Transverse Momentum Distribution of Vector Mesons Produced in Ultraperipheral Relativistic Heavy Ion Collisions

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We study the transverse momentum distribution of vector mesons produced in ultraperipheral relativistic heavy ion collisions (UPCs). In UPCs there is no strong interaction between the nuclei, and the vector mesons are produced in photon-nucleus collisions where the (quasireal) photon is emitted from the other nucleus. Exchanging the role of both ions leads to interference effects. A detailed study of the transverse momentum distribution, which is determined by the transverse momentum of the emitted photon, the production process on the target, and the interference effect, is done. We study the unrestricted cross section and the one with an additional electromagnetic excitation of one or both ions; in the latter case small impact parameters are emphasized.

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Because of the strong electromagnetic fields surrounding the heavy ions in relativistic collisions, the Relativistic Heavy Ion Collider (RHIC) and the Large Hadron Collider (LHC) can be seen as a factory of quasireal photons of high energies. One of the interesting photonuclear processes studied in these "ultraperipheral collisions" (UPC) is the coherent production of vector mesons, in particular ρ^0, which has been measured recently at RHIC [1,2]. The coherent production was identified through the transverse momentum distribution of the meson, which is enhanced for values $v_\perp \lesssim 1/R$ where R denotes the nuclear radius. We give a careful theoretical study of the process

$$A + A \rightarrow A^{(*)} + A^{(*)} + V \qquad (1)$$

with ("A^*") and without ("A") an electromagnetic excitation of either one or both ions, predominantly to the giant dipole resonance. This is of interest for the analysis of the RHIC experiments, as well as for future experiments at LHC where also heavier vector mesons like J/ψ and even Υ can be studied in order to explore, e.g., nuclear shadowing effects [3–5]. While the theory of UPC is generally in a good shape [5–9], the specific question of the *transverse momentum distribution* has been paid attention to only in less rigor [7,10,11].

Heavy ion scattering offers a unique possibility to study an important interference effect [10]. As is shown below, the transverse momentum distribution is very sensitive to this effect. It is the purpose of this Letter to give a careful study of this transverse momentum distribution.

The kinematics of the process given in Eq. (1) is denoted by (see Fig. 1)

$$p + k \rightarrow p' + k' + v. \qquad (2)$$

Because of the additional elastic photon exchanges which are schematically denoted by the open blobs in Fig. 1, the momenta Q and Δ are not related to the asymptotic momenta by $Q = p - p'$ and $\Delta = k - k'$ as, e.g., in the pp case (without rescattering) [12]. For small transverse momenta the longitudinal components of the photon momentum and the momentum transfer from the vector meson production ("Pomeron momentum") are given in the c.m. system by the mass m_V and the rapidity Y of the produced meson as

$$Q_0 = \beta Q_z = \frac{m_V}{2} e^Y, \qquad \Delta_0 = -\beta \Delta_z = \frac{m_V}{2} e^{-Y}. \qquad (3)$$

The momenta p and k (see Fig. 1(a)) are given by $p = m_A u_+$ (ion 1) and $k = m_A u_-$ (ion 2), where m_A is the ion mass and $u_\pm = \gamma(1, 0, 0, \pm\beta)$. In the exchange process (see Fig. 1(b)) the photon is emitted from ion 2, and the "Pomeron" from ion 1.

For the application of AuAu and PbPb collisions in this Letter the Coulomb parameter $\eta = \frac{Z_1 Z_2 e^2}{\hbar v} \approx Z_1 Z_2 \alpha$ is

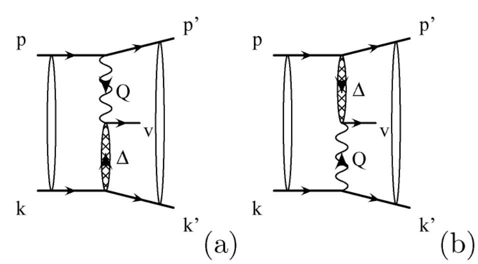

FIG. 1. A schematic Feynman diagram for the vector meson production in ultraperipheral heavy ion collisions (a). The corresponding exchange diagram is also shown (b).

much larger than 1, and we can use the semiclassical approximation [2,13,14]. We also show how this can be derived from eikonal or Glauber theory, more details will be given in a forthcoming publication [15]. Using a simple model for the meson production process, we are able to give analytical results. Implications for the current experiments at RHIC and for future experiments at the LHC [4,10] are discussed. In the semiclassical approximation the two ions move along a straight line, and the process is described by an impact parameter dependent amplitude $a(\vec{b}, \vec{v}_\perp, Y)$. In contrast to the momentum of the vector meson the momenta of the outgoing ions are not detected and the differential cross section is given by

$$\frac{d^3\sigma}{d^2v_\perp dY} = \frac{1}{2(2\pi)^3}\sum_{e_V}\int d^2b|a_{fi}(\vec{b}, \vec{v}_\perp, Y)|^2. \quad (4)$$

The integration over the impact parameter $\vec{b}$ corresponds to an integration over the unobserved momenta k' and p' of the scattered ions. The different processes that can occur according to Eq. (1) factorize [14]:

$$a_{fi}(\vec{b}, \vec{v}_\perp, Y) = a_{\rm nucl}(\vec{b})a_1(\vec{b})a_2(\vec{b})a_V(\vec{b}, \vec{v}_\perp, Y). \quad (5)$$

This factorization is clearly valid for the processes considered here; see [14] for details. The strong absorption due to the interaction of the ions for $b < 2R$ is given by $a_{\rm nucl}(\vec{b}) \approx \Theta(b - 2R)$ with the nuclear radius R. We use $R = 7$ fm in our calculation. While the nucleon radius of interaction increases significantly with energy, this increase is relatively small compared to the radius of a (heavy) nucleus. $a_V(\vec{b}, \vec{v}_\perp, Y)$ describes the vector meson production. Additional electromagnetic excitation amplitudes of ion 1 and/or 2 are denoted by $a_1(\vec{b})$ and $a_2(\vec{b})$. The background contribution of vector meson production by inelastic photon emission was shown to be small; see Sec. 3.2 of [8] or [14].

The cross section can be written as

$$\frac{d^3\sigma}{d^2v dY} = \frac{1}{2(2\pi)^3}\int_{2R}^{\infty} d^2b f_{ij}(b)\sum_{e_V}|a_V(\vec{b}, \vec{v}, Y)|^2. \quad (6)$$

where $f_{ij}(b)$ takes the triggering condition of the measurement into account. Without any condition imposed we have $f_{00}(b) = 1$. For vector meson production together with the electromagnetic excitation of one of the ions one has either $f_{10}(b) = P_1(b) = |a_1(b)|^2$ or $f_{01}(b) = P_2(b) = |a_2(b)|^2$, for the mutual excitation of both ions $f_{11}(b) = P_1(b)P_2(b)$.

The electromagnetic excitation probabilities for the range of b of interest here are given by $P_i(b) = \frac{S}{b^2}$ with $S \approx 5.45 \times 10^{-5} Z^3 N A^{-2/3}$ fm^2 [16].

In the semiclassical treatment the amplitude of an electromagnetic process can be written as [17,18]

$$a(\vec{b}) = \int \frac{d^4Q}{(2\pi)^4} A^{\mu}_{\rm ext}(\vec{b}, Q)J_\mu(Q), \quad (7)$$

where

$$A^{\mu}_{\rm ext}(b, Q) = 2\pi Ze Q^{\mu}_{\perp}\delta(Qu_+)\frac{\gamma}{Q_0}\frac{F(Q^2)}{Q^2}\exp(-i\vec{Q}_\perp\vec{b}) \quad (8)$$

is the Liènard-Wiechert potential.

For the elastic form factor $F(Q^2)$ we choose $F(Q^2) = \exp(Q^2R_\gamma^2)$, with $R_\gamma \approx \sqrt{\langle r^2\rangle/6} \approx 2.2$ fm. Alternatively, we can set $F(Q^2) = 1$, i.e., $R_\gamma = 0$, as the electric field outside a spherically symmetric charge distribution is the same as that of a corresponding point charge. Indeed, the whole effect of the nuclear distribution is encoded in one parameter R [8]. We find numerically the effect of the form factor to be rather small, justifying this assumption. We still keep it for completeness in the following equations.

In order to describe the meson production we need an expression for the electromagnetic current $J(A \to A + V)$. In the following we choose it as

$$J^\mu(Q) = e^\mu_V F_0(Y)\exp(-\Delta_\perp^2 R_V^2)\delta(\Delta u_-)(vu_-) \quad (9)$$

with $\Delta_\perp = v_\perp - Q_\perp$. e_V is the polarization of the outgoing vector meson, which by assuming s-channel helicity conservation is identical to the one of the incoming photon; see [19] for details. R_V is a parameter increasing with the photon energy, as well as depending on the target nucleus and the meson produced. We have chosen $R_V \approx 2.2$ fm throughout this Letter, which reproduces the slope of the transverse distribution for Pb and Au. This form agrees also with the one proposed in [20] and the Gaussian form agrees well with the results in Figs. 9 and 11 of [21]. It has been mainly chosen for ease of calculations, whereas the formalism given can be extended to more realistic forms, e.g., based on a full eikonal description. A simple extension is possible by using a sum of Gaussians for J_μ. This current can be related with the help of the vector dominance model, valid for the ρ meson, to the elastic scattering amplitude $V + A \to V + A$; see, e.g., [5,11,21].

In general, F_0 will be complex. For ρ production and the energies of RHIC and LHC the real part is small (of the order of 10%), whereas it can be substantial for J/ψ and Υ (up to 40%). For $Y \neq 0$ (and also for asymmetric collisions) there is a sensitivity to the phase of F_0. But for $Y = 0$ the result depends only on the absolute value $|F_0|$. We choose the value of $|F_0|$ to reproduce the cross sections given in [3,22] as $d\sigma/dY_\rho(\text{RHIC}) = 70$ mb and $d\sigma/dY_{J/\psi}(\text{LHC}) = 0.75$ mb.

Using these expressions, we get for the amplitude

$$a_V(\vec{b}, \vec{v}_\perp, Y) = \frac{ZeF_0}{(2\pi)^3}\exp(-Q_l^2R_\gamma^2)\int d^2Q_\perp(\vec{Q}_\perp\vec{e}_V) \times\frac{\exp(-Q_\perp^2R_\gamma^2)}{Q_\perp^2 + Q_l^2}\exp(-i\vec{Q}_\perp\vec{b}) \times\exp(-R_V^2(\vec{v}_\perp - \vec{Q}_\perp)^2) \quad (10)$$

with $Q_l^2 = Q_z^2 - Q_0^2 = (\frac{Q_0}{\beta\gamma})^2$. If we would neglect $\vec{Q}_\perp$ in $\exp(-R_V^2(\vec{v}_\perp - \vec{Q}_\perp)^2)$, the dependence of $|a_V|^2$ and $d^3\sigma/d^2 v_\perp dY$ on $v_\perp$ would be of the form $\exp(-2R_V^2 v_\perp^2)$, which is due to J alone. This coincides with the result for an incident photon of zero transverse momentum. The effect of the finite $Q_\perp$ distribution of the photon is to broaden this distribution. As the width of the $Q_\perp$ distribution depends on b via $\exp(-i\vec{Q}_\perp \vec{b})$, the effect of this broadening will depend on b. This effect is largest for small b, as the perpendicular momentum distribution of the photon is of the order $1/b$.

An analytic approximation for a_V can be found in the region of small b, that is, if $b < 1/Q_l$, which corresponds to the sudden limit. One gets (see [15] for details)

$$a_V(\vec{b}, \vec{v}_\perp, Y) \approx \frac{ZeF_0 i}{(2\pi)^2} \frac{(\vec{b} + 2i\vec{v}_\perp R_V^2)\vec{e}_V}{(\vec{b} + 2i\vec{v}_\perp R_V^2)^2} \exp(-Q_l^2 R_\gamma^2) \times \exp(-v_\perp^2 R_V^2) \times \left[\exp\left(\frac{-(\vec{b} + 2i\vec{v}_\perp R_V^2)^2}{4(R_V^2 + R_\gamma^2)}\right) - 1\right]. \quad (11)$$

The same final state can be obtained by exchanging the roles of both ions; see Figs. 1(a) and 1(b). The corresponding amplitude $a_V^X(\vec{b}, \vec{v}_\perp, Y)$ is given by

$$a_V^X(\vec{b}, \vec{v}_\perp, Y) = \int \frac{d^4Q}{(2\pi)^4} A_{\rm ext}^\mu(0, Q) J_\mu^X(Q), \quad (12)$$

where the impact parameter for $A_{\rm ext}$ is now $\vec{b} = 0$ and u_+ is replaced by u_-. The electromagnetic current J^X is now for vector meson production on an ion at position $\vec{b}$. One finds

$$J_\mu^X(Q) = J_\mu(Q)\exp(-i\vec{Q}_\perp \vec{b}) = J_\mu(Q)\exp(-i\vec{v}_\perp \vec{b} + i\vec{\Delta}_\perp \vec{b}). \quad (13)$$

We find that the exchange amplitude is of the form

$$a_V^X(\vec{b}, \vec{v}_\perp, Y) = a_V(-\vec{b}, \vec{v}_\perp, -Y)\exp(-i\vec{v}_\perp \vec{b}). \quad (14)$$

This has a simple interpretation: Y is replaced by $-Y$, the direction of $\vec{b}$ needs to be reversed, and, in addition, the origin needs to be shifted by $\vec{b}$, leading to the extra phase $\exp(-i\vec{v}_\perp \vec{b})$. This relation was also used in [10,23]. With a_V from Eq. (10) we finally get

$$a_V^{\rm tot}(\vec{b}, \vec{v}_\perp, Y) = a_V(\vec{b}, \vec{v}_\perp, Y) + e^{-i\vec{v}_\perp \vec{b}} a_V(-\vec{b}, \vec{v}_\perp, -Y). \quad (15)$$

The analytic expression in Eq. (11) allows us to discuss some properties of the transverse momentum distribution of the process: In the limit $v_\perp R_V^2 \ll b$ one has $a_V \sim \vec{b}\vec{e}_V$ and $a_V^X = -a_V$; i.e., the amplitudes have a relative sign of -1, leading to destructive interference at small b. In the other limit $v_\perp R_V^2 \gg b$ one has $a_V^X = a_V$, i.e., the same relative sign, but a_V and a_V^X are smaller than in the first case due to the last exponential in Eq. (11). The transverse momentum distribution is therefore more complex than treated in [10,23].

We can also derive the results starting from the eikonal or Glauber approach to multiphoton processes in UPC collisions; see [14]. In this case the scattering amplitude is given by

$$f_{fi,\rm Glauber}(\vec{K}) = \frac{i\pi}{k} \int d^2b \exp(i\vec{K}\vec{b})\langle f| \exp(i\chi(\vec{b}))|i\rangle, \quad (16)$$

where $\vec{K} = \vec{k}' - \vec{k}$ is the total momentum transfer to the "target" nucleus. The eikonal $\chi(b)$ takes care of all the different elastic and inelastic processes. In our case we have

$$\chi(\vec{b}) = \chi_{\rm nuc}(b) + \chi_C(b) + \chi_1(b) + \chi_2(b) + \chi_V(\vec{b}). \quad (17)$$

The term $\chi_{\rm nuc}(b)$ describes the effect coming from the nuclear interaction between the two ions. It can be approximated by $\exp(i\chi_{\rm nuc}(b)) \approx \Theta(b - 2R)$. The term $\chi_C \approx 2\eta \log(kb)$ describes the elastic Coulomb scattering. The last three terms describe the additional electromagnetic interactions: the possible excitation of the first and second nucleus and the vector meson production. The eikonal phase for the vector meson production process $\chi_V(b)$ can usually be treated in lowest order by expanding the exponential. The second order term would describe double ρ production, which is still sizable; see [11]. Bracketing with the initial and final states we get

$$\langle f|\exp(i\chi(\vec{b}))|i\rangle \approx i\exp(i\chi_{\rm nuc}(b))\exp(+i\chi_C(b)) \times \langle f_1|\exp(i\chi_1(b))|i_1\rangle \times \langle f_2|\exp(i\chi_2(b))|i_2\rangle\langle V, i_2|\chi_V(\vec{b})|i_2\rangle \quad (18)$$

with $|i\rangle = |i_1, i_2\rangle$ the initial (ground) states of the two ions and $|f\rangle = |f_1, f_2\rangle|V\rangle$ the final states of the ions and the meson. In order for this process to factorize, we made the reasonable assumption that the vector meson production on the excited nucleus is the same as the one on the nucleus in the ground state: $\langle V, i_2|\chi_V(b)|i_2\rangle \approx \langle V, f_2|\chi_V(b)|f_2\rangle$.

There is a correspondence of these terms to the different semiclassical amplitudes $a_{fi}(b)$ in Eq. (5), which was also explored in [24]. The χ_V is given by

$$\chi_V(\vec{b}) = -1 \int \frac{d^4Q}{(2\pi)^4} A_{\rm eik}^\mu(\vec{b}, z, Q)\hat{J}_\mu(Q) \quad (19)$$

with $A_{\rm eik}^\mu$ as given in Eq. (8) with $\delta(Qu_+)$ replaced with $\delta(\gamma(Q_0 - Q_z))$, which corresponds to the expression of the semiclassical amplitude a_V in Eq. (7) in the sudden limit.

The major difference between the two approaches is the presence of the Coulomb eikonal $\chi_C(b)$. For $\eta \gg 1$, $\chi_C(b)$ is a rapidly varying function, and one can evaluate the

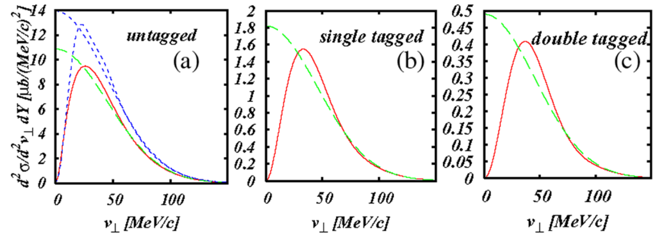

FIG. 2 (color online). The differential cross section $d\sigma/d^2v_\perp dY$ is shown for ρ^0 production at $Y = 0$ at RHIC (Au-Au collisions with $\gamma = 108$). The solid line is the result including the interference, and the dashed line the result from an incoherent adding of the two processes. The results for the three different tagging cases are given in (a), (b), and (c). The approximate result is shown as a dotted line only in the first (untagged) case. In the other two cases it cannot be distinguished from the full calculation.

Glauber expression by means of the well known saddle point approximation. As the momentum transfer to the ion is not measured in our case, one calculates "inclusive" cross sections by integrating over K. This gives the same result as in the semiclassical case of Eq. (5) in the sudden limit ($Q_l = 0$), as the Coulomb phase is purely imaginary.

We use both the exact expression Eqs. (6) and (10), as well as the approximate analytical result Eq. (11), to calculate results for the case studied in [10]. They are shown in Fig. 2. The analytic result is too large in the untagged case, but its shape agrees quite nicely with the full calculation. The effect of tagging for small b is a shift of the maximum of the curve to larger values of $v_\perp$ and a more pronounced interference structure. In Fig. 3 we also show the similar results for J/ψ production at the LHC.

Let us summarize our findings: We have put the transverse momentum distribution on a firm theoretical basis starting our derivation from the semiclassical approximation or alternatively from Glauber theory. The meson transverse momentum distribution was derived as a function of b and an analytic expression was given. The interference phenomenon was derived within this model. As the main outcome we find in this Letter that for a good understanding of the interference phenomenon a careful study of the transverse momentum distribution is essential. Whereas formally the results look similar to the one given in [2,10], differences appear both in the transverse momentum distribution as a function of b and in the form of the interference. This leads to a more complex result in the intermediate $v_\perp$ region. Our findings are important in analyzing the experimental data of STAR and also PHENIX. Results have also been given for future LHC measurements, which would be even more interesting for J/ψ or even Υ production.

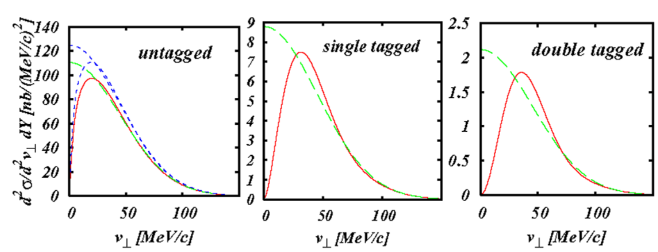

FIG. 3 (color online). The differential cross section $d\sigma/d^2v_\perp dY$ is shown for J/ψ production at $Y = 0$ at LHC (Pb-Pb ions with $\gamma = 3000$). The lines are the same as in Fig. 2.

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