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RAPID COMMUNICATIONS

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Energy loss of a heavy quark in the quark-gluon plasma

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The energy loss dE/dx for a heavy quark propagating through a quark-gluon plasma is calculated to leading order in the QCD coupling constant. Simple formulas for dE/dx are obtained in the regions $E \ll M_Q^2/T$ and $E \gg M_Q^2/T$, where M_Q is the mass of the heavy quark and T is the temperature. The crossover energy between the two regions is determined to be approximately $1.8M_Q^2/T$. Under conditions relevant to ultrarelativistic heavy-ion collisions, charm quarks and bottom quarks lie on opposite sides of the crossover energy and therefore experience significantly different energy losses.

Jets caused by high-energy quarks and gluons in ultrarelativistic heavy-ion collisions might provide a probe for the existence of a quark-gluon plasma (QGP). High-energy partons coming from initial hard collisions lose energy by propagating through the dense matter formed between the nuclei after collision. The energy loss is expected to be greater in AA collisions compared to pp or pA collisions, a phenomenon known as jet quenching. Gyulassy and Plümer [1] suggested that jet quenching should be suppressed if the dense matter consists of a QGP instead of hadrons. Their observation was based on an estimate of the energy loss of high-energy partons in a QGP by Bjorken [2] ($-dE/dx \sim 0.2$ GeV/fm for a 20-GeV quark at a temperature of about $T=0.25$ GeV), which is considerably smaller than in hadronic matter ($-dE/dx \sim 1$ GeV/fm). Bjorken considered the collisional energy loss of a massless quark due to elastic scattering off the quarks and gluons in the QGP. At tree level, there is a logarithmic infrared singularity in the integral over the three-momentum transfer q. Bjorken estimated the energy loss by keeping only the logarithmic term with physically reasonable upper and lower limits $q_{\min}$ and $q_{\max}$. In a complete calculation, the upper and lower cutoffs on q should be provided automatically by the physics of the energy-loss process. The purpose of this paper is to present a complete calculation of the energy loss of an energetic heavy quark to leading order in the QCD coupling constant.

Thoma and Gyulassy [3] eliminated the ambiguity due to the lower limit $q_{\min}$ of the momentum transfer by properly including the screening effects of the QGP. The logarithmic infrared divergence that arises in naive calculations is cut off at the *soft* momentum scale $g_s T$, where g_s is the QCD coupling constant. Unfortunately their calculation was incomplete in the region of *hard* momentum transfer ($q \sim T$) and therefore required the imposition of an upper limit $q_{\max}$. The resulting ambiguity was avoided in a complementary calculation carried out by Svetitsky [4] in a study of the diffusion of charm quarks in the QGP. Svetitsky's drag coefficient is directly related to the energy loss: $A(p^2)=(-dE/dx)/p$, where E and p are the energy and momentum of the heavy quark. In a straightforward tree-level calculation of the energy loss due to elastic scattering, the kinematics of the scattering process automatically sets an upper limit on the momentum transfer. To cut off the infrared divergence in the tree-level calculation, Svetitsky used the *ad hoc* prescription of introducing a gluon mass which violates gauge invariance and introduces an ambiguity equivalent to Bjorken's choice of $q_{\min}$.

To obtain the energy loss to leading order in g_s, the hard-momentum-transfer contribution can be calculated at tree level, but in the soft q region it is necessary to use a resummed perturbation expansion developed by Braaten and Pisarski [5]. The *hard thermal loop* corrections to the propagator of the exchanged gluon must be resummed in order to include the screening effects of the plasma, and this resummation reproduces the result of Thoma and Gyulassy [3] in the soft region. The contributions from the hard and soft regions must then be matched together consistently to give the complete energy loss to leading order in g_s. A general method for carrying out this matching has been developed by Braaten and Yuan [6]. One introduces an arbitrary intermediate momentum scale q^* satisfying $g_s T \ll q^* \ll T$, which is always possible in the weak-coupling limit $g_s \to 0$. For moderate values of g_s, this should be interpreted as merely a mathematical device for isolating all terms of leading order in g_s. The contribution from hard $q > q^*$ is calculated using tree-level Feynman diagrams while ignoring any screening due to the plasma. The logarithmic infrared divergences of the tree-level calculation manifest themselves as logarithms of q^*. The contribution from soft momentum transfers $q < q^*$ is calculated using the resummed perturbation expansion to take into account the effects of screening, and it also depends logarithmically on q^*. Adding the hard and soft contributions, the dependence on the arbitrary scale q^* cancels.

In Ref. [7], a quantum-field-theoretic formulation of the energy loss was developed. The method of Braaten and Yuan [6] was then used to compute the energy loss of a muon propagating through a plasma of electrons, positrons, and photons to leading order in the QED coupling constant e and in T/M, where M is the mass of the muon. In this paper, that calculation will be extended to obtain the energy loss of a relativistic heavy quark propagating through a QGP to leading order in g_s and in T/M_Q, where M_Q is the mass of the heavy quark. We assume that the mass M_Q and the momentum p of the heavy quark are both much larger than the temperature T of the plasma. Simple results can be obtained in two energy regimes, $E \ll M_Q^2/T$ and $E \gg M_Q^2/T$. The maximum momentum transfer q from elastic scattering off a thermal quark or gluon with energy k is $q_{\max} = 2k(1+k/E)/(1-v+2k/E)$, where v is the velocity of the heavy quark. In the region $E \ll M_Q^2/T$, we can set $q_{\max} = 2k/(1-v)$, while for $E \gg M_Q^2/T$, the maximum momentum transfer is the energy of the heavy quark: $q_{\max} = E$. In the region $E \sim M_Q^2/T$, the formula for dE/dx is very complicated because the general expression for $q_{\max}$ must be used. A good approximation to dE/dx in this crossover region can be obtained by switching from the formula for $E \ll M_Q^2/T$ to the formula for the ultrarelativistic region $E \gg M_Q^2/T$ at a crossover energy $E_{\rm cross}$ which is determined by continuity.

At leading order in g_s, the energy loss of a heavy quark comes from elastic scattering from thermal quarks (Qq scattering) and thermal gluons (Qg scattering). We first consider the energy loss in the region $E \ll M_Q^2/T$. We assume that the heavy quark has a kinetic energy much greater than T. Some of the contributions to dE/dx can be obtained from the corresponding QED calculation in Ref. [7] by simple substitution. The soft contribution from Qg and Qq scattering is obtained from the QED case by replacing e by the QCD coupling constant g_s, multiplying by a color factor $\frac{4}{3}$, and replacing the thermal photon mass $m_\gamma = eT/3$ by the thermal gluon mass [8] $m_g = (g_s T/\sqrt{3})(1+n_f/6)^{1/2}$, where n_f is the number of active flavors in the QGP. The result is

$$\left(-\frac{dE}{dx}\right)_{\rm soft}^{Qg+Qq} = \frac{g_s^4 T^2}{6\pi}\left(1+\frac{n_f}{6}\right)\Bigg\{\left(\frac{1}{v}-\frac{1-v^2}{2v^2}\ln\frac{1+v}{1-v}\right)\ln\frac{q^*}{m_g} - \frac{1}{v^2}\int_0^v dx\, x^2\left[\ln\frac{3\pi x}{2}+\frac{1}{2}\ln[1+Q_l(x)^2]+Q_l(x)\left(\frac{\pi}{2}-\arctan Q_l(x)\right)\right] - \frac{1}{2v^2}\int_0^v dx\, x^2\frac{v^2-x^2}{1-x^2}\left[\ln\frac{3\pi x}{4}+\frac{1}{2}\ln[1+Q_t(x)^2] + Q_t(x)\left(\frac{\pi}{2}-\arctan Q_t(x)\right)\right]\Bigg\}, \quad (1)$$

where $Q_l(x) = \{-\ln[(1+x)/(1-x)]+2/x\}/\pi$ and $Q_t(x) = \{\ln[(1+x)/(1-x)]+2x/(1-x^2)\}/\pi$. The hard contribution to dE/dx from Qq scattering is obtained from the QED calculation in Ref. [7] by replacing e by g_s, multiplying by a color factor $\frac{2}{3}$, and summing over the n_f flavors of the initial thermal quark:

$$\left(-\frac{dE}{dx}\right)_{\rm hard}^{Qq} = \frac{g_s^4 T^2}{6\pi}\frac{n_f}{6}\Bigg\{\left(\frac{1}{v}-\frac{1-v^2}{2v^2}\ln\frac{1+v}{1-v}\right)\left(\ln\frac{4TE}{q^* M_Q}+\frac{3}{2}-\gamma+\frac{\zeta'(2)}{\zeta(2)}\right) - \frac{1-v^2}{4v^2}\left[\mathrm{Sp}\left(\frac{1+v}{2}\right)-\mathrm{Sp}\left(\frac{1-v}{2}\right)+\frac{1}{2}\ln\frac{1+v}{1-v}\ln\frac{1-v^2}{4}\right]-\frac{2}{3}v\Bigg\}, \quad (2)$$

where $\mathrm{Sp}(x) = -\int_0^x dt\,(1/t)\ln(1-t)$ is the Spence function, $\gamma = 0.57722$ is Euler's constant, and $\zeta(z)$ is the Riemann zeta function: $\zeta'(2)/\zeta(2) = -0.56996$. Note that the logarithm of q^* in (2) cancels against the $\ln(q^*)$ term proportional to n_f in (1).

The only new calculation that is required to obtain dE/dx for a heavy quark is the hard contribution from Qg scattering. The rate of energy loss $-dE/dx$ is obtained by computing the interaction rate of the heavy quark weighted by the factor $(E-E')/v$, where v is the velocity of the heavy quark and E' is its energy after the collision. The hard contribution can be isolated by imposing a constraint $q>q^*$ on the magnitude of the three-momentum transfer $\mathbf{q}=\mathbf{k}'-\mathbf{k}$. The hard contribution from Qg scattering is

$$\left[-\frac{dE}{dx}\right]^{Qg}=\frac{1}{2E}\int\frac{d^3p'}{(2\pi)^3 2E'}\int\frac{d^3k}{(2\pi)^3 2k}n_B(k)\int\frac{d^3k'}{(2\pi)^3 2k'}[1+n_B(k')]\times(2\pi)^4\delta^4(P+K-P'-K')16\langle|\mathcal{M}|^2\rangle\frac{E-E'}{v}\theta(q-q^*)\,, \tag{3}$$

where $P=(E,\mathbf{p})$ and $K=(k,\mathbf{k})$ are the four-momenta of the incoming quark and gluon and P' and K' are the momenta of the outgoing quark and gluon. The phase space is weighted by a Bose distribution $n_B(k)=(e^{k/T}-1)^{-1}$ for the incoming gluon and a Bose-enhancement factor $1+n_B(k')$ for the outgoing gluon. The factor 16 is the number of spin and color states of the thermal gluon, and $\langle|\mathcal{M}|^2\rangle$ is the square of the matrix element, averaged over initial and summed over final spins and colors.

The matrix element $\mathcal{M}=\mathcal{M}_s+\mathcal{M}_t+\mathcal{M}_u$ for Qg scattering is given by the sum of the three Feynman diagrams in Fig. 1. There are s- and u-channel diagrams that correspond to Compton scattering in QED and a t-channel diagram involving the three-gluon vertex. In QED, Compton scattering does not contribute to dE/dx for a heavy lepton at leading order in T/M due to a cancellation between the s- and u-channel diagrams. In QCD, the cancellation is upset by the non-Abelian coupling of quarks to gluons. The square of the amplitude $\mathcal{M}_s+\mathcal{M}_u$ reduces in the limit $k,k'\ll E$ to

$$\langle|\mathcal{M}_s+\mathcal{M}_u|^2\rangle=g_s^4\frac{M_Q^4}{(P\cdot K)(P\cdot K')}\,. \tag{4}$$

The interference term $\mathrm{Re}\langle\mathcal{M}_t(\mathcal{M}_s+\mathcal{M}_u)^*\rangle$ between the t-channel and the s- and u-channel amplitudes vanishes in this limit. In the Feynman gauge, the square of the t-channel amplitude reduces in the limit $k,k'\ll E$ to

$$\langle|\mathcal{M}_t|^2\rangle=8g_s^4\frac{(P\cdot K)(P\cdot K')-M_Q^2(K\cdot K')}{Q^4}\,, \tag{5}$$

where $Q=K'-K$ is the four-momentum of the exchanged gluon. One must take care in computing (5) to sum only over the physical polarization states of the gluons. If a covariant expression is used for the polarization sum, then a contribution to (5) from quark-ghost scattering must be included in order to cancel that of unphysical gluon polarization states [9].

The squared matrix elements (4) and (5) must be inserted into (3). The integrals can be evaluated analytically using the methods of Ref. [7]. Scattering in the s and u channels does not give rise to an infrared divergence from the small-q region, so the infrared cutoff q^* can be set to zero in this calculation. The result is

$$\left[-\frac{dE}{dx}\right]^{Qg(s+u)}_{\text{hard}}=\frac{g_s^4T^2}{12\pi}\left[\frac{1}{v}-\frac{1-v^2}{2v^2}\ln\frac{1+v}{1-v}\right]. \tag{6}$$

There is an infrared divergence from the t-channel diagram, so the cutoff $q>q^*$ is necessary. Using $q^*\ll T$, the result is

$$\left[-\frac{dE}{dx}\right]^{Qg(t)}_{\text{hard}}=\frac{g_s^4T^2}{6\pi}\left\{\left[\frac{1}{v}-\frac{1-v^2}{2v^2}\ln\frac{1+v}{1-v}\right]\left[\ln\frac{2TE}{q^*M_Q}+1-\gamma+\frac{\zeta'(2)}{\zeta(2)}\right]-\frac{1-v^2}{4v^2}\left[\mathrm{Sp}\left[\frac{1+v}{2}\right]-\mathrm{Sp}\left[\frac{1-v}{2}\right]+\frac{1}{2}\ln\frac{1+v}{1-v}\ln\frac{1-v^2}{4}\right]-\frac{2}{3}v\right\}. \tag{7}$$

The complete result for dE/dx for a relativistic heavy quark with energy $E\ll M_Q^2/T$ is obtained by adding the soft contribution (1) and the hard contributions (2), (6), and (7). It can be written in the form

$$-\frac{dE}{dx}=\frac{8\pi\alpha_s^2T^2}{3}\left[1+\frac{n_f}{6}\right]\left[\frac{1}{v}-\frac{1-v^2}{2v^2}\ln\frac{1+v}{1-v}\right]\times\ln\left[2^{n_f/(6+n_f)}B(v)\frac{ET}{m_gM_Q}\right], \tag{8}$$

where $B(v)$ is a smooth function of the velocity that increases monotonically from $B(0)=0.604$ at $v=0$ to a maximum of 0.731 at $v=0.88$, and then decreases to $B(1)=0.629$ at $v=1$.

We next consider scattering in the ultrarelativistic region $E\gg M_Q^2/T$. Here again several of the contributions can be extracted from the QED calculation published in Ref. [7]. The soft contribution from Qg and Qq scattering is

$$\left[-\frac{dE}{dx}\right]^{Qg+Qq}_{\text{soft}}=\frac{g_s^4T^2}{6\pi}\left[1+\frac{n_f}{6}\right]\left[\ln\frac{q^*}{m_g}-0.843\right]. \tag{9}$$

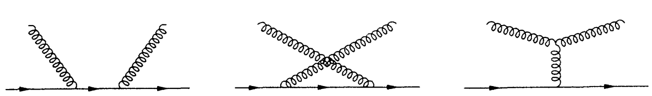

FIG. 1. Tree-level Feynman diagrams for Qg scattering (s channel, u channel, and t channel).

The hard contribution from Qq scattering is

$$\left[-\frac{dE}{dx}\right]_{\text{hard}}^{Qq}=\frac{g_s^4T^2}{12\pi}\frac{n_f}{6}\left[\ln\frac{2TE}{(q^*)^2}+\frac{8}{3}-\gamma+\frac{\zeta'(2)}{\zeta(2)}\right]. \tag{10}$$

The hard contribution from Qg scattering requires a new calculation. The s- and u-channel contributions vanish, and the t channel gives

$$\left[-\frac{dE}{dx}\right]_{\text{hard}}^{Qg}=\frac{g_s^4T^2}{12\pi}\left[\ln\frac{TE}{(q^*)^2}+\frac{8}{3}-\gamma+\frac{\zeta'(2)}{\zeta(2)}\right]. \tag{11}$$

Adding up the contributions to dE/dx in (9)–(11), the total energy loss for energies $E\gg M_Q^2/T$ is

$$-\frac{dE}{dx}=\frac{8\pi\alpha_s^2T^2}{3}\left[1+\frac{n_f}{6}\right]\ln\left[2^{n_f/2(6+n_f)}0.920\frac{\sqrt{ET}}{m_g}\right]. \tag{12}$$

In the region $E\sim M_Q^2/T$, the formula for the energy loss must cross over from the $v\rightarrow 1$ limit of (8) to (12). It should be a good approximation to simply use (8) up to some crossover energy E_{cross} and then switch to (12). By demanding that dE/dx remain continuous at $E=E_{\text{cross}}$, we determine the crossover energy to be $E_{\text{cross}}=1.80\times M_Q^2/T$ for $n_f=2$ active flavors of quarks.

The energy-loss formulas (8) and (12) are illustrated by the solid lines in Figs. 2 and 3 for conditions of relevance to ultrarelativistic heavy-ion collisions. We take the temperature of the plasma to be $T=250$ MeV, and we take the strong coupling constant at that temperature to be $\alpha_s=0.2$, which is in accordance with some lattice QCD calculations [10]. In Fig. 2, we show the energy loss of the charm quark, assuming a mass $M_c=1.5$ GeV. The crossover energy is $E_{\text{cross}}=16$ GeV, so that most of the range of energy relevant for jets in heavy-ion collisions is covered by the ultrarelativistic formula (12). The discontinuity in slope at the crossover point could be avoided by a more complete calculation for the region $E\sim M_Q^2/T$. In Fig. 3, we show the energy loss of a bottom quark with mass $M_b=5.0$ GeV. The crossover energy is $E_{\text{cross}}=180$ GeV, so the entire range of energy is covered by the formula (8). Note that the energy loss for a bottom quark in a QGP has a different momentum dependence from the charm quark, and is significantly smaller in magnitude. For example, for a momentum of 20 GeV, we find $-dE/dx=0.3$ GeV/fm for a charm quark and $-dE/dx=0.15$ GeV/fm for a bottom quark, both of which are considerably smaller than the energy loss in hadronic matter of about 1 GeV/fm. Bottom-quark jets should therefore suffer significantly less quenching than charm-quark jets if a QGP is created in a heavy-ion collision.

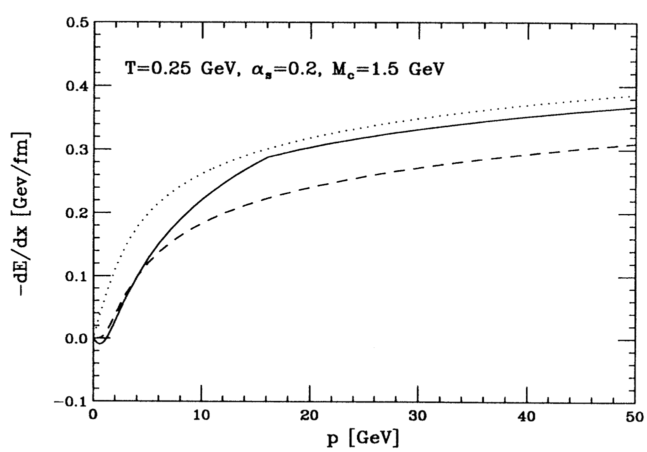

FIG. 2. Energy loss dE/dx of a charm quark as a function of its momentum for $T=250$ MeV and $\alpha_s=0.2$. The complete result to leading order in g_s (solid curve) is compared to previous calculations by Thoma and Gyulassy (dashed curve) and Bjorken (dotted curve).

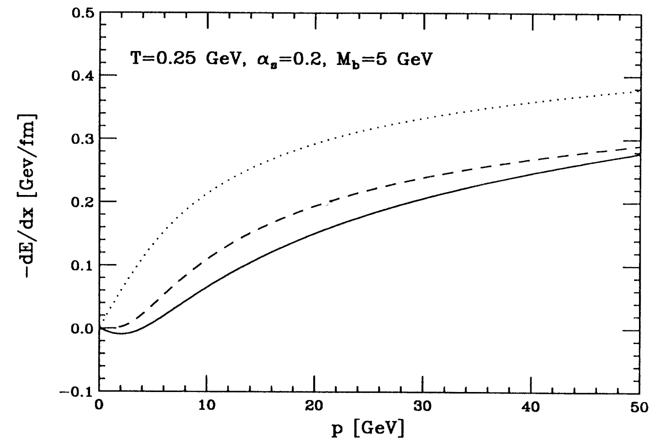

FIG. 3. Energy loss dE/dx of a bottom quark as a function of its momentum for $T=250$ MeV and $\alpha_s=0.2$. The complete result to leading order in g_s (solid curve) is compared to previous calculations by Thoma and Gyulassy (dashed curve) and Bjorken (dotted curve).

Also shown for comparison in Figs. 2 and 3 are the results of previous calculations of the energy loss. The dotted curves are the estimates of Bjorken [2] for light quarks adapted to the case of heavy quarks:

$$-\frac{dE}{dx}=\frac{8\pi\alpha_s^2T^2}{3}\left[1+\frac{n_f}{6}\right]\left[\frac{1}{v}-\frac{1-v^2}{2v^2}\ln\frac{1+v}{1-v}\right]\times\ln\frac{q_{\max}}{q_{\min}}. \tag{13}$$

For the upper limit on the momentum transfer, we follow Bjorken in using $q_{\max}=\sqrt{4TE}$, while for the lower limit we use the Debye screening mass $q_{\min}=\sqrt{3}m_g$. The dashed curves in Figs. 2 and 3 are the calculations of Thoma and Gyulassy [3]. For the charm quark in Fig. 2, the complete leading-order calculation is in reasonable agreement with that of Thoma and Gyulassy at low energies, but then crosses over to a form that is closer to Bjorken's estimate (13) at high energies.

The formula (8) for the energy loss breaks down at thermal energies $v\sim(T/M_Q)^{1/2}$. The energy loss must change sign in this region, because a quark with $v=0$ can only gain energy in a collision. The methods used above to compute dE/dx at high energies can also be used to compute it in the limit $v\rightarrow 0$. For weak coupling g_s,

$-dE/dx$ is negative corresponding to energy gain and it has a kinematic $1/v$ divergence:

$$-\frac{dE}{dx}=-\frac{16\pi\alpha_s^2T^3}{3M_Qv}\left[1+\frac{n_f}{6}\right]\ln\left[2^{n_f/(6+n_f)}0.604\frac{T}{m_g}\right]. \tag{14}$$

Note that the coefficient inside the logarithm agrees with the one in (8) in the limit $v\rightarrow 0$. A comparison of (14) with the $v\rightarrow 0$ limit of (8) can be used to make a semiquantitative estimate of the velocity at which the energy-loss changes sign $v\simeq(3T/M_Q)^{1/2}$. With a plasma temperature of 250 MeV as used in Figs. 2 and 3, this corresponds to a momentum of 1.5 GeV for charm quarks and 2.1 GeV for bottom quarks.

The formula (8) for dE/dx also breaks down at low energies for another reason. For coupling constants $g_s > 1.08$, the argument of the logarithm in (8) is less than 1 for $E<0.93g_sM_Q$ and the energy loss is negative, implying a gain of energy. If M_Q is sufficiently large, the energy below which (8) turns negative is much greater than the thermal energy. This unphysical behavior is due to a failure of the extrapolation from the weak-coupling limit $g_s\rightarrow 0$ to the physical value of g_s. This failure is not of great practical importance for the properties of heavy-quark jets in ultrarelativistic heavy-ion collisions. For the charm quark as shown in Fig. 2, (8) turns negative at a momentum of about 1.2 GeV, while for the bottom quark as illustrated in Fig. 3, it turns negative around 4.1 GeV. These are both lower than the energies of experimental interest. Nevertheless, this result serves as a warning that for properties of the QGP that have a logarithmic sensitivity to different energy scales, calculations to leading order in g_s may be meaningless. The usefulness of such a calculation can be determined only after computing the multiplicative constant inside the logarithm. In particular, any calculation that keeps only the leading-order $\ln(1/g_s)$ term is meaningless until the constant is also calculated.

The negative values of $-dE/dx$ predicted by (8) arise from isolating the terms of leading order in g_s, without any contamination from higher orders in g_s. This difficulty can be circumvented by using an effective propagator for soft gluons not only at soft momentum transfers q as in Ref. [7], but for all q. In the hard q region, this corresponds to including a subset of corrections that are higher order in g_s. Consistency and gauge invariance then demand that the effective vertices of Braaten and Pisarski [5] be used for the coupling of the virtual gluon to the thermal quarks and gluons and that the thermal quarks and gluons be placed on the mass shells of their respective effective propagators. This would of course greatly increase the complexity of the calculation of dE/dx.

We discuss briefly the energy loss of a high-energy light quark in the QGP. The energy loss of an ultrarelativistic heavy quark in (12) was calculated under the assumption $T\ll M_Q\ll E$. To obtain the energy loss of a light quark to leading order in g_s, the calculation should be repeated under the assumption $m_q\ll T\ll E$. A recent calculation by Mrowczynski [11] improves on previous calculations but is incomplete. Mrowczynski recognized the need to include contributions from both hard and soft momentum transfer. His separation of the integral over the transverse component q_T of the momentum transfer into two regions $q_T<k_0$ and $q_T>k_0$ is similar to our separation of q into soft $(q<q^*)$ and hard $(q>q^*)$ regions. In contrast with our calculation of dE/dx for a heavy quark where the dependence on q^* cancels, Mrowczynski's final result depends on k_0. He makes the arbitrary choice of setting k_0 equal to the Debye screening mass $\sqrt{3}m_g$. The failure of the k_0 dependence to cancel in the final answer is a symptom of an incomplete calculation. Mrowczynski's calculation also suffers an additional ambiguity in the choice of the maximum energy transfer, which should be imposed automatically by the physics in a complete calculation. Very large energy transfers arise from the energy of the quark being transferred to a collinear gluon, in which case the energy is not dispersed into the plasma. For light particles, it is most appropriate to think in terms of the energy of the jet rather than that of individual particles. Energy transfer to a collinear particle should then be interpreted not as energy loss, but instead as evolution of the jet due to its interactions with the plasma. A complete calculation of dE/dx for the jet will necessarily depend on the definition of a jet.

In this paper, we have only considered the collisional energy loss due to elastic Qq and Qg scattering. In addition, there is radiative energy loss due to bremsstrahlung processes, such as Qq scattering into Qqg with exchange of a virtual gluon. While a naive tree level calculation of the radiative energy loss gives a result that is higher order by a factor of g_s^2, it suffers from a quadratic infrared divergence in the integral over the momentum transferred through the virtual gluon. When the screening of the plasma at the scale g_sT is taken into account by resumming hard thermal loop corrections to the gluon propagator, the quadratic divergence is replaced by a factor of $1/g_s^2T^2$. Thus the radiative process contributes to the energy loss at the same order in g_s as elastic scattering. Unfortunately, as is the case with damping rate for a high-energy particle [3,12], the screening of the plasma at the scale g_sT only softens the quadratic infrared divergence of the tree level calculation into a logarithmic one. Because of the lack of screening of the static magnetic interaction at the scale g_sT, the radiative energy loss has a logarithmic infrared sensitivity to the smaller momentum scale g_s^2T. A complete leading order calculation of the radiative energy loss must therefore await the development of more powerful resummation techniques than those needed to calculate the energy loss from elastic scattering.

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