - \vec{b}| |\vec{x}'_\perp|} \right)^2 ; \quad (2)$$

for $\tilde{n}_{\perp}$ the dot product in the last bracket has to be replaced with a cross product. Here $n(\omega_1, \vec{x}_\perp + \vec{x}'_\perp - \vec{b})$ represents the equivalent photon distribution, $g(\omega_2, \vec{x}'_\perp)$ stands for the gluon distribution of a nucleon and $N(\vec{x}_\perp)$ is the nucleon number distribution projected onto the transverse plane; detailed expressions are given later. The geometrical factor given by the bracket in eq. (2) reflects the dependence of the photon-gluon interaction on their mutual polarisation and has been derived in ref. [10] for photon-photon fusion processes; for the parallel case ($\parallel$) the polarisations are parallel whereas for the perpendicular case ($\perp$) the polarisations are perpendicular. $\sigma_{\gamma g \rightarrow q\bar{q}}^{\parallel}$ and $\sigma_{\gamma g \rightarrow q\bar{q}}^{\perp}$ represent the corresponding polarised photon-gluon cross sections and add up to yield the total unpolarised photon-gluon cross section

$$\sigma_{\gamma g \rightarrow q\bar{q}} = \frac{1}{2} \left(\sigma_{\gamma g \rightarrow q\bar{q}}^{\parallel} + \sigma_{\gamma g \rightarrow q\bar{q}}^{\perp} \right) . \quad (3)$$

At this stage we make the usual approximation that the dependence of the gluon distribution on the longitudinal and transverse coordinates factorizes. We parametrize the gluon distribution as

$$g(\omega_2, \vec{x}'_\perp) = g(\omega_2) h(|\vec{x}'_\perp|) \quad (4)$$

with $\int d\vec{x}'_\perp h(|\vec{x}'_\perp|) = 1$. As a model for $h(|\vec{x}'_\perp|)$ we use, again as usually done, a gaussian distribution

$$h(|\vec{x}'_\perp|) = \frac{1}{\pi r^2} \exp\left(-\frac{\vec{x}'_\perp{}^2}{r^2}\right) \quad , \quad (5)$$

where r represents the radius of a nucleon. For peripheral collisions the inequality

$$|\vec{x}'_\perp| \ll |\vec{x}_\perp - \vec{b}| \quad (6)$$

holds; only for very close peripheral collisions it might not be fulfilled. The photon-gluon distribution of eq. (2) then becomes

$$\begin{aligned} \hat{n}_\parallel(\omega_1, \omega_2; \vec{b}) &= \int d\vec{x}_\perp n(\omega_1, \vec{x}_\perp - \vec{b}) N(\vec{x}_\perp) g(\omega_2) \int d\vec{x}'_\perp h(|\vec{x}'_\perp|) \left(\frac{(\vec{x}_\perp - \vec{b}) \cdot \vec{x}'_\perp}{|\vec{x}_\perp - \vec{b}| |\vec{x}'_\perp|} \right)^2 \\ &= \frac{1}{2} g(\omega_2) \int d\vec{x}_\perp N(\vec{x}_\perp) n(\omega_1, \vec{x}_\perp - \vec{b}) \\ &= \hat{n}_\perp(\omega_1, \omega_2; \vec{b}) \quad , \end{aligned} \quad (7)$$

so that the total cross section (1) transforms into

$$\sigma = \int_{2R}^\infty d^2b \int_0^\infty d\omega_1 \int_0^p d\omega_2 g(\omega_2) \left(\int d\vec{x}_\perp N(\vec{x}_\perp) n(\omega_1, \vec{x}_\perp - \vec{b}) \right) \sigma_{\gamma g \rightarrow q\bar{q}}(\omega_1, \omega_2) \quad . \quad (8)$$

This form of the cross section has been used in two previous publications [3, 4].

For a later convenience it is appropriate to introduce a change of integration variables. Instead of the frequencies ω_1 and ω_2 we will now use the center of momentum rapidity $y = \frac{1}{2} \ln(\omega_1/\omega_2)$ and the invariant mass $M = \sqrt{4\omega_1\omega_2}$ of the photon-gluon system. The total cross section (8) then reads:

$$\begin{aligned} \sigma &= \int_{2m_q}^\infty \frac{M}{2} dM \int_{\ln \frac{M}{2p}}^\infty dy g\left(\frac{M}{2} e^{-y}\right) \left[\int_{2R}^\infty d^2b \int d\vec{x}_\perp N(\vec{x}_\perp) n\left(\frac{M}{2} e^y, \vec{x}_\perp - \vec{b}\right) \right] \times \\ &\quad \times \sigma_{\gamma g \rightarrow q\bar{q}}\left(\frac{M}{2} e^y, \frac{M}{2} e^{-y}\right) \quad . \end{aligned} \quad (9)$$

In particular from the last expression it becomes obvious that the double differential cross section $\frac{d^2\sigma}{dM dy}$ with respect to the invariant mass and to the center of momentum rapidity of the produced pair of heavy quarks is directly proportional to the gluon distribution $g(\omega_2)$ of a nucleon inside a nucleus. In other words, if $\frac{d^2\sigma}{dM dy}$ could be determined experimentally and the photonic part of the cross section would be precisely known theoretically, the in-medium gluon distribution could be deduced.

For the equivalent photon distribution $n(\omega, \vec{b})$, which describes the number of photons with energy ω at a distance $|\vec{b}|$ from the center of the nucleus, we use [10]

$$n(\omega, \vec{b}) = \frac{4Z^2\alpha}{\omega} \left| \int \frac{d^2k_\perp}{(2\pi)^2} \vec{k}_\perp \frac{F(\vec{k}_\perp^2 + \omega^2/\gamma^2)}{\vec{k}_\perp^2 + \omega^2/\gamma^2} e^{-i\vec{b}\vec{k}_\perp} \right|^2. \quad (10)$$

Here α is the fine structure constant, Z is the nuclear charge number and $\gamma = (1-v^2/c^2)^{1/2}$ is the Lorentz contraction factor. Two different versions will be used for the nuclear charge formfactor F , namely that of a homogeneously charged sphere and that of a Gaussian charge distribution; confer ref. [10] for more details. – We use three different variations for the nucleon number density $N(\vec{x}_\perp)$. A very crude ansatz is that of a point nucleus with

$$N_{\text{pn}}(\vec{x}_\perp) = A_{\text{eff}} \delta(\vec{x}_\perp) \quad , \quad (11)$$

where $A_{\text{eff}} = A^{0.9}$ simulates nuclear shadowing [11]. A more relevant distribution, already used in ref. [4], is given by the projection of a uniform sphere onto a transverse plane

$$N_{\text{hs}}(\vec{x}_\perp) = 2\sqrt{R^2 - \vec{x}_\perp^2} \frac{A_{\text{eff}}}{\frac{4\pi}{3}R^3} \quad ; \quad (12)$$

it captures the large nonlinearities of the equivalent photon distribution (10) at small impact parameters. Still, nuclear shadowing is not treated correctly, since a quantity such as an absorption length is missing in the expression (12). Therefore we propose

$$N_{\text{cs}}(\vec{x}_\perp) = \frac{A}{\frac{4\pi}{3}R^3} \begin{cases} 2\sqrt{R^2 - \vec{x}_\perp^2} & \text{for } R \geq |\vec{x}_\perp| \geq \sqrt{R^2 - r_{12}/4} \\ r_{12} & \text{for } 0 \leq |\vec{x}_\perp| \leq \sqrt{R^2 - r_{12}/4} \end{cases} \quad (13)$$

Here a sphere with radius R is cut from another sphere with radius R at a center distance r_{12} and the leftover is projected onto a transverse plane (confer fig. 2). The distance $r_{12} = 0.82941424 \cdot R$, which simulates the absorption length, is determined in such a way that $\int N_{\text{cs}}(\vec{x}_\perp) d\vec{x}_\perp = A_{\text{eff}}$. Even the expression (13) is still very simple since it does not account for a frequency dependence of the absorption length. – It should be clear that the actual quantity we will deduce from an experiment is the product of the effective nucleon number and the gluon distribution and not the gluon distribution alone. The dependence of the cross section on the actual form of $N(\vec{x}_\perp)$ will be small; therefore we have introduced the three parametrisations given above to test this sensitivity. The only unknown here is the effective nucleon number A_{eff} , which determines the absolute value of the cross section but not the shape of the relevant differential cross section; it could be deduced from the asymptotic large x behaviour of the gluon distribution in comparison with other independent experimental data. Here we simply assume to know A_{eff} .

In second order perturbation theory the elementary polarised cross section for photon-gluon fusion are given by

$$\sigma_{\gamma g \rightarrow q\bar{q}}^{\parallel} = \frac{2\pi\alpha\alpha_s(M^2)e_q^2}{2s} \left\{ \left[\frac{5}{2} + \beta^2 - \frac{3}{2}\beta^4 \right] \ln \left(\frac{1+\beta}{1-\beta} \right) - \beta(5 - 3\beta^2) \right\} \Theta(s - 4m_q^2) \quad ,$$

$$\sigma_{\gamma g \rightarrow q\bar{q}}^{\perp} = \frac{2\pi\alpha\alpha_s(M^2)e_q^2}{2s} \left\{ \left[\frac{7}{2} - \beta^2 - \frac{1}{2}\beta^4 \right] \ln \left(\frac{1+\beta}{1-\beta} \right) - \beta(3-\beta^2) \right\} \Theta(s-4m_q^2), \quad (14)$$

where $\beta = (1 - 4m_q^2/s)^{1/2}$ and $s = 4\omega_1\omega_2 = M^2$. The corresponding elementary unpolarised cross section $\sigma_{\gamma g \rightarrow q\bar{q}}$ follows from eqs. (3) and (14).

The elementary cross sections indicated above are relevant, if the detector is able to identify both outgoing quarks. As for the proposed ALICE detector at LHC, it will only have a coverage of $\pi/4 \leq \theta \leq 3\pi/4$ for the polar angle [6], which translates into

$$|\tan \theta| \geq a_{\text{cut}} = 1 \quad ; \quad (15)$$

this covers only a fraction of the full phase space of the outgoing quarks. Therefore it is very important to reduce the cross section (9) only to those events, where both quarks fall into the polar angle coverage and thus are detected. As a consequence we rewrite the elementary second order photon-gluon fusion cross sections as

$$\sigma_{\gamma g \rightarrow q\bar{q}} = \int_{-y'}^{y'} dy' \frac{d\sigma_{\gamma g \rightarrow q\bar{q}}}{dy'}, \quad (16)$$

where y' represents the single rapidity of one outgoing quark in the photon-gluon center of momentum system and $\bar{y}' = \frac{1}{2} \ln \left(\frac{1+\beta}{1-\beta} \right)$. For the detection of both produced quarks we then require the cut-condition (15), which can also be expressed in terms of rapidities as:

$$\frac{\sinh |y \pm y'|}{\sqrt{1 - (1 - \beta^2) \cosh^2 y'}} \leq \frac{1}{a_{\text{cut}}} \quad . \quad (17)$$

This leads to a modification of the expression (9):

$$\begin{aligned} \sigma^{\text{cut}} &= \int_{2m_q}^{\infty} \frac{M}{2} dM \int_{-y_{\text{cut}}^*}^{y_{\text{cut}}} dy g \left(\frac{M}{2} e^{-y} \right) \left[\int_{2R}^{\infty} d^2b \int d\vec{x}_{\perp} N(\vec{x}_{\perp}) n \left(\frac{M}{2} e^y, \vec{x}_{\perp} - \vec{b} \right) \right] \times \\ &\quad \times \sigma_{\gamma g \rightarrow q\bar{q}}^{\text{cut}}(M, y, a_{\text{cut}}) \end{aligned} \quad (18)$$

with

$$y_{\text{cut}} = \text{arcsinh} \left(\frac{\beta}{a_{\text{cut}}} \right) \quad , \quad (19)$$

$$-y_{\text{cut}}^* = \begin{cases} -y_{\text{cut}} & \text{for } \ln \frac{M}{2p} \leq -y_{\text{cut}} \\ \ln \frac{M}{2p} & \text{for } -y_{\text{cut}} \leq \ln \frac{M}{2p} \leq y_{\text{cut}} \\ y_{\text{cut}} & \text{for } \ln \frac{M}{2p} \geq y_{\text{cut}} \end{cases} \quad (20)$$

and

$$\sigma_{\gamma g \rightarrow q\bar{q}}^{\text{cut}} = \frac{2\pi\alpha\alpha_s e_q^2}{s} \left\{ (3 - \beta^4) \bar{y}'_{\text{cut}} - \tanh \bar{y}'_{\text{cut}} - (1 - \beta^2)^2 \sinh \bar{y}'_{\text{cut}} \cosh \bar{y}'_{\text{cut}} \right\} \quad , \quad (21)$$

for $\bar{y}'_{\text{cut}} \leq \bar{y}'$, or

$$\sigma_{\gamma g \rightarrow q\bar{q}}^{\text{cut}} = \sigma_{\gamma g \rightarrow q\bar{q}} \quad , \quad (22)$$

for $\bar{y}'_{\text{cut}} \geq \bar{y}'$, respectively, where

$$\bar{y}'_{\text{cut}} = \frac{1}{2} \ln \left(\frac{(1 + \beta^2 + a_{\text{cut}}^2) + 2\sqrt{\beta^2(1 + a_{\text{cut}}^2) - a_{\text{cut}}^2(1 - \beta^2) \sinh^2 y}}{(1 - \beta^2 + a_{\text{cut}}^2 e^{2|y|})} \right) \quad (23)$$

follows from the requirement (17).

3 Results

3.1 The photon part of the cross section

First of all we investigate the differences between the cross section (1) and its approximated version (8). Since the underlying approximation (6) depends on the impact-parameter, we compare the corresponding differential cross sections with respect to the impact-parameter. For the calculations we have used a Gaussian nuclear charge form-factor entering the equivalent photon distribution (10), the nucleon number distribution (13) of a cut sphere, the parametrisation of the gluon distribution from ref. [12] and the Born expressions (14) and (3) for the elementary photon-gluon fusion cross section into a $b\bar{b}$ -quark pair. The Lorentz contraction factor has been set equal to $\gamma = 3000$ as proposed for Pb+Pb collisions at LHC energies and the mass of the b -quark has been chosen to be $m_b = 4.5$ GeV. Various differential cross sections with respect to the impact-parameter are shown in fig. 3. The solid line represents $d\sigma/db$ of eq. (8) and is identical to $d\sigma/db$ of eq. (1) in the limit $r \rightarrow 0$, where the nucleon radius approaches zero, so that the approximation (6) is perfectly fulfilled for all impact-parameters. For a realistic nucleon radius of $r = 0.8$ fm (dotted line) the differential cross section $d\sigma/db$ of eq. (1) is practically identical to $d\sigma/db$ of eq. (8) for larger impact-parameters; small deviations occur only at small impact-parameters, where the former differential cross section is somewhat smaller than the latter one. Assuming a larger nucleon radius of $r = 1.2$ fm (dashed line), this tendency is increased. Nevertheless, for impact-parameters larger than about twice the nuclear radius $R = 7.1$ fm of Pb, which corresponds to peripheral collisions, all versions of $d\sigma/db$ are practically the same. Consequently we can use the approximated form (8) of the cross section without introducing sizeable errors.

The double differential cross section (8) with respect to the photon- and gluonfrequencies ω_1 and ω_2 is directly proportional to the gluon distribution $g(\omega_2)$. The proportionality factor is given by the ω_1 and ω_2 dependent function

$$\xi(\omega_1, \omega_2) = \int_{2R}^{\infty} d^2b \int d\vec{x}_{\perp} N(\vec{x}_{\perp}) n(\omega_1, \vec{x}_{\perp} - \vec{b}) \sigma_{\gamma g \rightarrow q\bar{q}}(\omega_1, \omega_2) \quad , \quad (24)$$

which we will simply denote as a structure function; here the nucleon number density $N(\vec{x}_\perp)$, the impact-parameter dependent equivalent photon distribution $n(\omega_1, \vec{x}_\perp - \vec{b})$ as well as the elementary photon-gluon fusion cross section $\sigma_{\gamma g \rightarrow q\bar{q}}$ enter. Suppose, the double differential cross section could be measured independently, then in order to extract the gluon distribution the structure function $\xi(\omega_1, \omega_2)$ has to be controlled theoretically. Therefore we will study the sensitivity of $\xi(\omega_1, \omega_2)$ upon changes of some of the quantities entering the expression (24) next.

The impact-parameter dependent equivalent photon distribution (10) is well known. We will use the nuclear charge formfactor of a Gaussian charge distribution. Slight modifications result for a formfactor of a homogeneously charged sphere for impact-parameters smaller than about the nuclear radius [10]. Since we are only interested in impact-parameters larger than twice the nuclear radius, these modifications will turn out to be very small.

Next, we study the influence of the nucleon number density $N(\vec{x}_\perp)$. At very large impact-parameters the three different ansätze (11)–(13) should yield the same results. Therefore and for simplicity we consider the differential cross section (8) with respect to the impact parameter and not the impact-parameter dependence of the structure function (24), which also depends on the frequencies ω_1 and ω_2 . We choose the Born expression (14) for the elementary photon-gluon fusion cross section (3) for $b\bar{b}$ quark pair production with $m_b = 4.5$ GeV and the parametrisation of the gluon distribution from ref. [12] and consider a Pb+Pb collision with $\gamma = 3000$. The results for $d\sigma/db$ are depicted in fig. 4. In comparison to the point-nucleus distribution (11) the impact-parameter dependent cross section with the uniform sphere distribution (12) is larger for impact-parameters between about 10 and 30 fm. This is because the nucleons closer to the source of the electromagnetic field interact more with equivalent photons of high frequency than more distant nucleons, which encounter equivalent photons of lower frequency. Since the equivalent photon distribution is strongly nonlinear for this impact-parameter regime, the contributions to $d\sigma/db$ of adjacent nucleons dominate over those of the more distant nucleons by up to about 20%. For larger impact parameters this nonlinearity decreases and, as expected, the impact-parameter dependent cross section converges for the two models for the nucleon number density we studied. At $b = R$ the impact-parameter dependent cross section with the ansatz (11) has its maximum and deviates extensively from the corresponding one with the ansatz (12). Since we are only interested in impact parameters larger than twice the nuclear radius, $b > 2R$, these deviations are irrelevant. – The effect of nuclear shadowing has already been incorporated for the two ansätze (11) and (12) on a very elementary level. No absorption length, which may also be frequency dependent, has been considered. In order to find out how large the modifications for $d\sigma/db$ would be for such a refinement, we have also studied the results for the nucleon number density (13) of a cut sphere; confer again fig. 2. It turns out that the results following from the ansatz (13) do not deviate significantly from those with the ansatz (12) as can be realized in fig. 4. – We conclude, that refined models of the nucleon number

density do not lead to a noticeably different impact-parameter dependent cross section. For the remaining calculations we will rely upon the "cut sphere" ansatz (13).

Throughout this publication we use the Born expression (14) for the elementary photon gluon fusion cross section, which also enters in the structure function (24). QCD corrections to the Born cross section will not be considered here. We are aware that these modifications are important. They are represented by the so-called K-factor, which is of order two. In this respect it would also be of interest to study how much the K-factor depends on the nuclear medium. This will be the subject of further investigations. – Another point, which also needs further investigation, is the lower cutoff $b_{\min} = 2R$ of the integration over the impact parameter appearing in the expression (24) for the structure function. Even when the two colliding nuclei overlap there is still the probability of a diffractive process, which means that still only one photon fuses with one gluon. Furthermore, already for impact parameters larger than twice the nuclear radius two gluons from the tails of the nucleon distributions of the nuclei might fuse to produce a quark pair and thus represent a small background to the photon-gluon fusion process. – Those effects have to be under control in order to deduce the in-medium gluon distribution with sufficiently high precision.

3.2 Deduction of the gluon distribution – part I

For the discussion in this section we assume that we know from which nucleus the gluon originates. We do this because the cross section for the case that one ignores which nucleon contributes a photon respectively a gluon is most easily obtained from the results we shall present here. – In practice we can not distinguish to which nucleus the gluon belongs: After the photon-gluon fusion the two outgoing quarks will dress themselves to be colourless. The net colour, resulting from the ingoing gluon, has to be reabsorbed by the nucleus in form of a soft gluon and will lead to its disintegration. Since the energy scale involved is not so much different compared to electromagnetic disintegration, both processes should be indistinguishable for the case that the other nucleus dissociates electromagnetically, which is mainly nucleon emission after excitation of the giant resonance [13].

The double differential cross section $\frac{d^2\sigma}{dMdy}$ of eq. (9) for photon-gluon fusion into a $b\bar{b}$ -quark pair ($m_b = 4.5$ GeV) in a peripheral Pb+Pb collision at LHC energies ($\gamma = 3000$) is shown in fig. 5; here the Born expression (14) of the elementary photon-gluon fusion cross section (3), a Gaussian nuclear charge formfactor for the equivalent photon distribution (10), the nucleon number distribution (13) of a cut sphere and the parametrisation of the gluon distribution from ref. [12] have been used. A contour line of $10 \mu\text{b}/\text{GeV}$ for the double differential cross section translates with an assumed LHC luminosity of $\mathcal{L} = 2 \cdot 10^{27} \text{ cm}^{-2}\text{sec}^{-1}$, a running time of $T = 10^6 \text{ sec/year}$, a rapidity window of $\Delta y = 0.1$ and an invariant mass resolution of $\Delta M = 1 \text{ GeV}$ into a production rate of $2 \cdot 10^3 \text{ } b\bar{b}\text{-pairs per year}$. – Because we have chosen the gluon to come from the right,

$\frac{d^2\sigma}{dMdy}$ is not symmetric around central rapidity $y = 0$; the main contributions result from small negative rapidities. This is due to the fact that the equivalent photon distribution has a different frequency spectrum than the gluon distribution, so that in average the gluon carries more momentum than the equivalent photon.

Now let us set a window $(\Delta y, \Delta M)$ around the point (y, M) and assume that the double differential cross section $\frac{d^2\sigma}{dMdy}$ has been measured for that window. To avoid a misunderstanding we emphasize: Here $\frac{d^2\sigma}{dMdy}$ has been calculated with a given gluon distribution, just to demonstrate the method how the gluon distribution could be deduced experimentally. Then according to eqs. (9) and (24) the gluon distribution can be identified as the ratio of the double differential cross section and the structure function $\xi(\omega_1, \omega_2)$,

$$g(\omega_2 = \frac{M}{2}e^{-y}) = \frac{\left(\frac{d^2\sigma}{dMdy}\right)_{\text{exp}}}{\frac{M}{2}\xi(\omega_1 = \frac{M}{2}e^y, \omega_2 = \frac{M}{2}e^{-y})} \quad . \quad (25)$$

The structure function $\frac{M}{2}\xi(\omega_1, \omega_2)$, which is determined theoretically, is depicted in fig. 6. Besides the point (y, M) there are many other points, which yield the same constant value for $\omega_2 = \overline{M}/2$; all points along the curve $M = \overline{M}e^y$ correspond to the same value of ω_2 . It follows, that $g(\omega_2 = \overline{M}/2)$ could also be determined from any other window along the curve $M = \overline{M}e^y$. Therefore we can conclude:

$$g(\omega_2 = \overline{M}/2) = \frac{1}{\overline{N}} \sum_{M=\overline{M}e^y} \frac{\left(\frac{d^2\sigma}{dMdy}\right)_{\text{exp}}}{\frac{M}{2}\xi(\omega_1 = \frac{M}{2}e^y, \omega_2 = \frac{M}{2}e^{-y})} \quad , \quad (26)$$

where $\overline{N}$ represents the number of windows chosen along $M = \overline{M}e^y$. This will increase the statistics.

Usually the gluon distribution is parametrized in terms of the dimensionless variable $x = \omega_2/p = \frac{M}{2p}e^{-y}$, so that $g(\omega_2) = \frac{1}{p}g(x = \frac{\omega_2}{p})$. For Pb+Pb collisions at LHC energies the nucleon momentum is equal to $p = 3000$ GeV; hence $x = \frac{M}{6000 \text{ GeV}}e^{-y}$. Then the two points $(y = -3.5, M = 15 \text{ GeV})$ and $(y = 1.5, M = 15 \text{ GeV})$ on the contour line $10 \mu\text{barn/GeV}$ of fig. 5 correspond to x -values of $x = 0.083$ and $x = 0.00056$. The points $(y = -4, M = 20 \text{ GeV})$, $(y = -2, M = 35 \text{ GeV})$, $(y = 0, M = 35 \text{ GeV})$ and $(y = 2, M = 20 \text{ GeV})$, all belonging to the contourline $1 \mu\text{barn/GeV}$ of fig. 5, correspond to x -values of $x = 0.182, 0.043, 0.0058$ and 0.00045 , respectively. This demonstrates that the photon-gluon fusion process into a $b\bar{b}$ quark pair occurring in peripheral ultrarelativistic heavy-ion collisions at LHC represents a very good tool to determine the low x and very low x regime of the gluon distribution in a nuclear medium.

The experimental detector setup might not cover the full phase space. In fact the proposed ALICE detector at LHC will only cover the polar angular region $\pi/4 \leq \Theta \leq 3\pi/4$ [6]. Since it is essential for our analysis to determine the invariant mass and the center of momentum rapidity, both produced heavy quarks have to be detected. Thus we have required condition (17), which has led to the modifications (18) of the total

cross section and (21),(22) of the elementary photon-gluon fusion cross section of lowest order. For the double differential cross section $\frac{d^2\sigma^{\text{cut}}}{dMdy}$ the rapidity y now only runs from $-y_{\text{cut}}^* \leq y \leq y_{\text{cut}}$, where y_{cut}^* follows from eq. (20). For the LHC case the lower boundary will always be $-y_{\text{cut}}^* = -y_{\text{cut}}$, because the invariant mass will stay much smaller than the nucleon momentum, i.e. $M \ll 2p$, so that $\ln(M/2p) \ll -1 \leq -y_{\text{cut}}$. The double differential cross section $\frac{d^2\sigma^{\text{cut}}}{dMdy}$ is exemplified in fig. 7; the same equivalent photon, nucleon number and gluon distribution have been chosen as in fig. 5. Due to the condition (17) the double differential cross section goes to zero for $y = \pm y_{\text{cut}}(M)$. The edges appearing at $y = 0$ are also a result of the same condition. It is noticed that $\frac{d^2\sigma^{\text{cut}}}{dMdy}$ is now more symmetric around $y = 0$, which is in contrast to fig. 5; this can be understood in the following way: A pair of quarks will be produced with a certain center of momentum rapidity y . Since both quarks should be detected, their two single rapidities $y \pm y'$ have to fall into the tiny experimental accessible rapidity interval $|y \pm y'| \leq y_{\text{cut}}$; therefore $y = 0$ plays a special role, so that $\frac{d^2\sigma^{\text{cut}}}{dMdy}$ appears to be almost symmetric. – For the analysis to extract the gluon distribution, the modified structure function

$$\xi^{\text{cut}}(\omega_1, \omega_2) = \int_{2R}^{\infty} d^2b \int d\vec{x}_{\perp} N(\vec{x}_{\perp}) n(\omega_1, \vec{x}_{\perp} - \vec{b}) \sigma_{\gamma g \rightarrow q\bar{q}}^{\text{cut}}(\omega_1, \omega_2) \quad (27)$$

is needed; compare with eq. (24). It is shown in fig. 8. As a result of the rapidity cuts, the values of ξ do not increase for decreasing negative rapidities as was the case in fig. 6; indeed ξ drops to zero as $|y| \rightarrow y_{\text{cut}}(M)$. The edges at $y = 0$ in fig. 8b are due to the edge at $y = 0$ for the cut condition (17). Apart from that, the deduction of the gluon distribution is analogous to eqs. (25) and (26) and will not be repeated here. – The points $(y = -0.42, M = 12 \text{ GeV})$ and $(y = 0.38, M = 12 \text{ GeV})$ on the contour line $10 \mu\text{barn/GeV}$ of fig. 7 correspond to x -values of $x = 0.0030$ and 0.0014 , respectively; the points $(y = -0.7, M = 16 \text{ GeV})$ $(y = 0, M = 28 \text{ GeV})$ and $(y = 0.6, M = 20 \text{ GeV})$ on the contourline $1 \mu\text{barn/GeV}$ correspond to $x = 0.0054, 0.0047$ and 0.0018 , respectively. This clearly demonstrates, that the rapidity cut proposed for the ALICE detector at LHC will only allow to deduce the inmedium gluon distribution for the very low x regime $0.001 \leq x \leq 0.01$, whereas a full rapidity coverage would allow $0.0005 \leq x \leq 0.15$ as we have demonstrated before.

3.3 Deduction of the gluon distribution – part II

In a realistic experiment it will be extremely difficult to identify the nucleus from which the gluon came. The best and only trigger for photon-gluon fusion events would be two jets and nothing else. For this case the strategy sketched in the previous section to deduce the gluon distribution will fail. The reason for this is the following: Let us choose again a window $(\Delta y, \Delta M)$ around the point (y, M) . If the gluon came from the right, we could assign a gluon frequency $\omega_2 = \frac{M}{2} e^{-y} = \frac{\overline{M}}{2}$ to that window. But if the gluon belonged to the nucleus coming from the left, we ought to assign a gluon frequency $\omega_2' = \frac{M}{2} e^y = \frac{\overline{M}'}{2}$

to that window, because then the role of the frequencies ω_1 and ω_2 have interchanged. Therefore the double differential cross section $\frac{d^2\tilde{\sigma}}{dMdy}$ is composed of two parts:

$$\frac{d^2\tilde{\sigma}}{dMdy} = g(\omega_2)\xi(\omega_1, \omega_2) + g(\omega'_2)\xi(\omega'_1, \omega'_2) \quad , \quad (28)$$

where $\omega_1 = \omega'_2 = \frac{M}{2}e^y$ and $\omega_2 = \omega'_1 = \frac{M}{2}e^{-y}$; a tilde has been introduced for this cross section in order to distinguish it from the one before, for which the origin of the gluon was known. In general, the gluon frequencies are different, $\omega_2 \neq \omega'_2$, so that the double differential cross section $\frac{d^2\tilde{\sigma}}{dMdy}$ will no longer be proportional to the gluon distribution.

Only if $\omega_2 = \omega'_2$, the gluon distribution could be directly extracted from eq. (28). This is the case for zero rapidity, i.e. $y = 0$, where $\omega_1 = \omega_2 = \omega'_1 = \omega'_2$. In other words, if $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}$ is measured experimentally and divided by the theoretical structure function $\xi(\omega_1 = \omega_2 = \frac{M}{2})$, the gluon distribution $g(\omega_2 = \frac{M}{2})$ can be deduced, where the gluon frequency ω_2 is equal to the half of the invariant mass M of the produced $b\bar{b}$ quark pair. The structure function $\xi(\omega_1 = \omega_2 = \frac{M}{2})$ is depicted in fig. 6a as the dense dashed curve. Figure 9 displays the double differential cross section $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}$ for various parametrisations of the gluon distribution. The two old parametrisations of Duke and Owens [14] are used for the dense dotted and dashed curve, whereas the new parametrisation of Owens [12] has been used for the solid curve; the resolution momentum has been set equal to $Q^2 = M^2$. For the wide dotted curve the parametrisation of ref. [12] has been used again, but now with a rescaled $Q^2 = 2M^2$; this has been suggested in ref. [15] for the inmedium gluon distribution. For small invariant masses M the different parametrisations for the gluon distribution result in variations for $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}$ by a factor of 1.5; for larger invariant masses the deviations are smaller. For invariant masses in the range $10 \text{ GeV} \leq M \leq 50 \text{ GeV}$, which corresponds to the range $0.0017 \leq x \leq 0.0083$ for the dimensionless variable x used for the gluon momentum, the values of the double differential cross section $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}$ lie in the range of $100 \text{ } \mu\text{barn/GeV}$ down to $0.5 \text{ } \mu\text{barn/GeV}$. With the assumed LHC luminosity $\mathcal{L} = 2 \cdot 10^{27} \text{ cm}^{-2}\text{sec}^{-1}$ and running time $T = 10^6 \text{ sec/year}$ for the Pb+Pb collider mode, a rapidity width of $\Delta y = 0.1$ and a resolution of $\Delta M = 1 \text{ GeV}$ for the invariant mass, these cross sections translate into $2 \cdot 10^4 - 10^2$ produced $b\bar{b}$ quark pairs per year, which should be sufficient to get reasonable statistics.

The statistics might even be increased if cross checks for $(\Delta y, \Delta M)$ windows away from zero rapidity are used in the spirit of eq. (28): Suppose the gluon distribution $g(\omega_2)$ has been determined from the double differential cross section $\frac{d^2\tilde{\sigma}}{dydM}\big|_{y=0}$ in the windows $(\Delta y, \Delta M)$ with $y = 0$. These values could be inserted as initial values in the right hand side of the relation (28) together with the theoretically known structure function $\xi(\omega_1, \omega_2)$ to determine $\frac{d^2\tilde{\sigma}}{dydM}$ for windows $(\Delta y, \Delta M)$ away from $y = 0$; those determined values of $\frac{d^2\tilde{\sigma}}{dydM}$ can be compared to the actual experimentally measured ones in the windows $(\Delta y, \Delta M)$ with y unequal zero. This comparison results in an optimisation for the values of the gluon distribution and in an extension of the accessible x -regime.

In comparison to the double differential cross section $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}$ we now discuss the differential cross section $\frac{d\tilde{\sigma}}{dM}$, which is integrated over the rapidity spectrum. It has already been proposed and used in the references [3] and [4] to deduce the gluon distribution. The apparent disadvantage is that the in-medium gluon distribution cannot be deduced directly; only different parametrisations of the gluon distribution can be tested to agree or disagree with the experimental outcome. Differential cross sections $\frac{d\tilde{\sigma}}{dM}$ for the same parametrisations of the gluon distribution already used for fig. 9 are shown in fig. 10. The variations for different parametrisations are not as large as for $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}$, which is depicted in fig. 9. On the other hand, due to the full rapidity coverage the production rates are increased by a factor of about 50 with respect to the analysis at $y = 0$. Here it remains to clarify how much of this gain will survive once rapidity cuts for the experimental detector are considered; this will be the next topic.

The introduction of the rapidity cuts (17) for the proposed LHC detector has only a small impact on the double differential cross section $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}$ with respect to rapidity and invariant mass. $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}^{\text{cut}}$ with cuts included is shown in fig. 11. Compared to $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}$ without cuts, $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}^{\text{cut}}$ is reduced by a factor of about 1.5 for small invariant masses. For larger invariant masses this factor increases as more and more pairs of quarks lie outside the cut region; for $M = 30$ GeV this factor is already about 2.5. The sensitivity on the numerous gluon distributions used as before is practically identical to the previous case without cuts, so that the deduction of the gluon distribution would proceed with the same quality. On the other side the rapidity cuts have a major impact on the differential cross section $\frac{d\tilde{\sigma}}{dM}$, which is integrated over the whole rapidity spectrum. $\frac{d\tilde{\sigma}}{dM}\big|_{y=0}^{\text{cut}}$ is depicted in fig. 12. We make the following observations: Compared to $\frac{d\tilde{\sigma}}{dM}$ without cuts, $\frac{d\tilde{\sigma}}{dM}\big|_{y=0}^{\text{cut}}$ is reduced by more than about a factor of 10; the tiny accessible rapidity coverage of the LHC detector represents only a small fraction of the whole rapidity spectrum of produced $b\bar{b}$ -quark pairs. Confer also figs. 5 and 7 again. As a consequence the ratio $\frac{d\tilde{\sigma}}{dM}/(\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0} \cdot \Delta y) \approx 50$ with $\Delta y = 0.1$ turns into $\frac{d\tilde{\sigma}}{dM}\big|_{y=0}^{\text{cut}}/(\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}^{\text{cut}} \cdot \Delta y) \approx 6$ now. Another consequence of the cuts seems to be that $\frac{d\tilde{\sigma}}{dM}\big|_{y=0}^{\text{cut}}$ appears to be more sensitive to the different choices for the gluon distribution; in fact the sensitivity almost matches the one for $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}^{\text{cut}}$. This can be explained by the specific cuts themselves, which more or less concentrate on central rapidity $y = 0$. We realize that the double differential cross section $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}^{\text{cut}}$ as well as the differential cross section $\frac{d\tilde{\sigma}}{dM}\big|_{y=0}^{\text{cut}}$ represent efficient quantities to deduce the gluon distribution. With the expected LHC luminosity of $\mathcal{L} = 2 \cdot 10^{27} \text{ cm}^{-2}\text{sec}^{-1}$, an assumed running time of $T = 10^6 \text{ sec/year}$, a rapidity window of $\Delta y = 0.1$ and an invariant mass resolution of $\Delta M = 1 \text{ GeV}$ the cross section $\frac{d^2\tilde{\sigma}}{dMdy}\big|_{y=0}^{\text{cut}}$ translates into 12000 (300) $b\bar{b}$ -quark pair events per year for $M = 10$ (30) GeV, whereas the cross section $\frac{d\tilde{\sigma}}{dM}\big|_{y=0}^{\text{cut}}$ translates into 70000 (2500) events for the same

invariant mass regime, which corresponds to x -values of $0.0017 \leq x \leq 0.005$. Here an efficient procedure would be first to deduce the gluon distribution directly from $\frac{d^2\tilde{\sigma}}{dMdy}\Big|_{y=0}^{\text{cut}}$, then find an adequate parametrisation and finally test this parametrisation in $\frac{d\tilde{\sigma}}{dM}\Big|_{y=0}^{\text{cut}}$.

4 Conclusions

In peripheral ultrarelativistic heavy-ion collisions a photon stemming from the electromagnetic field of one of the two colliding nuclei can penetrate into the other nucleus and interact with one of its hadrons. In particular the photon might fuse with a gluon to produce a pair of heavy quarks. Because of its semielectromagnetic character this process represents a good and clean opportunity to extract information about the in-medium gluon distribution. Adopting the picture of an equivalent photon - partonic gluon interaction for the description of the fusion process, the double differential cross section $\frac{d^2\tilde{\sigma}}{dMdy}\Big|_{y=0}$ with respect to the invariant mass and the rapidity of the produced quark pair at zero rapidity turns out to be directly proportional to the gluon distribution in the nuclear medium; the proportionality factor is governed by the equivalent photon spectrum, the nuclear density distribution and the elementary photon-gluon fusion cross section, which are all quantities that can be controlled theoretically. Thus, if $\frac{d^2\tilde{\sigma}}{dMdy}\Big|_{y=0}$ is measured somehow for $c\bar{c}$ or $b\bar{b}$ production and divided by this proportionality factor, we directly get the gluon distribution. For $b\bar{b}$ production in peripheral Pb+Pb collisions at LHC energies the cross sections lie in the range $100\text{-}0.5\ \mu\text{barn}$ depending on the invariant mass of the quark pair $10\ \text{GeV} \leq M \leq 50\ \text{GeV}$; this mass range corresponds to the range $0.0017 \leq x \leq 0.0083$ for the gluonic Bjorken variable used to parametrize the gluon distribution. With an assumed luminosity of $\mathcal{L} = 2 \cdot 10^{27}\ \text{cm}^{-2}\text{sec}^{-1}$, a running time of $T = 10^6\ \text{sec/year}$, a rapidity window of $\Delta y = 0.1$ and an invariant mass resolution of $\Delta M = 1\ \text{GeV}$ these cross sections translate into rates of $100\text{ - }20000\ b\bar{b}$ -pairs per year. The statistics might even be increased if cross checks of $\frac{d^2\tilde{\sigma}}{dMdy}$ away from zero rapidity are considered; be aware that $\frac{d^2\tilde{\sigma}}{dMdy}\Big|_{y \neq 0}$ is not proportional to the gluon distribution anymore. In this respect the x -range could also be extended to $0.0005\text{ - }0.15$. This of course requires a large rapidity coverage for the detector. As for the proposed ALICE detector at LHC [6] this is not the case: the polar angle is restricted to $\pi/4 \leq \Theta \leq 3\pi/4$. Since it is essential for our analysis to determine the invariant mass of the produced heavy quark pair, both quark jets have to be detected within these experimental cuts. Taking these cuts into consideration the $\frac{d^2\tilde{\sigma}}{dMdy}\Big|_{y=0}$ analysis is only affected to a small degree; on the other hand extensions away from central rapidity will not be possible anymore, so that one is narrowed to the $0.002 \leq x \leq 0.008$ regime for the deduction of the in-medium gluon distribution. The statistics might be increased for the analysis with rapidity cuts once the differential cross section $d\tilde{\sigma}/dM\Big|_{y=0}^{\text{cut}}$ is considered; although it is not directly proportional to the gluon distribution various parametrisations of the latter can be tested

sensitively.

This study cannot be more than an important intermediate step to achieve the wanted goal, i.e. the deduction of the in-medium gluon distribution. More refined studies are necessary. First of all, it is obvious to repeat the analysis presented here for $c\bar{c}$ -production at RHIC and LHC energies with and without rapidity cuts; the increased rates will allow further extensions for the accessible x -regime of the in-medium gluon distribution. Then, elementary photon-gluon fusion cross sections with QCD corrections should be considered; the deviations to the Born expressions are reflected in the K-factor, which is about two. This factor has to be controlled theoretically in order to determine the gluon distribution precisely. Referring to the used integration of the impact parameter from twice the nuclear radius to infinity, we should check contributions from diffractive processes and from gluon-gluon fusion, which results from the overlapping tails of the nucleon distributions at close impact parameters and which represents a small background. Another point to be clarified is, how much are the produced quark jets deflected, partially or totally absorbed by the remaining nucleus. Here parton cascade models [16] could provide an answer and might set a lower bound on the invariant mass for the analysis. Finally we want to mention, that also lighter heavy ions are of interest. For the system Ca+Ca, for example, we would lose a factor of $(Z^2 A_{\text{eff}})_{\text{Ca}} / (Z^2 A_{\text{eff}})_{\text{Pb}} \approx 1/60$ in the cross section compared to the system Pb+Pb, but according to present predictions [17] we would gain more than a factor of 1000 in luminosity. Thus, the in-medium gluon distribution could be studied in dependence of the nuclear mass number. Also let us remind that whatever the result will be photon-gluon fusion offers a new possibility to study the formation of heavy quarks in a nuclear medium, a process of great importance for the understanding of high energy heavy ion collisions.

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Figures

Figure 1: $\vec{b}$: impact parameter between the two nuclei; $\vec{x}_\perp$: distance of one nucleon with respect to the center of the nucleus; $\vec{x}'_\perp$: distance of one gluon with respect to the center of the corresponding nucleon.

Figure 2: The nucleon number density $N_{cs}(\vec{x}_\perp)$ of the cut sphere is proportional to the projection of the shaded volume onto the transverse plane.

Figure 3: Differential cross section $d\sigma/db$ of eq. (8) (solid line) for b-quark pair production compared to $d\sigma/db$ of eq. (1) with $r = 0.8$ fm (dotted line) and $r = 1.2$ fm (dashed line); for more details see text.

Figure 4: The influence of the various nucleon number densities (11)-(13) on the impact-parameter dependent cross section $d\sigma/db$ for b-quark production via photon-gluon fusion in a Pb+Pb collision with $\gamma = 3000$.

Figure 5: Double differential cross section $\frac{d^2\sigma}{dMdy}$ for b-quark pair production, where the direction of the ingoing gluon has been fixed artificially; for more details see text.

Figure 6: Structure function (24) entering the cross section (9) shown for constant rapidities (a) and for constant invariant masses (b); for more details see text.

Figure 7: Double differential cross section $\frac{d^2\sigma}{dMdy}\Big|_{y=0}^{\text{cut}}$ for b-quark pair production, where the direction of the ingoing gluon has been fixed artificially and where the LHC rapidity cut has been considered; for more details see text.

Figure 8: Structure function (27) entering the cross section (18) with LHC rapidity cuts (17) shown for constant rapidities (a) and for constant invariant masses (b); for more details see text.

Figure 9: Double differential cross section $\frac{d^2\tilde{\sigma}}{dMdy}\Big|_{y=0}$ for b-quark pair production, where the direction of the ingoing gluon is not known, shown for various parametrisations of the gluon distribution; for more details see text.

Figure 10: Differential cross section $d\tilde{\sigma}/dM$ for b-quark pair production shown for various parametrisations of the gluon distribution; for more details see text.

Figure 11: Same as fig. 9 with LHC rapidity cut (17) now included.

Figure 12: Same as fig. 10 with LHC rapidity cut (17) now included.

Figure 1

Figure 2

Figure 3

Figure 4

Figure 5

Figure 6a

Figure 6b

Figure 7

Figure 8a

Figure 9

Figure 10

Figure 11

Figure 12