QUARK ENERGY LOSS IN AN EXPANDING QUARK-GLUON PLASMA

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Abstract

We study the quark energy loss in an expanding quark-gluon plasma. The expanding plasma produced in high energy AA-collision is described by Bjorken's model. The dependence of the energy loss on the infrared cutoff for the radiated gluons, on the quark mass, and on the initial conditions of QCD plasma is investigated.

Study of the induced gluon radiation from a fast quark in a quark-gluon plasma (QGP) is of great importance in connection with the forthcoming experiments on high energy AA-collisions at the RHIC and LHC. It is expected that the energy loss of high-$p_\perp$ jets produced at the initial stage of AA-collision may be an important potential probe for formation of QGP.[1,2] In the recent works[3,4] the quark energy loss, ΔE_q, was estimated for a homogeneous QGP. Here we report on the evaluation of ΔE_q in expanding QGP in the light-cone path integral approach to the induced radiation.[5] As in previous works, QGP is modelled by a system of static Debye screened scattering centers.[1]

We consider a fast quark produced in the central rapidity region with a velocity perpendicular to the axis of AA-collision. We choose the z axis along the initial quark velocity, and the quark production point is assumed to be at $z = 0$. Then, the distance passed by the quark in QGP, $L = z$, is close to the expansion time, τ. For the $\tau(z)$-dependence of the temperature of QGP we use prediction of Bjorken's model[6] $T\tau^{1/3} = T_0\tau_0^{1/3}$.

The probability of radiation of a gluon with the fractional longitudinal momentum x from a fast quark produced at $z = 0$ is given by[5]

$$\frac{dP}{dx} = 2\mathrm{Re}\int\limits_0^\infty dz_1 \int\limits_{z_1}^\infty dz_2 \exp\left[\frac{i(z_1 - z_2)}{L_f}\right] g(z_1, z_2, x)\left[K(0, z_2|0, z_1) - K_v(0, z_2|0, z_1)\right] . \quad (1)$$

Here K is the Green's function for the Hamiltonian (acting in the transverse plane)

$$H = \frac{\mathbf{q}^2}{2\mu(x)} + v(\boldsymbol{\rho}, z), \quad (2)$$

$$v(\boldsymbol{\rho}, z) = -i\frac{n(z)\sigma_3(\boldsymbol{\rho}, x, z)}{2}, \quad (3)$$

and K_v is the Green's function for the Hamiltonian (2) with $v(\boldsymbol{\rho}, z) = 0$. In (2) the Schrödinger mass is $\mu(x) = E_q x(1-x)$, $L_f = 2E_q x(1-x)/[m_q^2 x^2 + m_g^2(1-x)]$ is the gluon formation length, here m_q is the quark mass, and m_g is the mass of the radiated gluon. We introduce the gluon mass to remove the contribution of the unphysical long-wave gluon excitations. In (3) $n(z)$ is the number density of QGP, and σ_3 is the cross section of interaction of color singlet $q\bar{q}g$ system with color center. Summation over triplet (quark) and octet (gluon) color states is implicit in (3). The z-dependence of σ_3 is connected with the one of the Debye screening mass.

The vertex factor $g(z_1, z_2, x)$, entering (1), reads (we neglect the spin-flip $q \to qg$ transitions, which give a small contribution to the quark energy loss)

$$g(z_1, z_2, x) = \frac{\alpha_s[4 - 4x + 2x^2]}{3x}\,\frac{\mathbf{q}(z_2)\cdot\mathbf{q}(z_1)}{\mu^2(x)} . \quad (4)$$

The Hamiltonian (1) describes evolution of the light-cone wave function of a fictitious $q\bar{q}g$ system.[5] Using relations between the transverse Green's functions and the light-cone functions for the transition $q \to qg$ in vacuum and inside medium[7] the radiation rate (1) after some algebra can be represented in another form

$$\frac{dP}{dx} = \int\limits_0^\infty dz\, n(z)\frac{d\sigma_{eff}^{BH}(x, z)}{dx}, \quad (5)$$

$$\frac{d\sigma_{eff}^{BH}(x,z)}{dx} = \mathrm{Re} \int d\boldsymbol{\rho}\, \psi^{*}(\boldsymbol{\rho},x)\sigma_3(\rho,x,z)\psi(\boldsymbol{\rho},x,z)\,, \tag{6}$$

where $\psi(\boldsymbol{\rho},x)$ is the light-cone wave function for the $q \to qg$ transition in vacuum, and $\psi(\boldsymbol{\rho},x,z)$ is the medium-modified light-cone wave function for $q \to qg$ transition in medium at the longitudinal coordinate z. In the low density limit and $z \to \infty$ $\psi(\boldsymbol{\rho},x,z)$ becomes close to $\psi(\boldsymbol{\rho},x)$, and (6) is reduced to formula for the Bethe-Heitler cross section for a quark incident on an isolated center, $d\sigma^{BH}(x)/dx$.[8] At $z \to 0$ the ratio $d\sigma_{eff}^{BH}(x,z)/dx/d\sigma^{BH}(x)/dx$ vanishes.[4,8] This is a direct consequence of the decrease at small z of the transverse size of the qg Fock component of the quark produced at $z=0$. Note that this effect is responsible for the L^2-dependence of ΔE_q at $E_q \to \infty$ for a homogeneous medium.[3,4]

The three-body cross section entering the imaginary potential (3) can be written as[9]

$$\sigma_3(\rho,x,z) = \frac{9}{8}[\sigma_2(\rho,z) + \sigma_2((1-x)\rho,z))] - \frac{1}{8}\sigma_2(x\rho,z)\,, \tag{7}$$

where $\sigma_2(\rho,z)$ is the dipole cross section of interaction with color center of color singlet $q\bar{q}$ system. The latter can be written as $\sigma_2(\rho,z) = C_2(\rho,z)\rho^2$, where $C_2(\rho,z)$ has a smooth (logarithmic) dependence on ρ at small ρ. As in Ref. 4 we approximate $C_2(\rho,z)$ by its value at $\rho \approx 1/m_g$. Then the Hamiltonian (2) takes the oscillator form with the z-dependent frequency

$$\Omega(z) = \frac{(1-i)}{\sqrt{2}}\left(\frac{n(z)C_3(x,z)}{E_q x(1-x)}\right)^{1/2}\,, \tag{8}$$

where $C_3(x,z) = \frac{1}{8}\left\{9[1+(1-x)^2] - x^2\right\} C_2(1/m_g,z)$. The coefficient C_2 was calculated in the double gluon approximation. In numerical calculations we evaluated the Green's function K using the approach previously developed in analysis of the vector meson photoproduction.[10]

It follows from (8) that even for a thermolized QGP the oscillator frequency is dominated by the gluon contribution. There is every indication[11,12] that for the RICH and LHC the hot QCD medium produced in AA-collisions at $\tau \sim \tau_0 \sim 0.1$ fm will be thermolized gluon plasma to a first approximation, and for quarks the chemical equilibration is not reached during the expansion of QGP. For this reason we take for the gluon fugacity $\lambda_g = 1$. To take into account the suppression of quarks we use the value $\lambda_q = 1/3$ for the quark fugacity. Numerical calculations were carried out with $\alpha_s = 1/3$. For τ_0 we use the value 0.1 fm.[11,12] At $z \le \tau_0$ we take $\Omega(z) = \Omega(\tau_0)$. We have evaluated ΔE_q for $T_0 = 1100$ expected for central Pb-Pb collisions at the LHC.[11] To study the dependence of the quark energy loss on T_0 we give also the results for $T_0 = 700$ MeV. The numerical predictions for ΔE_q as a function of E_q for $L = 3$, 6, 9 fm are shown in Fig. 1. The results were obtained for $m_q = 0.2$ GeV. As for a homogeneous QGP[4], our predictions have a weak dependence on m_q. To study the infrared sensitivity of ΔE_q we present in Fig. 1 the results for $m_g = 0.75$ and 0.375 GeV. These values are of the order of the Debye screening mass for QGP in the region $\tau \gtrsim 2-3$ fm, which, as our numerical calculations show, dominates the quark energy loss. For this reason the value of m_g within the above range seems to be a plausible estimate for the infrared cutoff in the considered problem. Note

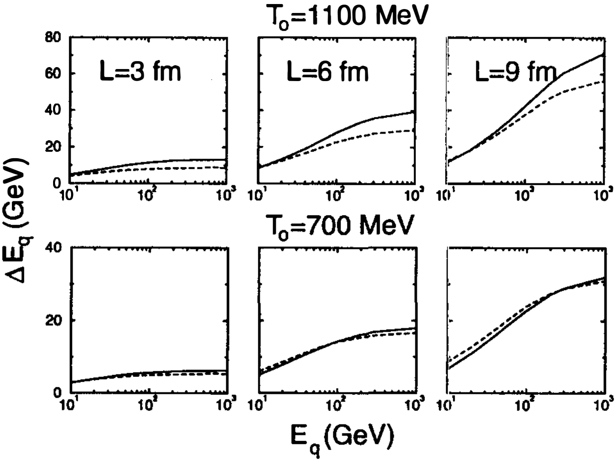

Figure 1: The quark energy loss as a function of the quark energy for m_q = 0.2 GeV, m_g = 0.75 GeV (solid line) and m_g = 0.375 GeV (dashed line).

that in the limit $E_q \to \infty$ ΔE_q has only a smooth (logarithmic) m_g-dependence connected with ρ-dependence of the coefficient C_2.[4)] Our numerical calculations really demonstrate that for $E_q \gtrsim 10$ GeV ΔE_q is not very sensitive to m_g. Fig. 1 shows that ΔE_q grows by a factor $\sim 2-3$ as E_q increases from ~ 10 to ~ 100 GeV. The predictions for $L = 6$ and $L = 9$ fm can be regarded as a plausible estimates for ΔE_q in central collisions of heavy nuclei. We have checked that the effect of the mixed phase (with $T_c \approx 150$ MeV) and the hadronic phase turns out to be relatively small. Comparison of the results for $T_0 = 1100$ and $T_0 = 700$ MeV demonstrates that the T_0-dependence of ΔE_q is not very strong. This is connected with an increase of Landau-Pomeranchuk-Migdal suppression and a decrease of the coefficient C_2 for a high density QGP.

To study the sensitivity of ΔE_q to dynamics of QGP at small τ we have also carried out the calculations taking $\Omega(z \lesssim a) = \Omega(a)$ with $a \sim 1-2$ fm. The results for these versions do not differ strongly from the ones in Fig. 1. This demonstrates that ΔE_q is insensitive to the dynamics of QGP at $\tau \lesssim 2$ fm. This fact is a consequence of the vanishing of the ratio $d\sigma_{eff}^{BH}(x,z)/dx/d\sigma^{BH}(x)/dx$ at small z.

To illustrate the quark mass dependence of ΔE_q in Fig. 2 we compare the results for ΔE_q for c-quark ($m_q = 1.5$ GeV) with the ones for light quarks ($m_q = 0.2$ GeV). As one can see the dependence of ΔE_q on m_q becomes weak at $E_q \gtrsim 50-100$ GeV.

In Fig. 3 we present the gluon spectra for $m_q = 0.2$ and 1.5 GeV at E_q =50, 100 and 200 GeV for $T_0 = 1100$ MeV. We also show in this figure the Bethe-Heitler spectra. As one can see the radiation rate is strongly suppressed by the medium and finite-size effects for the gluon momenta $k \gtrsim 10$ GeV (excluding a narrow region near $k \approx E_q$). For the heavy quark the radiation rate at $k \approx E_q$ is suppressed as compared to the light quark.

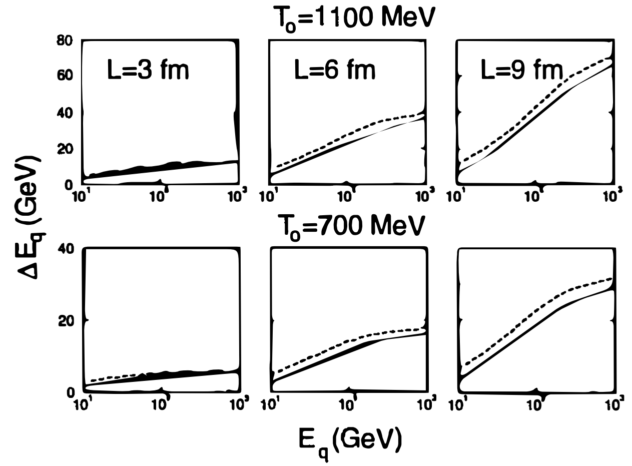

Figure 2: The quark energy loss as a function of the quark energy for m_g = 0.75 GeV, m_q = 1.5 GeV (solid line) and m_q = 0.2 GeV (dashed line).

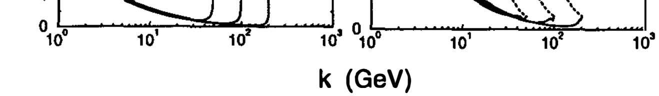

Figure 3: The gluon spectrum (solid line) as a function of the gluon momentum at E_q =50, 100 and 200 GeV for m_q = 0.2 (a) and 1.5 (b) GeV. In both the cases m_g = 0.75 GeV, and T_0 = 1100 MeV. The dashed line shows the Bethe-Heitler spectrum.

However, in the soft region the spectra for the light and heavy quarks are close to each other.

Our predictions for the absolute normalization of the energy loss must be regarded as rough estimates with uncertainties of a factor ~ 2. Nevertheless, the obtained rather large values of ΔE_q show that jet quenching may be an important probe for the formation of QGP in AA-collisions. The Monte-Carlo analysis of jet quenching[2] with ΔE_q close to the estimates of the present work demonstrates that the induced radiation must considerably modify the charged particle spectra. The induced gluon radiation can also lead to other interesting experimental consequences. For instance, fluctuations of the quark energy loss will generate additional transverse momentum (defined with respect to the AA-collision axis) in production of $c\bar{c}$ and $b\bar{b}$ pairs. In the case of $g \to gg$ transitions the production of the gg state in the color decuplet state can increase the cross section for production of baryon-antibaryon pairs through the mechanism analogous to the one previously discussed[13] in connection with $B\bar{B}$ annihilation at high energies.

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