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Nuclear geometry of jet quenching

I.P. Lokhtin, A.M. Snigirev

M.V. Lomonosov Moscow State University, D.V. Skobeltsyn Institute of Nuclear Physics, Moscow, Russia

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Abstract. The most suitable way to study jet quenching as a function of the distance traversed is varying the impact parameter b of the ultrarelativistic nucleus–nucleus collision (the initial energy density in the nuclear overlapping zone is almost independent of b up to $b \sim R_A$). It is shown that the b-dependences of the medium-induced radiative and collisional energy losses of a hard parton jet propagating through dense QCD matter are very different. The experimental verification of this phenomenon could be performed for a jet with non-zero cone size based on the essential difference between the angular distributions of the collisional and radiative energy losses.

1 Introduction

The experimental investigation of ultrarelativistic nuclear collisions offers a unique possibility of studying the properties of strongly interacting matter at a high energy density when the hadronic matter is expected to become deconfined and a gas of asymptotically free quarks and gluons is formed. This is called a quark–gluon plasma (QGP); the colour interactions between the partons are screened owing to collective effects (see, for example, the reviews in [1–4]).

In recent years, a great deal of attention has been paid to the study of "hard" probes of QGP – heavy quarkonia and hard partonic jets, which do not appear as constituents of the thermalised system, but can carry information about the earliest stages of its evolution. In particular, the strong suppression of the yield of heavy quark vector mesons as J/Ψ, Ψ' ($c\bar{c}$ states) and Υ, Υ', Υ'' ($b\bar{b}$ states) is one of the promising signatures of the quark–gluon plasma formation in heavy ion collisions [5]. An intriguing phenomenon is the "anomalously" small yield of Ψ-resonances, observed in Pb–Pb collisions in the NA50 experiment (CERN-SPS) [6], this being inconsistent with the conventional model of pre-resonance absorption in cold nuclear matter. Although the interpretation of this phenomenon as a result of the formation of a QGP is quite plausible [7], alternative explanations have also been put forward, for example, Ψ–h rescattering on comoving hadrons [8]. Thus, the nature of this "anomalous" suppression of Ψ-resonance production is not yet fully understood, and it should be completely explained in future [9]. For heavier ($b\bar{b}$) systems, a similar suppression effect in super-dense strongly interacting matter is expected at higher temperatures than for $c\bar{c}$, which are expected to be reached in central collisions of heavy ions at the RHIC at BNL and LHC at CERN colliders.

Along with the suppression of heavy quarkonia, one of the processes which may give information about the earliest stages of evolution of the dense matter formed in ultrarelativistic nuclear collisions is the passage through the matter of hard jets of colour-charged partons, pairs of which are created at the very beginning of the collision process (typically, at $\lesssim 0.01\,\mathrm{fm}/c$) as a result of individual initial hard nucleon–nucleon (parton–parton) scatterings. Such jets pass through the dense parton matter formed due to mini-jet production at larger time scales ($\sim 0.1\,\mathrm{fm}/c$), and interact strongly with the comoving constituents in the medium, changing its original properties as a result of additional rescatterings. The inclusive cross section for hard jet production processes is still very small for performing a systematic analysis at the SPS energies ($s^{1/2} \simeq 20\,\mathrm{GeV}$ per nucleon pair), but it increases fast with the energy of the collided nuclei. Thus these will play an important role in the formation of the initial state at the energies of RHIC ($s^{1/2} = 200\,\mathrm{GeV}$ per nucleon pair) and LHC ($s^{1/2} = 5.5\,\mathrm{TeV}$ per nucleon pair) colliders.

The actual problem is to study the energy losses of a hard jet evolving through dense matter. We know two possible mechanisms of energy losses:

(1) radiative losses due to gluon "bremsstrahlung" induced by multiple scattering [10–15], and
(2) collisional losses due to the final state interactions (elastic rescatterings) of high-p_T partons off the medium constituents [16–18]. Since the jet rescattering intensity strongly increases with temperature, the formation of a super-dense and hot partonic matter in heavy ion collisions (with initial temperature up to $T_0 \sim 1\,\mathrm{GeV}$ at LHC [19]) should result in significantly larger jet energy losses as compared with the case of "cold" nuclear matter or hadronic gas at $T \lesssim 0.2\,\mathrm{GeV}$.

Although the radiative energy losses of a high energy parton have been shown to dominate over the collisional losses by up to an order of magnitude [11], a direct experimental verification of this phenomenon remains an open problem. Indeed, with increasing the hard parton energy the maximum of the angular distribution of the bremsstrahlung gluons has a shift towards the parent parton direction. This means that measuring the jet energy as a sum of the energies of the final hadrons moving inside an angular cone with a given finite size θ_0 will allow the bulk of the gluon radiation to belong to the jet and thus makes it possible for the major fraction of the initial parton energy to be reconstructed. Therefore, the medium-induced radiation will, in the first place, soften the particle energy distributions inside the jet and increase the multiplicity of the secondary particles, but will not affect the total jet energy. It was recently shown [12,13] that the radiation of energetic gluons in a QCD medium is essentially different from the Bethe–Heitler independent radiation pattern. Such gluons have formation times exceeding the mean free path for QCD parton scattering in the medium. In these circumstances the coherent effects play a crucial role leading to a strong suppression of the medium-induced gluon radiation. This coherent suppression is a QCD analogue of the Landau–Pomeranchuk–Migdal effect in QED. It is important to notice that the coherent LPM radiation induces a strong dependence of the jet energy on the jet cone size θ_0 [20,21].

On the other hand, the collisional energy losses represent an incoherent sum over all rescatterings. It is almost independent of the initial parton energy. Meanwhile, the angular distribution of the collisional energy loss is essentially different from that of the radiative one. The bulk of "thermal" particles knocked out of the dense matter by elastic scatterings fly away in an almost transverse direction relative to the hard jet axis. As a result, the collisional energy loss turns out to be practically independent on θ_0 and emerges outside the narrow jet cone. Thus the relative contribution of collisional losses would likely become significant for jets with finite cone size propagating through the QGP [20].

In the search for experimental evidence in favour of the medium-induced energy losses a significant dijet quenching (a suppression of high-p_{T} jet pair yield) [22] and a monojet-to-dijet ratio enhancement [23] were proposed as possible signals of dense matter formation in ultrarelativistic collisions of nuclei. Other possible signatures that could directly measure the energy losses involve tagging the hard jet opposite a particle that does not interact strongly as a Z-boson [24] (mostly $q+g \to q+Z(\to \mu^+\mu^-)$, but also $q+\bar{q} \to q+Z$) or a photon [25] (mostly $q+g \to q+\gamma$, but also $q+\bar{q} \to q+\gamma$). The jet energy losses in dense matter should result in a non-symmetric shape of the distribution of differences in P_{T} between the Z-boson (γ) and jet. The above phenomena can be studied in heavy ion collisions [26] with the Compact Muon Solenoid (CMS), which is the general purpose detector designed to run at the LHC [27]. Note that using the $\gamma+\text{jet}$ channel in this case is complicated due to the large background from jet + jet production when one of the particles in the jet in an event is misidentified as a photon (the leading π^0). However, the shape of the distribution of differences in E_{T} between the γ and jet is very different for signal and background, and is still sensitive to the jet energy losses [26].

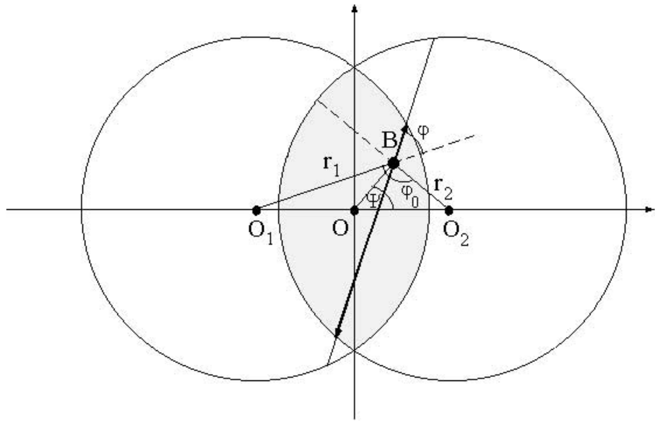

Fig. 1. Jet production in a high energy symmetric nucleus–nucleus collision in the plane of the impact parameter $\mathbf{b}$. O_1 and O_2 are the nucleus centers, $OO_2 = -O_1O = b/2$. $B(r\cos\psi, r\sin\psi)$ is the jet (dijet) production vertex, r is the distance from the nuclear collision axis to B, r_1, r_2 are the distances between the nucleus centers (O_1, O_2) and B; φ is the jet azimuthal angle, and φ_0 is the azimuthal angle between the vectors $\mathbf{r}_1$ and $\mathbf{r}_2$

The advantage of $\gamma+\text{jet}$ and $Z(\to \mu^+\mu^-)+\text{jet}$ channels is that one can determine the average initial transverse momentum of the hard jet, $\langle P_{\mathrm{T}}^{\mathrm{jet}} \rangle \approx \langle P_{\mathrm{T}}^{\gamma,Z} \rangle$. It gives the attractive opportunity to search for coherent effects in a QCD medium: the dependence of the energy losses of the distance traversed can be studied experimentally in different bins of the impact parameter distribution of the nucleus–nucleus collision, or by varying collided ions and selecting the most central collisions. The intriguing prediction associated with the coherence pattern of the medium-induced radiation is that radiative energy losses per unit distance $\mathrm{d}E/\mathrm{d}x$ depend on the total distance traversed, L [12,13]. The value $\mathrm{d}E/\mathrm{d}x$ approaches a behaviour of being proportional to L for a static medium [12], and it has a weaker L-dependence for the case of an expanding medium [13]. The main goal of the present paper is to analyse the possibility of observing the L-dependence of jet energy losses $\mathrm{d}E/\mathrm{d}x$ for a realistic nuclear geometry. In particular, we study the impact parameter dependence of the collisional and radiative jet energy losses in dense QCD matter, created in ultrarelativistic heavy ion collisions.

2 The geometrical model for jet production in nuclear collisions

Let us to consider the simple geometrical model of jet production and jet passing through a dense matter in a high energy symmetric nucleus–nucleus collision. Figure 1 shows the essence of the problem in the plane of the impact parameter $\mathbf{b}$ of two colliding nuclei A–A. The impact parameter b here is the transverse distance between the nucleus centers O_1 and O_2, $OO_2 = -O_1O = b/2$. Let $B(r\cos\psi, r\sin\psi)$ be denoted as a jet (dijet) production vertex, r being the distance from the nuclear collision axis to the B. Then the distance between the nucleus centers (O_1, O_2) and the vertex B can be found to be

$$r_{1,2} = \sqrt{r^2 + \frac{b^2}{4} \pm rb\cos\psi}. \qquad (1)$$

The distribution over the jet production vertex $B(r,\psi)$ at a given impact parameter b is written as

$$P_{AA}(\mathbf{r}, b) = \frac{T_A(r_1)\cdot T_A(r_2)}{T_{AA}(b)}, \qquad (2)$$

where

$$T_{AA}(b) = \int \mathrm{d}^2\mathbf{s} T_A(\mathbf{s}) T_A(\mathbf{b}-\mathbf{s}) = \int_0^{2\pi} \mathrm{d}\psi \int_0^{r_{\max}} r\mathrm{d}r T_A(r_1) T_A(r_2) \qquad (3)$$

is the nuclear overlap function, $T_A(\mathbf{r}) = A \int_{-\infty}^{+\infty} \rho_A(\mathbf{r}, z)\mathrm{d}z$ is the nuclear thickness function with nucleon density distribution $\rho_A(\mathbf{r}, z)$. The maximum possible value of r in the nuclear overlapping zone can be estimated from the equation

$$\max\{r_1(r = r_{\max}), r_2(r = r_{\max})\} = R_A \qquad (4)$$

(R_A is the radius of the nucleus A). This gives

$$r_{\max} = \min\left\{\sqrt{R_A^2 - \frac{b^2}{4}\sin^2\psi} + \frac{b}{2}\cos\psi, \sqrt{R_A^2 - \frac{b^2}{4}\sin^2\psi} - \frac{b}{2}\cos\psi\right\}. \qquad (5)$$

In particular, for the uniform nucleon density distribution, $\rho_A^{\mathrm{un}}(\mathbf{R}) = \rho_0 \cdot \Theta(R_A - |\mathbf{R}|)$, the nuclear overlap function is equal to $T_A^{\mathrm{un}}(r) = 3A(R_A^2 - r^2)^{1/2}/(2\pi R_A^3)$. Then the distribution $P_{AA}^{\mathrm{un}}(\mathbf{r}, b)$ is proportional to

$$P_{AA}^{\mathrm{un}}(\mathbf{r}, b) \propto \sqrt{R_A^2 - r_1^2(r,\psi,b)} \cdot \sqrt{R_A^2 - r_2^2(r,\psi,b)}. \qquad (6)$$

For central AA collisions ($b = 0$, $r_{\max} = R_A$) we simply get $P_{AA}^{\mathrm{un}}(\mathbf{r}, b = 0) \propto (R_A^2 - r^2)$.

It is straightforward to evaluate the time $\tau_L = L$ it takes for a jet to traverse the dense zone:

$$\tau_L = \min\left\{\sqrt{R_A^2 - r_1^2\sin^2\varphi} - r_1\cos\varphi, \sqrt{R_A^2 - r_2^2\sin^2(\varphi-\varphi_0)} - r_2\cos(\varphi-\varphi_0)\right\}, \qquad (7)$$

where φ is the azimuthal angle which determines the direction of the jet motion in the transverse plane, and φ_0 is the angle between the vectors $\mathbf{r}_1$ and $\mathbf{r}_2$. The expression for

$$\varphi_0 = \arccos\frac{r^2 - b^2/4}{r_1 r_2} \qquad (8)$$

can be obtained from the condition

$$r_1 r_2 \cos\varphi_0 = \mathbf{r}_1 \cdot \mathbf{r}_2 = (-b/2 - r\cos\psi)\cdot(b/2 - r\cos\psi) + r^2\sin^2\psi = r^2 - b^2/4. \qquad (9)$$

Finally, we are going to estimate the dependence of the initial energy density in the nuclear overlapping zone on the impact parameter of the collision. At collider energies the minijet system (the semi-hard gluons, quarks and antiquarks with $p_{\mathrm{T}} \gtrsim p_0 \sim 1 \div 2\,\mathrm{GeV}/c$) in the central rapidity region is typically formed in parton–parton scatterings at very early times, $\tau_0 \sim 1/p_{\mathrm{T}} \lesssim 1/p_0 \sim 0.1\,\mathrm{fm}/c$, and this will then serve as the initial condition for the further evolution of the system [19]. Strictly speaking, the soft particle production mechanisms (like the decay of the colour field) can also contribute to the initial conditions in nuclear interactions. However, the relative strength of the soft part decreases strongly with increasing c.m.s. energy of the ion beams. In particular, at LHC energies, $s^{1/2} = 7\,\mathrm{TeV}\times(2Z/A)$ per nucleon pair, the hard and semi-hard processes contribute over 80% to the transverse energy in heavy ion collisions [19]. Moreover, soft processes with small momentum transfer $Q^2 \sim \Lambda_{QCD}^2 \simeq (200\,\mathrm{MeV})^2 \ll p_0^2$ can be partially or fully suppressed, owing to screening of the colour interaction in the dense parton matter produced from the system of minijets in the early stages of the reaction [28]. Therefore, at LHC energies, we will consider only the dominant semi-hard contribution to the formation of the initial state.

The initial energy density inside the comoving volume of longitudinal size $\Delta z = \tau_0 \cdot 2\Delta y$ can be estimated using the Bjorken formula [29,19] as

$$\varepsilon(\tau = \tau_0) = \frac{\langle E_{\mathrm{T}}^A(|y| < \Delta y)\rangle}{S(b)\cdot\Delta z} = \frac{\langle E_{\mathrm{T}}^A(|y| < \Delta y)\rangle \cdot p_0}{S(b)\cdot 2\Delta y}, \qquad (10)$$

where

$$S_{AA}(b) = \int_0^{2\pi} \mathrm{d}\psi \int_0^{r_{\max}} r\mathrm{d}r = \left(\pi - 2\arcsin\frac{b}{2R_A}\right)R_A^2 - b\sqrt{R_A^2 - \frac{b^2}{4}} \qquad (11)$$

is the effective transverse area of the nuclear overlapping zone at impact parameter b. The total initial transverse energy deposition in the mid-rapidity region can be calculated [19] to be

$$\langle E_{\mathrm{T}}^A(b, \sqrt{s}, p_0, |y| < \Delta y)\rangle = T_{AA}(b)\cdot\sigma_{NN}^{\mathrm{jet}}(\sqrt{s}, p_0)\cdot\langle p_{\mathrm{T}}\rangle, \qquad (12)$$

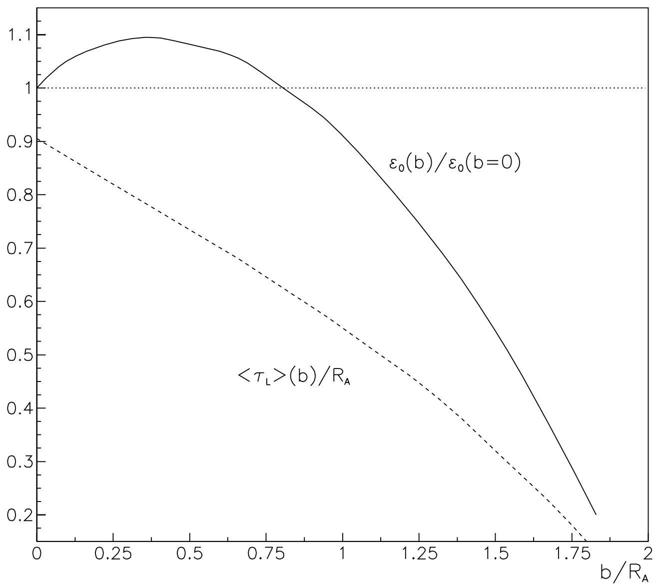

Fig. 2. The impact parameter dependence of the initial energy density $\varepsilon_0(b)/\varepsilon_0(b=0)$ in the nuclear overlapping zone (solid curve), and the average proper time $\langle\tau_L\rangle/R_A$ of a jet escaping from the dense matter (dashed curve) for a uniform nucleon density distribution

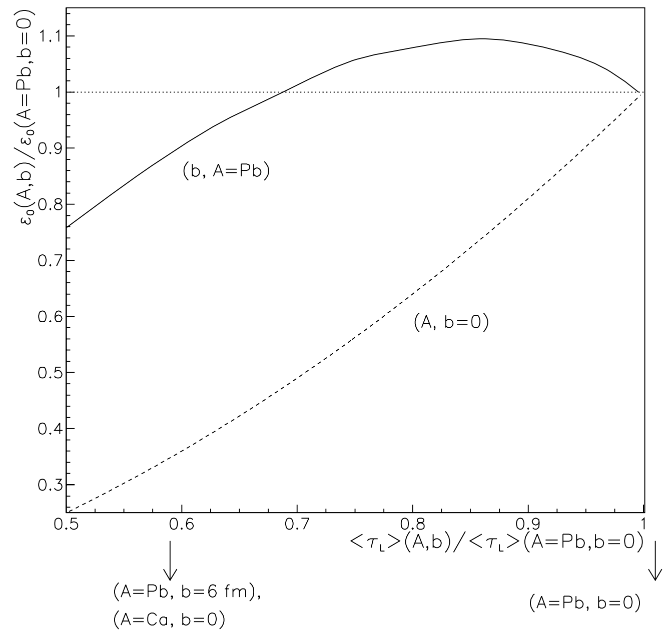

Fig. 3. The initial energy density $\varepsilon_0(A,b)/\varepsilon_0(A=\mathrm{Pb}, b=0)$ in the nuclear overlapping zone versus the average proper time $\langle\tau_L\rangle(A,b)/\langle\tau_L\rangle(A=\mathrm{Pb}, b=0)$ of a jet escaping from the dense matter for varying atomic weight A at fixed $b=0$ (dashed curve), and impact parameter b at fixed $A=\mathrm{Pb}=207$ (solid curve) for a uniform nucleon density distribution

where the first p_T-moment of the inclusive differential minijet cross section $\sigma_{NN}^{\mathrm{jet}} \cdot \langle p_\mathrm{T}\rangle$ is determined by the dynamics of the nucleon–nucleon interactions at the corresponding c.m.s. energy. Then the dependence of the initial energy density ε_0 in the nuclear overlapping zone on the impact parameter b has the form

$$\varepsilon_0(b) \propto T_{AA}(b)/S_{AA}(b), \tag{13}$$

or

$$\varepsilon_0(b) = \varepsilon_0(b=0)\frac{T_{AA}(b)}{T_{AA}(b=0)}\frac{S_{AA}(b=0)}{S_{AA}(b)}. \tag{14}$$

For central AA collisions we have $S_{AA}(b=0) = \pi R_A^2$ and $T_{AA}(b=0) = 9A^2/(8\pi R_A^2)$.

It is worth noting that although this simple geometrical model for jet production in nucleus–nucleus collisions is formally applicable up to an impact parameter of $b = 2R_A$, the major informative domain of our interest is central and semi-central collisions with $b \lesssim R_A$ only. We have the following reasons in favour of this.

(1) The contribution of such events to the total jet rate is dominant, although these events represent only a few percent of the total inelastic AA cross section [30]. For example, the Pb–Pb collisions with impact parameter $b < 0.9R_{\mathrm{Pb}} = 6\,\mathrm{fm}$ contribute $\approx 50\%$ to the total dijet rate at LHC energy, their relative fraction of total cross section being only $\approx 10\%$ in this case [26].

(2) In the most central heavy ion collisions the maximum initial energy density is expected to be achieved in a fairly large (compared with typical hadronic scales) volume, when the effect of super-dense and hot matter formation, like quark–gluon plasma, really may be observable.

The result for the impact parameter dependence of the initial energy density ε_0 (14) in the nuclear overlapping zone for a uniform nucleon density is shown in Fig. 2: it is very weakly dependent of b ($\delta\varepsilon_0 \lesssim 10\%$) up to $b \sim R_A$, and decreases rapidly at $b \gtrsim R_A$. On the other hand, the averaged (over all possible jet production vertices) proper time $\langle\tau_L\rangle$ (7) of a jet escaping from the dense zone is found to decrease almost linearly with increasing impact parameter b (see the second curve in Fig. 2). Therefore the variation of the impact parameter b of the nucleus–nucleus collision (which can be measured, for example, using the total transverse energy deposition detected in different parts of the calorimeters [31]) up to $b \sim R_A$ gives the possibility to study jet quenching as a function of distance traversed without significantly changing the initial energy density ε_0.

Meanwhile, the weakness of the b-dependence of ε_0 gives us an advantage as compared with using beams of different ions at a fixed bin of the impact parameter distribution, when the scaling $\langle\tau_L\rangle(b=0) \propto R_A \propto A^{1/3}$ exists. Equation (13) gives $\varepsilon_0(b=0) \propto A^2/R_A^4$, i.e. $\varepsilon_0(b=0) \propto A^{2/3}$. Figure 3 illustrates the change of the average time $\langle\tau_L\rangle$ of jet travel and the initial energy density ε_0 in the dense zone with variation of the impact parameter b (at fixed $A = 207$, Pb) and atomic weight A (at fixed $b = 0$). For example, decreasing $\langle\tau_L\rangle(\mathrm{Pb}, b=0) \simeq 6\,\mathrm{fm}$ by the factor $\simeq 1.7$ can be obtained by

(a) increasing b up to $b = 0.9R_{\mathrm{Pb}} \simeq 6\,\mathrm{fm}$ (at the expense of only $\sim 10\%$ of ε_0 reduction);

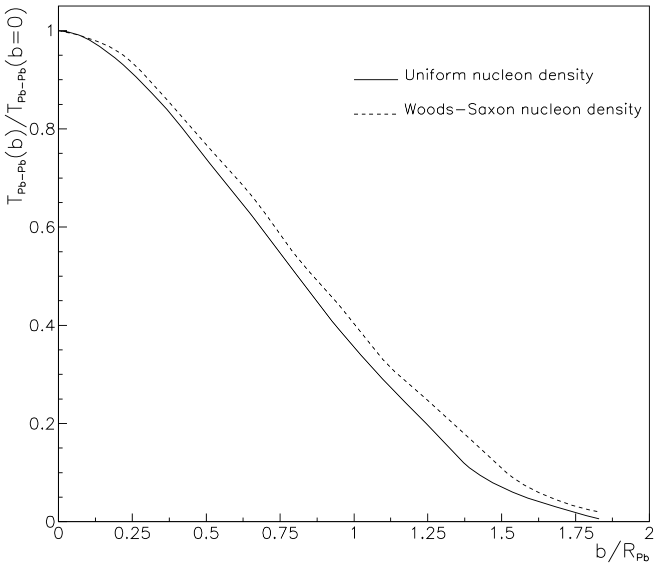

Fig. 4. The nuclear overlap profile function $T_{AA}(b)/T_{AA}(b = 0)$ for uniform (solid) and Woods–Saxon (dashed) nucleon densities in Pb–Pb collisions

(b) decreasing A down to $A = 40$ (Ca) (at the expense of ε_0 reduction by the factor ~ 3).
(3) It is well known that the uniform nucleon density distribution in the nucleus, $\rho_A^{\rm un}(\mathbf{R}) = \rho_0 \cdot \Theta(R_A - |\mathbf{R}|)$, can serve as a good approximation for central and semi-central collisions (see Fig. 4, which shows the nuclear overlap function profile for the uniform[1] and the standard Woods–Saxon nucleon densities). The edge effects near the surface of the nucleus, the impact parameter dependence of the nuclear parton structure functions ("nuclear shadowing") [32], the early transverse expansion of the system and other potentially important phenomena for peripheral ($b \sim 2R_A$) collisions are beyond our considerations here.

3 Impact parameter dependence of the jet energy losses

The intensity of the final state rescattering and collisional and radiative energy losses of hard jet partons in dense QCD matter, created in the nuclear overlapping zone, are sensitive to their initial parameters (energy density, formation time) and space-time evolution [18]. In order to analyse the impact parameter dependence of the jet energy losses and jet quenching, we treat the medium as a boost-invariant longitudinally expanding quark–gluon fluid, and partons as being produced on a hyper-surface of equal proper times $\tau = (t^2 - z^2)^{1/2}$ [29]. We expect that this is an adequate approximation for central and semi-central collisions for our semi-qualitative discussion.

The approach relies on accumulative energy losses, when both the initial and final state gluon radiation is associated with each scattering in the expanding medium, including the interference effect by the modified radiation spectrum as a function of the decreasing temperature $\mathrm{d}E/\mathrm{d}x(T)$. Note that recently the radiative energy losses of a fast parton propagating through expanding (according to Bjorken's model) QCD plasma have been explicitly evaluated in [13] as $\mathrm{d}E/\mathrm{d}x|_{\rm expanding} = c \cdot \mathrm{d}E/\mathrm{d}x|_{T_L}$ with a numerical factor $c \sim 2$ (6) for a parton created inside (outside) the medium, T_L being the temperature at which the dense matter was left [13].

The total energy losses in a transverse direction experienced by a hard parton due to multiple scattering in matter are the result of averaging over the jet production vertex $P_{AA}(\mathbf{r}, b)$, see (2), the transfer momentum squared t in a single rescattering and space-time evolution of the medium:

$$\langle \Delta E_{\rm T}(b) \rangle = \int\limits_0^{2\pi} \mathrm{d}\psi \int\limits_0^{r_{\max}} r \cdot \mathrm{d}r \frac{T_A(r_1) \cdot T_A(r_2)}{T_{AA}(b)} \times \int\limits_0^{2\pi} \frac{\mathrm{d}\varphi}{2\pi} \int\limits_{\tau_0}^{\tau_L} \mathrm{d}\tau \left(\frac{\mathrm{d}E}{\mathrm{d}x}^{\rm rad}(\tau) + \sum_b \sigma_{ab}(\tau) \cdot \rho_b(\tau) \cdot \nu(\tau) \right). \quad (15)$$

Here τ_0 and τ_L, see (7), are the proper time of the plasma formation and the time of the jet escaping from the dense zone, respectively; $\rho_b \propto T^3$ is the density of the plasma constituents of type b at temperature T, σ_{ab} is the integral cross section of scattering of the jet parton a off the comoving constituent b (with the same longitudinal rapidity y), and ν and $\mathrm{d}E/\mathrm{d}x^{\rm rad}$ are the thermally averaged collisional energy loss of a jet parton due to single elastic scattering and the radiative energy losses per unit distance, respectively.

If the mean free path of a hard parton is larger than the screening radius in the QCD medium, $\lambda \gg \mu_D^{-1}$, the successive scatterings can be treated as independent [11]. The transverse distance between successive scatterings, $\Delta r_i = (\tau_{i+1} - \tau_i) \cdot v_{\rm T} = (\tau_{i+1} - \tau_i) \cdot p_{\rm T}/E$, is determined in linear kinetic theory according to the probability density:

$$\frac{\mathrm{d}P}{\mathrm{d}(\Delta r_i)} = \lambda^{-1}(\tau_{i+1}) \cdot \exp \left(- \int\limits_0^{\Delta r_i} \lambda^{-1}(\tau_i + s) \mathrm{d}s \right), \quad (16)$$

where the mean inverse free path is given by $\lambda_a^{-1}(\tau) = \sum_b \sigma_{ab}(\tau) \rho_b(\tau)$.

The dominant contribution to the differential cross section $\mathrm{d}\sigma/\mathrm{d}t$ for scattering of a parton with energy E

[1] Moreover, in this case the explicit form of $T_{AA}(b)$ can be obtained:

$$T_{AA}^{\rm un}(b) = T_{AA}^{\rm un}(b=0) \left[1 - \tilde{b} \left[1 + \left(1 - \frac{\tilde{b}}{4} \right) \ln \frac{1}{\tilde{b}} + 2 \left(1 - \frac{\tilde{b}}{4} \right) \times \left(\ln (1 + \sqrt{1 - \tilde{b}}) - \frac{\sqrt{1 - \tilde{b}}}{1 + \sqrt{1 - \tilde{b}}} \right) - \frac{\tilde{b}(1 - \tilde{b})}{2(1 + \sqrt{1 - \tilde{b}})^2} \right] \right],$$

$\tilde{b} = b^2/(4R_A^2)$, the weak b-dependence of ε_0 and approximately linear drop of $\langle \tau_L \rangle(b)$ being derived for $b \lesssim R_A$ analytically off the "thermal" partons with energy (or effective mass) $m_0 \sim 3T \ll E$ at temperature T can be written as [11,33]

$$\frac{\mathrm{d}\sigma_{ab}}{\mathrm{d}t} \cong C_{ab} \frac{2\pi\alpha_s^2(t)}{t^2}, \tag{17}$$

where $C_{ab} = 9/4, 1, 4/9$ for gg, gq and qq scatterings, respectively;

$$\alpha_s = \frac{12\pi}{(33 - 2N_f)\ln(t/\Lambda_{\mathrm{QCD}}^2)} \tag{18}$$

is the QCD running coupling constant for N_f active quark flavours, and Λ_{QCD} is the QCD scale parameter which is of the order of the critical temperature, $\Lambda_{\mathrm{QCD}} \simeq T_c$. The integrated parton scattering cross section,

$$\sigma_{ab} = \int\limits_{\mu_D^2(\tau)}^{m_0(\tau)E/2} \mathrm{d}t \frac{\mathrm{d}\sigma_{ab}}{\mathrm{d}t}, \tag{19}$$

is regularised by the Debye screening mass squared μ_D^2.

The collisional energy losses due to elastic scattering with high-momentum transfer have originally been estimated by Bjorken in [16], and recalculated later in [17] taking also into account the loss with low-momentum transfer dominated by the interactions with plasma collective modes. Since the latter process contributes to the total collisional energy losses without the large factor $\sim \ln(E/\mu_D)$ in comparison with high-momentum scattering, and since it can be effectively "absorbed" by the redefinition of minimal $t \sim \mu_D^2$ under the numerical estimates, we shall concentrate on collisional energy losses with high-momentum transfer only[2]. The thermal average of such losses can be written as

$$\nu = \left\langle \frac{t}{2m_0} \right\rangle = \frac{1}{2}\left\langle \frac{1}{m_0} \right\rangle \cdot \langle t \rangle \simeq \frac{1}{4T\sigma_{ab}} \int\limits_{\mu_D^2}^{3TE/2} \mathrm{d}t \frac{\mathrm{d}\sigma_{ab}}{\mathrm{d}t} t. \tag{20}$$

The value ν is independent of the total distance traversed and is determined by temperature, with roughly $\nu \propto T$. Then the total collisional energy losses integrated over the whole jet path are estimated as $\langle \Delta E_{\mathrm{col}} \rangle \propto T_0^2 \propto (\varepsilon_0)^{1/2}$, as has been pointed out in [16]. The τ_L-dependence of ΔE_{col} can be weaker than linear for an expanding medium ($\Delta E_{\mathrm{col}} \propto \tau_L$ for static matter).

The energy spectrum of coherent medium-induced gluon radiation and the corresponding dominated part of the radiative energy losses, $\mathrm{d}E/\mathrm{d}x$, were analysed in [12, 13] by means of the Schrödinger-like equation whose "potential" is determined by the single-scattering cross section of the hard parton in the medium. For the quark produced in the medium it gives [13,14][3]

$$\frac{\mathrm{d}E}{\mathrm{d}x}^{\mathrm{rad}} = \frac{2\alpha_s C_R}{\pi\tau_L} \int\limits_{\omega_{\min}}^{E} \mathrm{d}\omega \left[1 - y + \frac{y^2}{2}\right] \times \ln|\cos(\omega_1\tau_1)|, \tag{21}$$

$$\omega_1 = \sqrt{\mathrm{i}\left(1 - y + \frac{C_R}{3}y^2\right)\bar{\kappa}\ln\frac{16}{\bar{\kappa}}} \quad \text{with} \quad \bar{\kappa} = \frac{\mu_D^2\lambda_g}{\omega(1-y)}. \tag{22}$$

Here $\tau_1 = \tau_L/(2\lambda_g)$, and $y = \omega/E$ is the fraction of the hard parton energy carried by the radiated gluon, and $C_R = 4/3$ is the quark colour factor. A similar expression for the gluon jet can be obtained by substituting $C_R = 3$ and a proper change of the factor in the square bracket in (21), see [13]. The integral (21) is carried out over all energies from $\omega_{\min} = E_{LPM} = \mu_D^2\lambda_g$ (λ_g is the gluon mean free path), the minimal radiated gluon energy in the coherent LPM regime, up to initial jet energy E. The complex form of the expression (21) does not allow us in the general case to extract the explicit form of the τ_L- and T-dependences of $\mathrm{d}E/\mathrm{d}x^{\mathrm{rad}}$. In the limit of a "strong" LPM effect, $\omega \gg \mu_D^2\lambda_g$, we have [12,13,21] $\mathrm{d}E/\mathrm{d}x^{\mathrm{rad}} \propto T^3$ and $\mathrm{d}E/\mathrm{d}x^{\mathrm{rad}} \propto \tau_L$ with logarithmic accuracy. Then the total radiative energy losses $\langle \Delta E_{\mathrm{rad}} \rangle = \int \mathrm{d}\tau \cdot \mathrm{d}E/\mathrm{d}x^{\mathrm{rad}}$ are estimated as $\langle \Delta E_{\mathrm{rad}} \rangle \propto T_0^3 \propto \varepsilon_0^{3/4}$ and $\Delta E_{\mathrm{rad}} \propto \tau_L^{\beta}$, where $\beta \lesssim 2$ for an expanding medium ($\beta \sim 2$ in the case of static matter).

In order to simplify numerical calculations (and not to introduce new parameters) we omit the transverse expansion and viscosity of the fluid using the well-known scaling solution due to Bjorken [29] for a temperature and density of QGP at $T > T_c \simeq 200\,\mathrm{MeV}$:

$$\varepsilon(\tau)\tau^{4/3} = \varepsilon_0\tau_0^{4/3},\ T(\tau)\tau^{1/3} = T_0\tau_0^{1/3}, \quad \rho(\tau)\tau = \rho_0\tau_0. \tag{23}$$

Let us remark that the influence of the transverse flow, as well as of the mixed phase at $T = T_c$, on the intensity of jet rescattering (which is a strongly increasing function of T) seems to be inessential for high initial temperatures $T_0 \gg T_c$ [18]. On the contrary, the presence of viscosity slows down the cooling rate, which leads to a jet parton spending more time in the hottest regions of the medium. As a result the rescattering intensity goes up, i.e., in fact the effective temperature of the medium is increased as

[2] Anyway, high- and low-momentum parts of the collisional energy losses have the same dependence on distance traversed

[3] The gluons with formation times τ_f exceeding the time $\tau_L = L$ that are formed outside the medium (the factorization *medium-independent* component) carry away a fraction of the initial parton energy proportional to $\alpha_s(E)$. This part of the gluon radiation produces the standard jet energy profile which is identical to that of a jet produced in a hard process in vacuum. Hereafter we shall concentrate on the *medium-dependent* effects and will not include the "vacuum" part of the jet profile

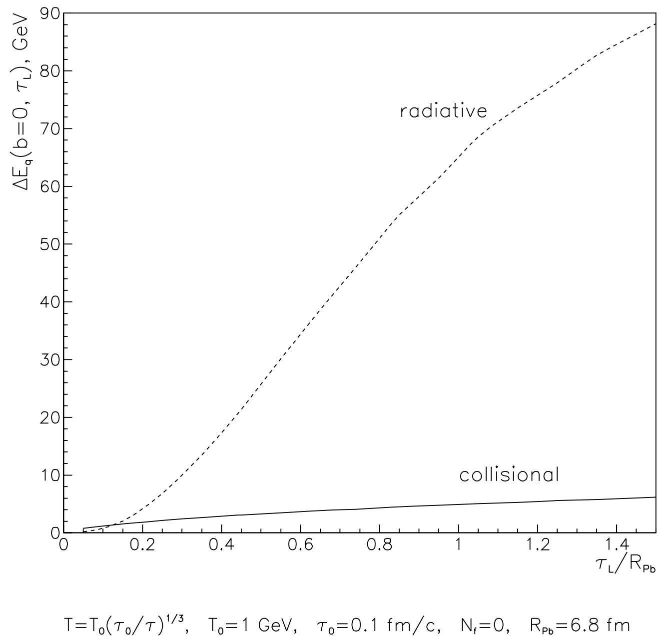

Fig. 5. The medium-induced radiative (dashed) and collisional (solid) energy losses of a quark-initiated jet with initial energy $E_{\rm T}^q = 100\,{\rm GeV}$ versus the average proper time $\langle \tau_L \rangle / R_{\rm Pb}$ of a jet escaping from the dense matter

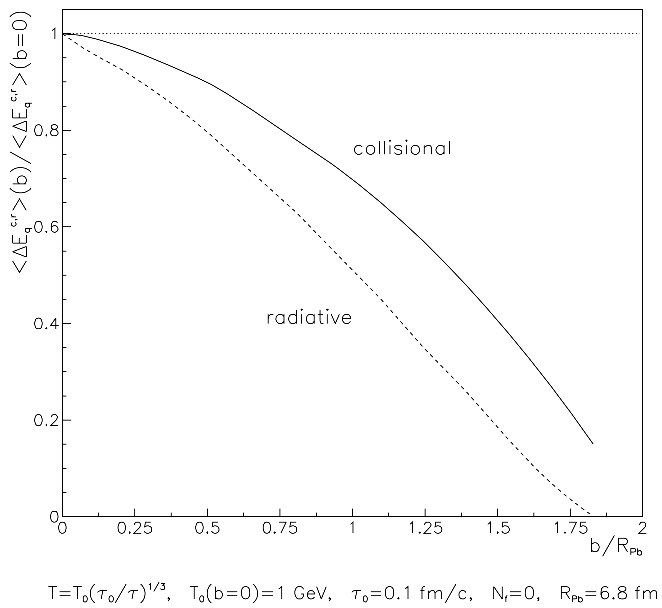

Fig. 6. The impact parameter dependence of the medium-induced radiative (dashed) and collisional (solid) energy losses of a quark-initiated jet with initial energy $E_{\rm T}^q = 100\,{\rm GeV}$ normalised to the corresponding average values at zero impact parameter

compared with the perfect QGP case [18]. We also do not take into account here the probability of jet rescattering in nuclear matter, because the intensity of this process and the corresponding contribution to the total energy losses are not significant due to the much smaller energy density in "cold" nuclei. For certainty we used the initial conditions for the gluon-dominated plasma formation ($N_f \approx 0$, $\rho_q \approx 1.95 T^3$) expected for central Pb–Pb collisions at LHC [19]: $\tau_0 \simeq 0.1\,{\rm fm}/c$, $T_0 \simeq 1\,{\rm GeV}$[4].

Figure 5 represents the calculated τ_L-dependence of coherent medium-induced radiative and collisional energy losses of a quark-initiated jet with initial energy $E_{\rm T}^q = 100\,{\rm GeV}$. We see that the τ_L-dependence of radiative and collisional losses is very different: $\Delta E_{\rm rad}(\tau_L)$ grows somewhat stronger than linearly; meanwhile $\Delta E_{\rm col}(\tau_L)$ looks rather logarithmic. This results in the corresponding difference in the impact parameter dependence of radiative and collisional losses, the normalised profiles of which are presented in Fig. 6. To make the plot more visual the energy losses $\langle \Delta E_{\rm T}(b) \rangle$ are normalised to the corresponding average values at zero impact parameter, $\langle \Delta E_{T\ {\rm rad}}^q (b = 0) \rangle \sim 45\,{\rm GeV}$ and $\langle \Delta E_{T\ {\rm col}}^q (b = 0) \rangle \sim 5\,{\rm GeV}$ for the parameters used.

For example, increasing of the impact parameter from $b = 0$ to $b = R_{\rm Pb}$ gives a reduction of $\sim 30\%$ collisional and $\sim 50\%$ radiative losses. We have also found that the form of the b-dependence of collisional losses is almost independent of the scenarios of the space-time evolution of QGP (perfect or viscous fluid), the b-dependence of the radiative losses being somewhat more sensitive to these effects.

Note that the choice of the scale for a minimal jet energy $E_{\rm T}^q \sim 100\,{\rm GeV}$ corresponds to the threshold for "true" QCD jet recognition against the "thermal" background jets (statistical fluctuations of the transverse energy flux) with reconstruction efficiency close to 1 in heavy ion collisions at LHC [26,27,34]. We hope that the separation of the collisional and the radiative contribution to the total energy losses, when doing the experimental data analysis for jets with finite cone size, could be performed based on the essential difference of their angular distribution [20,21]: the radiative losses are expected to dominate at small jet cone size θ_0, while the relative contribution to collisional losses grows with increasing θ_0.

4 Impact parameter dependence of dijet production rate

In the previous section we have analysed the impact parameter dependence of the jet energy losses, which can be directly observed in $\gamma +$jet and $Z(\to \mu^+\mu^-)+$jet production processes. Another observable effect is a suppression of the high-$p_{\rm T}$ jet pair yield (dijet quenching) due to final state rescattering and energy losses. In connection with this, we would like to estimate the impact parameter dependence of the jet + jet production rates in heavy ion collisions. The observed number of $\{ij\}$ type dijets with transverse momenta $p_{\rm T1}, p_{\rm T2}$ produced in initial hard scat-

[4] These estimates are of course rather approximate and model-depending: the discount of higher order $\alpha_{\rm s}$ terms, uncertainties of the structure functions in the low-x region, and nuclear shadowing can result in variations of the initial energy density [19]

tering processes in minimum bias AA collisions is written as

$$\frac{\mathrm{d}N_{ij}^{\mathrm{dijet}}}{\mathrm{d}p_{\mathrm{T}1}\mathrm{d}p_{\mathrm{T}2}} = \int\limits_0^\infty \mathrm{d}^2 b \frac{\mathrm{d}^2\sigma_{\mathrm{jet}}^0}{\mathrm{d}^2 b} \cdot \frac{\mathrm{d}N_{ij}^{\mathrm{dijet}}}{\mathrm{d}p_{\mathrm{T}1}\mathrm{d}p_{\mathrm{T}2}}(b) \Bigg/ \int\limits_0^\infty \mathrm{d}^2 b \frac{\mathrm{d}^2\sigma_{\mathrm{jet}}^0}{\mathrm{d}^2 b}, \quad (24)$$

$$\frac{\mathrm{d}N_{ij}^{\mathrm{dijet}}}{\mathrm{d}p_{\mathrm{T}1}\mathrm{d}p_{\mathrm{T}2}}(b) = \int\limits_0^{2\pi} \mathrm{d}\psi \int\limits_0^{r_{\max}} r\mathrm{d}r T_A(r_1) T_A(r_2) \times \int\limits_0^{2\pi} \frac{\mathrm{d}\varphi}{2\pi} \int \mathrm{d}p_{\mathrm{T}}^2 \frac{\mathrm{d}\sigma_{ij}}{\mathrm{d}p_{\mathrm{T}}^2} \delta(p_{\mathrm{T}1} - p_{\mathrm{T}} + \Delta E_{\mathrm{T}}^i(r,\psi,\varphi,b))\delta(p_{\mathrm{T}2} - p_{\mathrm{T}} + \Delta E_{\mathrm{T}}^j(r,\psi,\pi-\varphi,b)), \quad (25)$$

where the parton differential cross section $\mathrm{d}\sigma_{ij}/\mathrm{d}p_{\mathrm{T}}^2$ is calculated in perturbative QCD:

$$\frac{\mathrm{d}\sigma_{ij}}{\mathrm{d}p_{\mathrm{T}}^2} = K \int \mathrm{d}x_1 \int \mathrm{d}x_2 \int \mathrm{d}\hat{t} f_i(x_1, p_{\mathrm{T}}^2) f_j(x_2, p_{\mathrm{T}}^2) \times \frac{\mathrm{d}\hat{\sigma}_{ij}}{\mathrm{d}\hat{t}} \delta\left(p_{\mathrm{T}}^2 - \frac{\hat{t}\hat{u}}{\hat{s}}\right); \quad (26)$$

$\mathrm{d}\widehat{\sigma_{ij}}/\mathrm{d}\hat{t}$ expresses the differential cross section for parton–parton scattering as a function of the kinematical Mandelstam variables $\hat{s}$, $\hat{t}$ and $\hat{u}$, $f_{i,j}$ are the structure functions, x is the nucleon-momentum fraction carried by a parton, and the correction factor K takes into account higher order contributions. We have performed tests with the program of Ellis et al. [35] to verify that next-to-leading order (NLO) corrections are insignificant ($K \sim 1$) for jets with $p_{\mathrm{T}} \geq 50 \div 100\,\mathrm{GeV}/c$ and the reasonable cone radius in the (y,ϕ)-plane of $R = 0.3 \div 0.5$ (see also [36]). Note also that the region of sufficiently hard jets, $x_{1,2} \sim (\hat{s}/s)^{1/2} \gtrsim 0.2$, hardly is affected by the initial state nuclear interactions, like gluon depletion ("nuclear shadowing" of the nucleon structure functions) [37]. Anyway, the dijet rate, integrated above the threshold value $p_{\mathrm{T}}^{\mathrm{cut}}$,

$$R_{AA}^{\mathrm{dijet}}(p_{\mathrm{T}1}, p_{\mathrm{T}2} > p_{\mathrm{T}}^{\mathrm{cut}}) = \int\limits_{p_{\mathrm{T}}^{\mathrm{cut}}} \mathrm{d}p_{\mathrm{T}1} \int\limits_{p_{\mathrm{T}}^{\mathrm{cut}}} \mathrm{d}p_{\mathrm{T}2} \sum_{i,j} \left(\frac{\mathrm{d}N_{ij}^{\mathrm{dijet}}}{\mathrm{d}p_{\mathrm{T}1}\mathrm{d}p_{\mathrm{T}2}}\right)_{AA}, \quad (27)$$

in AA relative to pp collisions can be studied by introducing a reference process, unaffected by energy losses and with a production cross section proportional to the number of nucleon–nucleon collisions, such as Drell–Yan dimuons or (suitable for LHC [24]) $Z(\to \mu^+\mu^-)$ production,

$$R_{AA}^{\mathrm{dijet}}/R_{pp}^{\mathrm{dijet}} = \left(\sigma_{AA}^{\mathrm{dijet}}/\sigma_{pp}^{\mathrm{dijet}}\right) / \left(\sigma_{AA}^{\mathrm{DY}(Z)}/\sigma_{pp}^{\mathrm{DY}(Z)}\right). \quad (28)$$

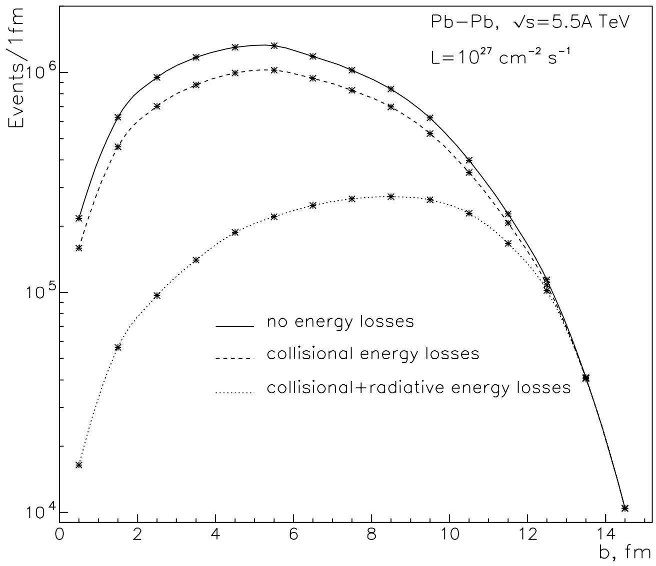

Fig. 7. The jet + jet rates for $E_{\mathrm{T}}^{\mathrm{jet}} > 100\,\mathrm{GeV}$ and $|y^{\mathrm{jet}}| < 2.5$ in different impact parameter bins for various cases: without energy losses (solid curve), with collisional losses only (dashed curve), and with collisional and radiative losses (dotted curve). The rates are normalised to the expected number of events produced in Pb–Pb collisions during two weeks of LHC running at a luminosity of $L = 10^{27}\,\mathrm{cm}^{-2}\mathrm{s}^{-1}$

The cross section $\mathrm{d}^2\sigma_{\mathrm{jet}}^0/\mathrm{d}^2 b$ for initially produced jets in AA collisions at given b can be written as [26,30]

$$\frac{\mathrm{d}^2\sigma_{\mathrm{jet}}^0}{\mathrm{d}^2 b}(\mathbf{b},\sqrt{s}) = T_{AA}(\mathbf{b})\sigma_{NN}^{\mathrm{jet}}(\sqrt{s})\frac{\mathrm{d}^2\sigma_{in}^{AA}}{\mathrm{d}^2 b}(\mathbf{b},\sqrt{s}), \quad (29)$$

where the nucleon–nucleon collision cross section of the hard process $\sigma_{NN}^{\mathrm{jet}}$ has been computed with the PYTHIA model [38]. The differential inelastic AA cross section is calculated to be

$$\frac{\mathrm{d}^2\sigma_{\mathrm{in}}^{AA}}{\mathrm{d}^2 b}(\mathbf{b},\sqrt{s}) = \left[1 - \left(1 - \frac{1}{A^2}T_{AA}(\mathbf{b})\sigma_{NN}^{\mathrm{in}}(\sqrt{s})\right)^{A^2}\right], \quad (30)$$

with the inelastic non-diffractive nucleon–nucleon cross section $\sigma_{NN}^{\mathrm{in}}$ ($\simeq 60\,\mathrm{mb}$ for $s^{1/2} = 5.5\,\mathrm{TeV}$).

Figure 7 shows the dijet rates $\sigma_{AA}^{\mathrm{in}} R_{AA}^{\mathrm{dijet}} L$ ($E_{\mathrm{T}}^{\mathrm{jet}} > p_{\mathrm{T}}^{\mathrm{cut}} = 100\,\mathrm{GeV}$, with the rapidity window $|y^{\mathrm{jet}}| < 2.5$) in different impact parameter bins for three cases:

(i) without energy losses,
(ii) with collisional losses only,
(iii) with collisional and radiative losses.

The rates are normalised to the expected number of events produced in Pb–Pb collisions during two weeks ($1.2 \times 10^6\,\mathrm{s}$) of a LHC run time, assuming a luminosity $L = 10^{27}\,\mathrm{cm}^{-2}\mathrm{s}^{-1}$ [27]. The total initial dijet rate with $E_{\mathrm{T}}^{\mathrm{jet}} > 100\,\mathrm{GeV}$ is estimated to be 1.1×10^7 events ($gg \to gg \simeq 60\%$, $qg \to qg \simeq 30\%$, $qq, gg \to qq \simeq 10\%$). Since the dijet quenching is much stronger in central collisions than in peripheral ones, the maximum and mean values of the $\mathrm{d}N^{\mathrm{dijet}}/\mathrm{d}b$ distribution get shifted towards larger b. The corresponding result for jets with non-zero cone size θ_0 is expected to be somewhere between the cases (iii) ($\theta_0 \to 0$) and (ii). The observation of a dramatic change in the b-dependence of the dijet rates in heavy ion collisions as compared to what is expected from an independent nucleon–nucleon interactions pattern would indicate the existence of medium-induced parton rescattering.

As we have mentioned above, the measurement of the centrality of events can be performed from total transverse energy deposition $E_{\mathrm{T}}^{\mathrm{tot}}$ in the calorimeters, which strongly decreases from central to peripheral collisions [31], roughly as $\overline{E_{\mathrm{T}}^{\mathrm{tot}}}(b) \propto T_{AA}(b)$. If the jet energy losses $\langle \Delta E_{\mathrm{T}}^{\mathrm{jet}} \rangle$ (15) or the dijet production rates R^{dijet} (27) and (28) are measured in different bins of $E_{\mathrm{T}}^{\mathrm{tot}}$, then one can relate the b- and $E_{\mathrm{T}}^{\mathrm{tot}}$-dependences of $F = (\Delta E_{\mathrm{T}}^{\mathrm{jet}}, R^{\mathrm{dijet}})$ using the $E_{\mathrm{T}}^{\mathrm{tot}}$–$b$ correlation functions C_{AA}:

$$F(E_{\mathrm{T}}^{\mathrm{tot}}) = \int \mathrm{d}^2 b F(b) C_{AA}(E_{\mathrm{T}}^{\mathrm{tot}}, b),$$

$$C_{AA}(E_{\mathrm{T}}^{\mathrm{tot}}, b) = \frac{1}{\sqrt{2\pi}\sigma_{E_{\mathrm{T}}}(b)} \exp\left(-\frac{\left(E_{\mathrm{T}}^{\mathrm{tot}} - \overline{E_{\mathrm{T}}^{\mathrm{tot}}}(b)\right)^2}{2\sigma_{E_{\mathrm{T}}}^2(b)}\right), \quad (31)$$

$$F(b) = \int \mathrm{d}E_{\mathrm{T}}^{\mathrm{tot}} F(E_{\mathrm{T}}^{\mathrm{tot}}) C_{AA}(b, E_{\mathrm{T}}^{\mathrm{tot}}),$$

$$C_{AA}(b, E_{\mathrm{T}}^{\mathrm{tot}}) = \frac{1}{\sqrt{2\pi}\sigma_b(E_{\mathrm{T}}^{\mathrm{tot}})} \exp\left(-\frac{\left(b - \overline{b}(E_{\mathrm{T}}^{\mathrm{tot}})\right)^2}{2\sigma_b^2(E_{\mathrm{T}}^{\mathrm{tot}})}\right). \quad (32)$$

The estimate with the HIJING model [39] accuracy of the determination of the impact parameter $\sigma_b(E_{\mathrm{T}}^{\mathrm{tot}}) \sim 1-2\,\mathrm{fm}$ in AA collisions at LHC [40] seems to be enough to observe the above effects.

5 Conclusions

To summarise, we have considered the impact parameter dependence of medium-induced radiative and collisional jet energy losses in dense QCD matter, created in ultrarelativistic heavy ion collisions. We have found that this b-dependence is very different for each mechanism due to coherence effects (the dependence of the radiative energy losses per unit distance $\mathrm{d}E/\mathrm{d}x$ of total distance traversed). As a consequence, the radiative losses are more sensitive to the impact parameter of the nucleus–nucleus collision, which determines the effective volume of the nuclear overlapping dense zone, and the space-time evolution of the medium.

A possible way to directly observe the energy losses at different impact parameter (or total detected E_{T} deposition) bins, involves tagging the hard jet opposite a particle that does not interact strongly, like in γ + jet and $Z(\to \mu^+\mu^-)$ + jet production processes. Since the initial energy density ε_0 in the dense zone depends on b very slightly ($\delta\varepsilon_0 \lesssim 10\%$) up to $b \sim R_A$, studying the b-dependence appears to be advantageous rather than using different ions at fixed impact parameter $b \sim 0$ (when $\varepsilon_0(b \sim 0) \propto A^{2/3}$). We hope that the separation of the collisional and the radiative contribution to the total energy losses when doing the experimental data analysis for jets with finite cone size could be performed based on the essential difference in their angular distributions.

Another process of interest is high-p_{T} jet pair production. The expected statistics for dijet rates in heavy ion collisions at LHC will be large enough to study the impact parameter dependence. Since suppression of the dijet yield (jet quenching) due to medium-induced energy losses should be much stronger in central collisions than in the peripheral ones, the maximum and mean values of the $\mathrm{d}N^{\mathrm{dijet}}/\mathrm{d}b$ distribution predicted to be shifted towards the larger b.

Finally, the study of the impact parameter dependences in the hard jet production processes (jet+jet, γ+jet and Z+jet channels) is important for extracting information about the properties of super-dense QCD matter to be created in heavy ion collisions at LHC.

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