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Parton energy loss in an expanding quark-gluon plasma: Radiative vs collisional

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We perform a comparison of the radiative and collisional parton energy losses in an expanding quark-gluon plasma. The radiative energy loss is calculated within the light-cone path integral approach [4]. The collisional energy loss is calculated using the Bjorken method with an accurate treatment of the binary collision kinematics. Our numerical results demonstrate that for RHIC and LHC conditions the collisional energy loss is relatively small in comparison to the radiative one. We find an enhancement of the heavy quark radiative energy loss as compared to that of the light quarks at high energies.

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1. The jet quenching phenomenon in AA-collisions observed at RHIC (for a review, see [1]) is widely believed to be due to the parton energy loss in the hot quark-gluon plasma (QGP) produced at the initial stage of nucleus-nucleus collision. The parton energy loss may come from the collisional energy loss and the induced gluon radiation. The first estimate of the collisional energy loss has been done by Bjorken [2]. The radiative energy loss has been under active investigation in the last years [3 – 9] (for a review, see [10]). The calculations within the BDMPS [3, 10] and the light-cone path integral (LCPI) [4, 6, 7]) approaches demonstrate that for high energy partons the energy loss is likely dominated by the induced gluon radiation. The estimates given in [11] show that the collisional energy loss can roughly increase the energy loss by 30–40% for RHIC energies. However, it has recently been claimed [12 – 14] that the collisional energy loss may be as important as the radiative one, or even dominate at low energies. But an accurate comparison of the two mechanisms of the energy losses so far has not been performed, say, in [14] even for the plasma with a uniform density the radiative energy loss was calculated incorrectly. The authors have used the kinematical suppression factor for the radiative energy loss obtained in [8] which strongly overestimates the kinematic suppression [15]. The question of the collisional energy loss is becoming of special current interest in connection with the recent data on the non-photonic electrons [16] indicating that the nuclear suppression for heavy quarks may be similar to that for light ones. This fact seems to be inconsistent with purely radiative energy loss which, at low energies, should be suppressed for heavy quarks by the mass effects [17]. Unfortunately, presently for both the radiative and collisional losses the uncertainties in the theoretical predictions are rather big. Say, the results are very sensitive to the choice of α_s (running or constant [11, 18]), to the infrared effects [11]. For clarifying the situation with the relative contributions of the radiative and collisional energy loss it is important to perform the calculations with the same parametrizations of the coupling constant and the infrared cutoffs. In the present paper we present the results of such calculations for expanding QGP for RHIC and LHC conditions.

To evaluate the radiative energy loss we use the LCPI formalism [4, 6, 7] which treats accurately the mass, finite-size and Landau-Pomeranchuk-Migdal (LPM) effects [19]. We evaluate the collisional energy loss using Bjorken's method [2]. However, contrary to the Bjorken analysis we treat accurately the kinematics of the binary collisions and fluctuations of the momentum transfer due to the thermal parton motion in the QGP. For the infrared momentum cutoff we use the Debye mass included into the gluon propagator. A more accurate treatment of the infrared region of collective excitations is possible in the pQCD Hard Thermal Loop (HTL) resummation technique [20], which, however, is unreliable for RHIC and LHC conditions when the plasma temperature $T \lesssim (2-3)T_c$ ($T_c \approx 170\,\mathrm{MeV}$ is the temperature of the deconfinement phase transition) and the coupling constant $g \sim 1.5 - 2$. Say, in the leading order in g the HTL approximation predicts zero magnetic screening mass, but the lattice calculations [21] show that the magnetic mass may be of the order of the electric Debye mass. Also, the HTL approach to the collisional energy loss [22, 23] does not incorporate the running α_s which turns out to be very important [18]. In our calculations of the radiative and collisional energy loss the infrared cutoffs for the t-channel gluon exchanges are performed in the same way. This should minimize the theoretical uncertainties associated with the collective excitations in the ratio of the radiative to collisional energy loss. Numerical calculations are performed with the Debye mass obtained in the lattice calculations.

Our results show that for RHIC and LHC conditions the effect of the collisional energy loss is relatively small, and cannot be crucial for the difference in the nuclear suppression of the heavy and light quark jets. For c-quark we find the radiative energy loss which is very close to that for the light quarks. For b-quark the radiative energy loss is significantly suppressed at $E \lesssim 20\,\mathrm{GeV}$, but at high energies it exceeds the charm (and light quark) energy loss. Our results show that the observed at RHIC suppression of the non-photonic electrons [16] may be naturally described if in the experimentally studied region $p_T \lesssim 8$ GeV the non-photonic electrons are dominated by charm production. Although the pQCD calculations [24] predict that the bottom contribution dominates at $p_T \gtrsim 5\,\mathrm{GeV}$, this possibility cannot be excluded since the pQCD results in the charm mass region are very fragile. Say, the pQCD calculations [24] underestimate the normalization of the electron and D-meson spectra by a factor of about 5 [25, 26]. Thus, presently, it is not clear at all whether the pQCD calculations predict correctly the ratio charm/bottom for the kinematical region studied at RHIC.

2. We begin with the radiative energy loss. For definiteness we consider the case of an energetic quark. We assume that the fast quark, Q, produced in a hard process at $z = 0$, passes through a length L of an expanding QGP (we choose the z-axis along the initial quark momentum). We define the energy loss as

$$\Delta E = E \int\limits_{x_{\min}}^{x_{\max}} dx x \frac{dP}{dx}, \tag{1}$$

where E is the initial quark energy, x is the Feynman variable for the radiated gluon, and dP/dx is the probability distribution of the induced gluon emission. Since for hard gluons with $x \gtrsim 0.5$ the jet really does not disappear, from the point of view of the jet quenching, a reasonable choice for the upper limit of x-integration is $x_{\max} = 0.5$. For $x_{\min}$ we use the value m_g/E, where hereafter m_g is the gluon quasiparticle mass (the quark quasiparticle mass will be denoted m_q). The gluon spectrum may be written as [17, 11]

$$\frac{dP}{dx} = \int\limits_0^L dz\, n(z) \frac{d\sigma_{\mathrm{eff}}^{BH}(x,z)}{dx}. \tag{2}$$

Here $n(z)$ is the medium number density, $d\sigma_{\mathrm{eff}}^{BH}/dx$ is an effective Bethe-Heitler cross section accounting for the LPM and finite-size effects given by [11]

$$\frac{d\sigma_{\mathrm{eff}}^{BH}(x,z)}{dx} = -\frac{\alpha_s C_F[1+(1-x)^2]}{\pi x M} \times \\ \times \mathrm{Im} \int\limits_0^z d\xi \frac{\partial}{\partial \rho}\left(\frac{F(\xi,\rho)}{\sqrt{\rho}}\right)\Big|_{\rho=0}, \tag{3}$$

where $M = Ex(1-x)$ is the reduced "Schrödinger mass", the function F is the solution to the radial Schrödinger equation for the azimuthal quantum number $m = 1$

$$i\frac{\partial F(\xi,\rho)}{\partial \xi} = \Big[-\frac{1}{2M}\left(\frac{\partial}{\partial\rho}\right)^2 - \\ - i\frac{n(z-\xi)\sigma_3(\rho,x)}{2} + \frac{4m^2-1}{8M\rho^2} + \frac{1}{L_f}\Big]F(\xi,\rho) \tag{4}$$

with the boundary condition $F(\xi=0,\rho) = \sqrt{\rho}\sigma_3(\rho,x) \times \times \epsilon K_1(\epsilon\rho)$ (K_1 is the Bessel function). Here $L_f = 2M/\epsilon^2$ with $\epsilon^2 = m_q^2 x^2 + m_g^2(1-x)^2$, $\sigma_3(\rho,x)$ is the cross section of interaction of the $q\bar{q}g$ system with a medium constituent which reads $\sigma_3(\rho,x) = \frac{9}{8}[\sigma_{q\bar{q}}(\rho) + \sigma_{q\bar{q}}((1 - - x)\rho)] - \frac{1}{8}\sigma_{q\bar{q}}(x\rho)$ [27, 4], where

$$\sigma_{q\bar{q}}(\rho) = \alpha_s^2 C_T C_F \int d\mathbf{q}_\perp \frac{[1-\exp(i\mathbf{q}_\perp b\boldsymbol{\rho})]}{(q_\perp^2+\mu_D^2)^2} \tag{5}$$

is the dipole cross section for the color singlet $q\bar{q}$ pair ($C_{F,T}$ are the color Casimir for the quark and thermal parton (quark or gluon), μ_D is the Debye mass).

The above equations are written for fixed μ_D and α_s. For z-dependent μ_D the dipole cross section, besides ρ, depends on z as well. The incorporation of the z-dependence of the dipole cross section does not lead to any problems. To account for the effect of running α_s on the dipole cross section we include $\alpha_s(q^2)$ in the integrand in (5). The inclusion of the running α_s for the emitted gluon is a more delicate question since the gluon spectrum is written in the coordinate representation, and there is no explicit dependence on the parton transverse momenta in the vertex $q \to gq$. To generalize our formulas to the running α_s we use the Schrödinger diffusion relation $\rho \sim \sqrt{\xi/M}$ connecting the longitudinal scale ξ and ρ scale (in $F(\xi=0,\rho)$). In terms of the transverse separation between quark and gluon ρ the transverse gluon momentum can be estimated via the uncertainty relation $q \sim 1/\rho$. Thus it seems quite reasonable to use for the virtuality scale in the splitting vertex the parametrization $q^2 \approx aM/\xi$. We adjusted the coefficient a to reproduce in our formalism the $N = 1$ rescattering contribution (which dominates the gluon spectrum) evaluated in the diagrammatic approach in momentum representation [15]. For our parametrization of the running α_s (discussed in Sec. 4) it gives $a \approx 1.85$.

The induced gluon spectrum can also be used to estimate the effect of the energy gain due to absorption of the thermal gluons. For the plasma with fixed (z-independent) temperature in the collinear gluon approximation one can obtain for the energy gain

$$\Delta E_{\text{gain}} \approx \\ \approx E \int\limits_{x_{\min}}^{x_{\max}} dx x(1+x) n_B(Ex) \frac{dP}{dy}(y = x/(1+x), E'), \quad (6)$$

where $E' = E(1+x)$ is the parton energy after thermal gluon absorption, $n_B(p) = (e^{p/T}-1)^{-1}$ is the thermal gluon distribution. An accurate evaluation of this effect for an expanding plasma in the situation when the plasma density changes strongly at the gluon formation scale L_f is a complicated problem. To estimate the effect of gluon absorption we have simply used for the thermal distribution the distribution averaged over z. Our numerical calculations show that for RHIC and LHC conditions the effect of the gluon absorption is suppressed by about two orders of magnitude as compared to the gluon emission. For this reason it can be safely neglected.

3. Let us now discuss the collisional energy loss. We calculate it through the energy transfer in the binary elastic collisions of the fast parton with the thermal quarks and gluons treated as free particles. In this approximation the collisional energy loss per unit length can be written as

$$\frac{dE}{dz} = \frac{1}{2Ev} \sum_{p=q,g} g_p \int \frac{d\mathbf{p}'}{2E'(2\pi)^3} \int \frac{d\mathbf{k}\, n_p(k)}{2k(2\pi)^3} \times \\ \times \int \frac{d\mathbf{k}'[1+\epsilon_p n_p(k')]}{2k'(2\pi)^3} (2\pi)^4 \delta^4(P+K-P'-K') \times \\ \times \omega \langle |M(s,t)|^2 \rangle \theta(\omega_{\max} - \omega), \quad (7)$$

where $v \approx 1$ is the velocity of the fast quark, $P = (E, \mathbf{p})$ and $K = (k, \mathbf{k})$ are the momenta for incoming partons, $P' = (E', \mathbf{p}')$ and $K' = (k', \mathbf{k}')$ are the momenta for outgoing partons, $\omega = E - E'$ is the energy transfer, $M(s,t)$ is the matrix element for $Qp \to Qp$ scattering ($s = (P+K)^2$, $t = (P-P')^2$ are the Mandelstam variables), $n_q(k) = (e^{k/T}+1)^{-1}$ and $n_g(k) = (e^{k/T}-1)^{-1}$ are the Fermi and Bose distributions, $\epsilon_q = -1$, $\epsilon_g = 1$, $g_q = 4N_cN_f$, $g_g = 2(N_c^2-1)$. In (7) $\omega_{\max}$ is the upper limit of the energy loss. Similarly to the radiative energy loss we take $\omega_{\max} = E/2$. After integrating over the $\mathbf{p}'$, azimuthal angle of the transverse momentum $\mathbf{k}_\perp$ and k'_z, (7) takes the form

$$\frac{dE}{dz} = \frac{1}{16EE'v(2\pi)^4} \sum_{p=q,g} g_p \int \frac{dk_z dk_\perp k_\perp\, n_p(k)}{k} \times \\ \times \int \frac{d\mathbf{q}_\perp [1+\epsilon_p n_p(k')]}{k'} \frac{\omega}{J} \langle |M(s,t)|^2 \rangle \theta(\omega_{\max} - \omega), \quad (8)$$

$$J = \left| \frac{\partial}{\partial q_z}(E'+k') \right| = \left| \frac{k'_z}{k'} + \frac{q_z - P_z}{\sqrt{m_Q^2 + (P_z - q_z)^2 + \mathbf{q}_\perp^2}} \right|, \quad (9)$$

where $\mathbf{q} = \mathbf{k}' - \mathbf{k}$. The longitudinal component q_z is determined from the energy conservation $E + k = E' + k'$. The ω may be written as (if we take $v = 1$)

$$\omega = \frac{-t - tk_z/E + 2\mathbf{k}_\perp \mathbf{q}_\perp}{2(k - k_z)}. \quad (10)$$

In the Bjorken analysis [2] the last two terms in the numerator of (10) have been neglected. Note that for heavy quark at $E \lesssim m_q^2/T$ (in this energy region the heavy quark becomes nonrelativistic in the centre of mass system of the binary collision) the value of $\omega_{\max}$ may be smaller than $E/2$ due to the kinematical limits. In numerical calculations we have used accurate formulas for the matrix elements [28, 29].

4. To apply our formulas we need to specify the parametrization of α_s and mass parameters. We parametrize $\alpha_s(Q^2)$ by the one-loop expression which is frozen at some value α_s^{fr} at $Q \le Q_{fr}$. Such a form with $\alpha_s^{fr} \approx 0.7$ was used in the analyses of the low-x structure functions within the dipole approach [30, 31, 27]. A similar parametrization has been used in [32] in the analysis of the heavy quark jets, where the authors obtained $\int_0^{2\text{ GeV}} dQ \alpha_s(Q^2)/\pi \approx 0.36$ GeV. It gives $\alpha_s(Q < Q_{fr}) = \alpha_s^{fr} \approx 0.7$, and $Q_{fr} \approx 0.82$ GeV (for $\Lambda_{QCD} = 0.3$ GeV) which agree surprisingly well with the parameters of Refs. [30, 31]. In the vacuum the stopping of the growth of α_s at $Q \lesssim Q_{fr} \sim 1$ GeV may be caused by the nonperturbative effects [32, 33]. In the QGP thermal partons can give additional suppression of α_s at low momenta [34]. At present there is no robust information on $\alpha_s(Q^2)$ in the QGP for gluons interacting with energetic ($E \gg T$) partons which is necessary in our case. Available pQCD calculations are performed in the static limit (see, for example, [35–37] and references therein). The α_s obtained in [36, 37] has a pole at $Q/\Lambda_{QCD} \sim 3$ at $T \sim 250$ MeV. Thus, in pQCD, even for the static case, the situation with Q-dependence of the in-medium α_s is unclear. For this reason it seems reasonable to estimate α_s^{fr} from the lattice results on the thermal $\alpha_s(T)$. The lattice simulations [38] give $\alpha_s(T)$ smoothly decreasing from ~ 0.5 at $T \approx 175$ MeV to

~ 0.35 at $T \approx 400$ MeV. This is also consistent with the analysis of the lattice data within the quasiparticle model [39]. One can expect that the thermal $\alpha_s(T)$ should be somewhat smaller than α_s^{fr}. For this reason it seems reasonable to use $\alpha_s^{fr} \sim 0.5$ for RHIC and LHC conditions.

The collisional energy loss of light quarks and gluons is not sensitive to the quark and gluon quasiparticle masses. However, this is not the case for the induced gluon radiation which is especially sensitive to m_g. We use the quasiparticle masses obtained in Ref. [39] from the analysis of the lattice data within the quasiparticle model. For the relevant range of the plasma temperature $T \sim (1-3)T_c$ the analysis [39] gives $m_q \approx 0.3$ and $m_g \approx 0.4$ GeV. We perform calculations for a fixed and T-dependent Debye mass. In the first case we use μ_D obtained in [39] through the perturbative relation $\mu_D = \sqrt{2}m_g$ with m_g extracted from the quasiparticle fit of the lattice data, which gives approximately the T-independent value $\mu_D \approx 0.57$ GeV. For the T-dependent parametrization we take the Debye mass obtained in the lattice calculations for $N_f = 2$ [40] which give the ratio μ_D/T slowly decreasing with T ($\mu_D/T \approx 3$ at $T \sim 1.5T_c$, $\mu_D/T \approx 2.4$ at $T \sim 4T_c$).

We use the Bjorken model [41] with the longitudinal expansion of the QGP produced in AA-collisions which gives the proper time dependence of the plasma temperature $T^3\tau = T_0^3\tau_0$ (T_0 is the initial plasma temperature). Thus, we have the number density $n(z) \propto 1/z$ for $z > \tau_0$. In the mixed phase the fraction of the QGP was calculated according to the $1/\tau$ dependence of the entropy density. Rescatterings in the hadron phase giving a small contribution have been neglected. As in our earlier analysis [11] we assume that the QGP is in the thermal and chemical equilibrium (we take $N_f = 2.5$). For RHIC we performed calculations for $T_0 = 297$ MeV and $\tau_0 = 0.5$ fm which correspond to the initial conditions in successful hydrodynamic description of $Au+Au$ collisions at $\sqrt{s} = 200$ GeV [42], and agree with the total entropy extracted from the charged particle multiplicity. For $Pb+Pb$ collisions at LHC for $\sqrt{s} = 5500$ GeV we use $T_0 = 350$ MeV (with the same $\tau_0 = 0.5$ fm), which was obtained from the extrapolation of the RHIC data on the charged particle distribution to the LHC energy performed in [43].

5. We present the numerical results for $\alpha_s^{fr} = 0.5$. Figs.1, 2 show the radiative and collisional energy losses for the light quark and gluon for RHIC and LHC conditions. To illustrate the effect of the running coupling constant we present the results for the case $\alpha_s(Q) = \alpha_s^{fr}$ as well. The higher panels correspond to the T-independent $\mu_D = 0.57$ GeV [39], and the lower panels to the T-dependent Debye mass from Ref. [40]. The calculations were performed for $L = 5$ fm which is the typical parton pathlength in the QGP (and mixed) phase with life-time about $R_A \sim 6$ fm for the central heavy ion collisions. One sees that the fraction of the collisional energy loss is relatively small. At $E \sim 10$ GeV the ratio $\Delta E_{\mathrm{col}}/\Delta E_{\mathrm{rad}}$ is about 0.3-0.4 for quarks and 0.2-0.3 for gluons. The smaller fraction of the collisional energy loss for gluons is due to the inequality $\sigma_{gq\bar{q}} < \sigma_{ggg}$ which increases the radiative energy loss for gluons (in addition to the factor $C_A/C_F = 9/4$ coming from the radiation vertex). For this reason effect of the collisional energy loss should be weaker in the kinematical region where the hadron spectra are dominated by gluon jets. The T-dependent parametrization of the Debye mass gives somewhat smaller ΔE. It is due to suppression of the rescatterings in the initial high-temperature region with $z \lesssim 1$ fm where μ_D may be about 1 GeV. However, the sensitivity of the results to the Debye mass is relatively weak. Figs.1, 2 demonstrate that the running α_s leads to flattening the radiative and collisional energy losses at high energies.

Fig.3 shows the results for the charm ($m_c = 1.2$ GeV) and bottom ($m_b = 4.5$ GeV) quarks obtained with running α_s for the Debye mass from [40]. The difference in the radiative energy loss for light and charm quarks is small. The charm radiative energy loss is only suppressed by $\sim 10\%$ at $E \sim 10$ GeV compared to the light quarks. For the bottom quark the mass suppression at low energies is significant. At high energies the radiative energy loss for the bottom quark becomes larger than that for the charm quark, and the charm contribution exceeds slightly the light quark one. This fact is connected with emission of gluons at moderate values of x where the induced radiation for heavy quarks turns out to be enhanced in the so-called diffusion regime when $L \ll L_f$ [44].

A small difference in the energy loss for the light and charm quarks show that one can expect approximately the same nuclear suppression for light and c-quarks. This is in contradiction with the considerable suppression of the induced gluon radiation from charm predicted in [45]. However, in [45] there was not performed any accurate evaluation of the mass effects in the induced radiation. In [45] the mass suppression factor for charm was defined as the ratio of the vacuum gluon emission spectra for heavy and light quarks which evidently has noting to do with the mass modification of the induced gluon radiation.

To study the infrared sensitivity of our results we also performed computations for $m_g = 0.2$ and $m_g = 0.6$ GeV. These values of m_g give reasonable the lower and

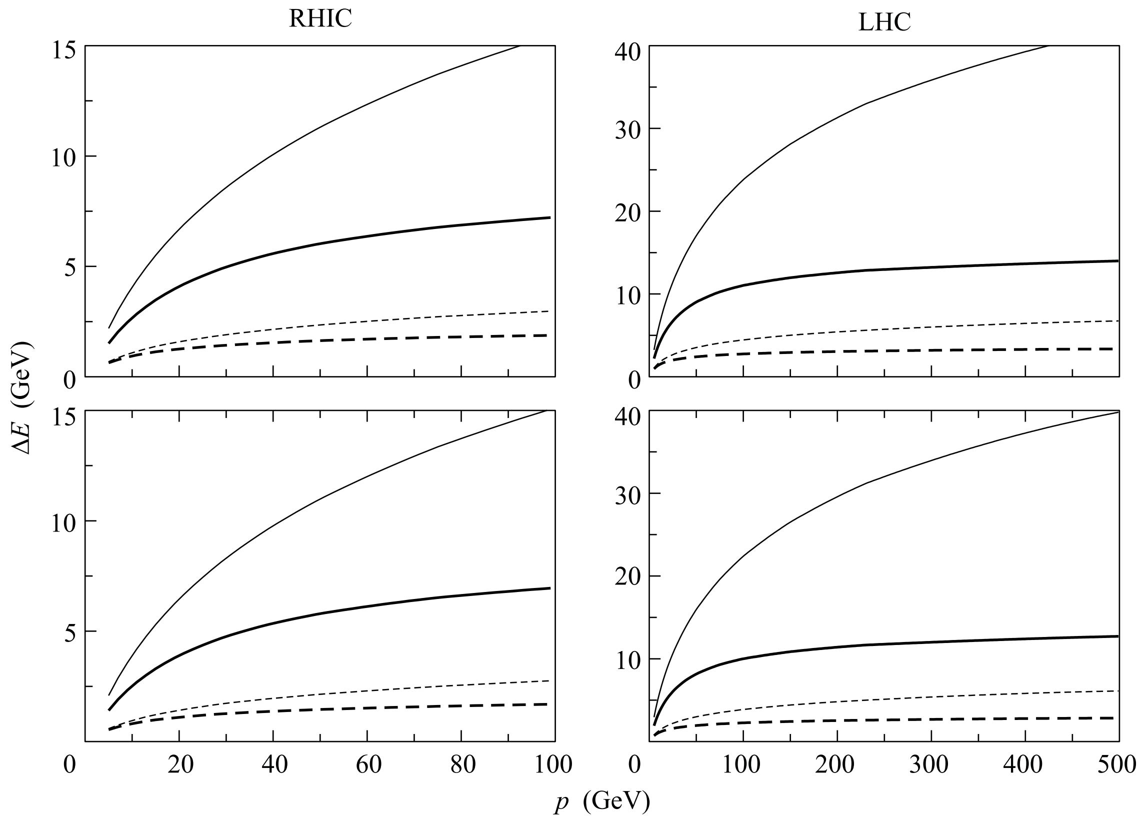

Fig.1. The light quark radiative (solid line) and collisional (dashed line) energy loss for RHIC (left) at $\sqrt{s} = 200$ GeV and LHC (right) at $\sqrt{s} = 5.5$ TeV conditions for $L = 5$ fm. The thick curves correspond to the running α_s, and thin curves to $\alpha_s = 0.5$. The higher panels show the results for the T-independent Debye mass $\mu_D \approx 0.57$ GeV [39], and the lower panels for the T-dependent Debye mass from the lattice calculations [40]

upper limits of the infrared cutoff for the induced gluon emission for RHIC and LHC conditions[1]. At $E \sim 10$ GeV for $m_g = 0.2$ GeV ΔE_{rad} is bigger by $\sim 20-30\%$ than that for $m_g = 0.4$ GeV, for $m_g = 0.6$ GeV the effect is of opposite sign. The effect of variation of the gluon mass in the above interval of m_g becomes small at higher energies ($\lesssim 10-15\%$ at $E \gtrsim 40$ GeV). To study the effect of variation of α_s we have also performed calculations using for α_s^{fr} the values 0.7 and 0.35. The first one neglects the in-medium suppression of coupling constant, and the second one, in the light of the lattice results, can be viewed as an low bound for α_s^{fr} for RHIC and LHC conditions. Numerically we obtained approximately the same fraction of the collisional energy loss as for $\alpha_s^{fr} = 0.5$.

Note that due to the dominance of the radiative energy loss modeling the jet quenching with the collisional energy loss alone within the model of a particle undergoing Brownian motion [12] does not make sense. Evidently for an accurate evaluation of the nuclear modification factor the radiative and collisional energy losses must be treated on an even footing. One can expect a nontrivial interplay of these two effects, say, it is clear that they cannot be additive, since the collisional energy loss will suppress the effective in-medium gluon formation length. The quantum nonlocal character of the induced gluon radiation makes the problem very complicated even at the level of the radiative energy loss. Presently the distribution in the induced energy loss which is necessary for evaluation of the nuclear modification factor is usually calculated assuming the independent gluon emission [47] which, however, has not any serious theoretical justification.

[1] Note that the analysis of the low-x proton structure function within the dipole BFKL equation [27, 31] gives the value of the effective gluon mass for gluon emission in the parton-nucleon interaction about 0.75 GeV. This value agrees well with the natural infrared cutoff for gluon emission in the vacuum $m_g \sim 1/R_c$, where $R_c \approx 0.27$ fm is the gluon correlation radius in the QCD vacuum [46]. One can expect that in AA-collisions the infrared cutoff will be approximately the same only for gluon emission in the developed mixed phase and for fast gluons with $L_f \gg L$ which give relatively small contribution to the total radiative energy loss.

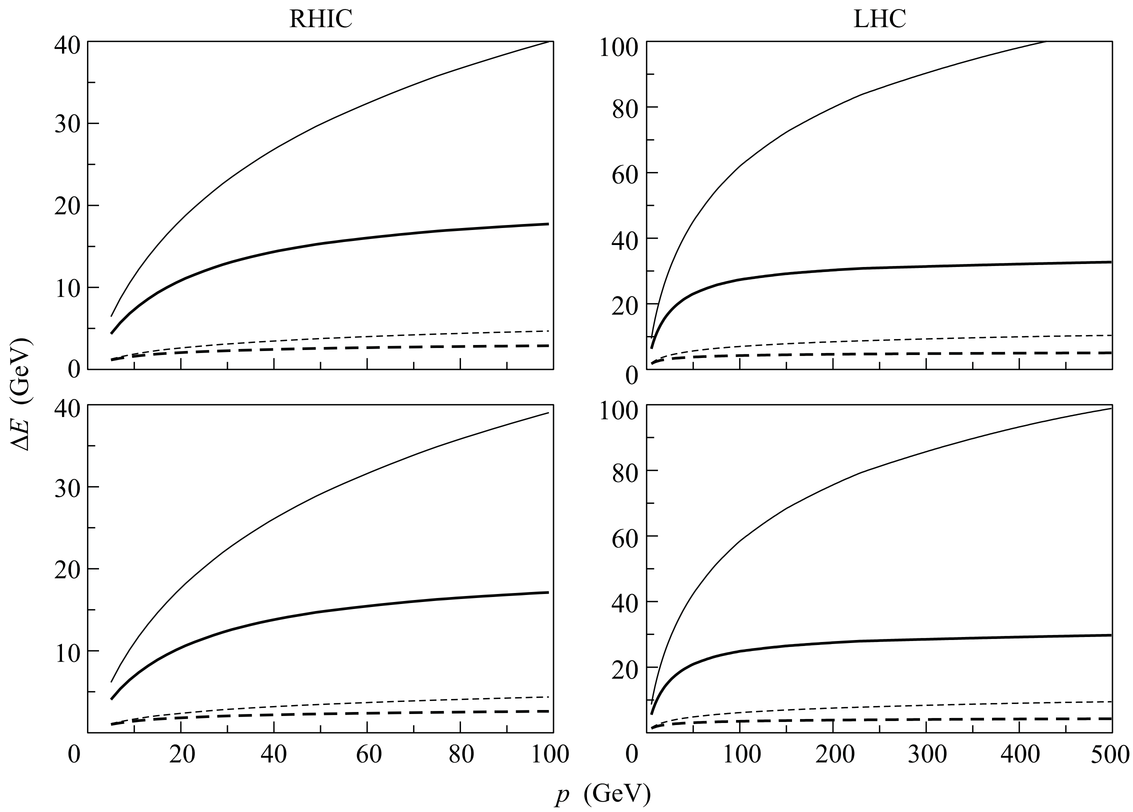

Fig.2. The same as in Fig.1 but for gluon

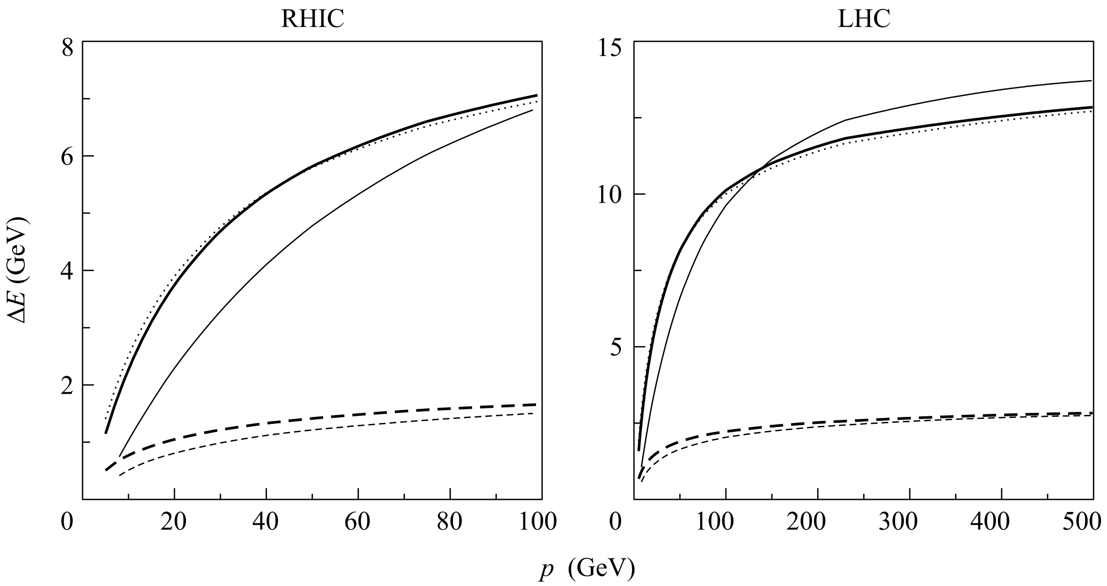

Fig.3. The charm (thick curves) and bottom (thin curves) quark radiative (solid line) and collisional (dashed line) energy loss for RHIC (left) at $\sqrt{s} = 200$ GeV and LHC (right) at $\sqrt{s} = 5.5$ TeV conditions for $L = 5$ fm, $m_c = 1.2$ GeV, $m_b = 4.5$ GeV. The dotted line shows the radiative energy loss for light quark. The calculations were performed with the running α_s and the T-dependent Debye mass from [40]

Although our approach does not treat accurately the region of soft momentum transfer $q \lesssim m_D$ we can expect that this inaccuracy should be small. Indeed, the relative contributions into collisional energy loss of the soft region with $q \lesssim 2\mu_D$ evaluated in our approach and in the HTL pQCD approach [23] with accurate treatment of the collective excitations are close. Also, even for low parton energy $E \sim 5 - 10$ GeV the soft region gives relatively small effect (about 30% for $T \sim 250$ MeV). In any case, since the infrared effects should modify the dipole cross section (5) and the probability of the collisional energy transfer approximately in the same way, one can expect a good stability of the ratio of the two mechanisms against the inaccuracy in the soft momentum region.

6. In summary, we have performed the comparison of the radiative and collisional energy losses of energetic quarks and gluons in an expanding quark-gluon plasma for RHIC and LHC energies. The radiative energy loss has been calculated within the LCPI approach [4]. To evaluate the collisional energy loss we have used the Bjorken model of elastic binary collisions with an accurate treatment of kinematics of the binary collisions. The calculations have been performed with the same infrared cutoffs and parametrization of the coupling constant for the radiative and collisional energy loss. We find that for RHIC and LHC conditions the fraction of the collisional energy loss is relatively small, and decreases with energy. For gluons it is smaller than for quarks. Our calculations show that the difference in the radiative energy loss for charm and light quarks is small. For this reason the nuclear modification factor for light hadrons and D-mesons should be approximately the same in the kinematical region where the light hadron spectra are dominated by the quark jets. At sufficiently large energies the heavy quark energy loss becomes bigger than that for light quark.

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