

Jet quenching in dense matter

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The quenching of hard jets in ultrarelativistic nuclear collisions is estimated emphasizing its sensitivity to possible changes in the energy loss mechanism in a quark-gluon plasma.

We calculate the attenuation of hard jet production in ultra-relativistic nuclear collisions and consider its sensitivity to changes in the energy loss mechanisms, dE/dx , in dense matter. In particular, we consider the consequences of a possible sudden *decrease* of dE/dx near the quark-gluon plasma phase transition temperature. We find that such a change decreases the rate of quenching with increasing A and may under favorable conditions lead to one of the signatures of that transition. In any case, hard jets provide a powerful "external" probe of the transient dense matter produced in such reactions because their production rates are calculable with perturbative QCD [1] up to a slowly varying correction factor $K \sim 2$. The main difficulty is that those rates are very small and large soft and mini-jet background fluctuations [2,3] may complicate their detection.

Previous jet calculations [4-6] for nuclear collisions considered enhanced acoplanarity of jets as a probe of multiple scattering in dense matter. Unfortunately, as emphasized in refs. [5,6], increased acoplanarity is expected to occur in both confined and deconfined phases of dense matter. A reduction of dE/dx with increasing density would, on the other hand, be a novel effect that could only occur if there were a dramatic change in the nature of quark interactions at very high densities. While we cannot prove that such a decrease is necessarily associated with the

QCD deconfinement transition, we motivate that possibility by reviewing recent lattice "data" and by showing that radiative energy loss is suppressed for high energy jets. The jet attenuation factor is then calculated varying generously dE/dx and taking into account large uncertainties in the time-evolution of ultrarelativistic nuclear collisions.

The sensitivity of jet cross sections to dE/dx in a quark-gluon plasma was first pointed out by Bjorken [7]. The collisional energy loss of a quark of energy E due to elastic interactions with partons in an ideal quark-gluon plasma at a temperature T was calculated to be

$$\left(\frac{dE}{dx}\right)_c \approx 6\alpha_s^2 T^2 \log\left(\frac{4ET}{M^2}\right) \exp\left(\frac{-M}{T}\right) \left(1 + \frac{M}{T}\right), \quad (1)$$

where M is an infrared cutoff on the order of the Debye mass. The energy loss for gluons is expected to be $\frac{2}{3}$ larger. The magnitude of the energy loss (1) is clearly very sensitive to the effective coupling strength, α_s , of fast quarks to low ($\sim M$) momentum transfer gluons. Near the deconfinement temperature, $T_c \sim 200$ MeV, one would expect that α_s could be quite large. However, recent QCD lattice studies [8,9] of the *static* heavy $q\bar{q}$ potential indicate that at least the coupling strength of heavy quarks may be surprisingly small just above T_c . In pure glue SU(3) on $24^3 \times (4-16)$ lattices the potential, fit [9] by a Debye screened form, $V(r) = \alpha \exp(-Mr)/r$, was found to be characterized by $M \approx 2T$ and $\alpha \approx 0.15 \pm$

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0.05 in the range $1.05 \lesssim T/T_c \lesssim 1.2$. Full QCD calculations [10] on smaller lattices $12^3 \times 4$ also indicate that $\alpha \lesssim 0.2$ may be surprisingly small. In addition present numerical evidence on the static string tension [9] below T_c indicates a possible reduction of the string tension as the critical temperature is approached from below. While these results all refer to static interactions in dense matter, they suggest the possibility that the dynamical coupling involved in eq. (1) may also be smaller than first expected in the plasma phase.

For $E \sim 20$ GeV jets in a plasma at temperature $T \sim 0.25$ GeV, a value of $\alpha_s \sim 0.2$ and $M = 0.5$ GeV, as used in ref. [7] and consistent with the lattice data, would imply that $(dE/dx)_c \sim 0.1$ GeV/fm! This energy loss is remarkably smaller than the conventional static and dynamic string tensions in the confined phase. The static string tension $\kappa_0 \approx 1$ GeV/fm, as determined by heavy quark spectroscopy, Regge trajectories, and lattice QCD. The dynamic string tension is the energy loss per unit length of a fragmenting high energy quark, $\kappa_d = 1/\tau_0 dE_\perp/dy \sim \kappa_0$ [10,11]. Here $dE_\perp/dy \sim 1$ GeV is the empirical transverse energy radiated per unit rapidity and $\tau_0 \sim 1$ fm is the characteristic proper formation time of hadrons. Deep inelastic lepton-nucleus data [12,13] have in fact revealed evidence for large $dE/dx \sim \kappa_d$ of jets in nuclear matter. The observed uniform (\sim factor 2) suppression of produced hadrons with $0.1 < x < 0.8$ from $E \sim 10$ GeV jets could not be explained as due to secondary inelastic cascading even in the limit of vanishing formation length. In ref. [12] we showed that additional energy loss $\sim \kappa_d$ was required to account for the suppression of the low x secondaries. The mechanism of jet energy loss in that case was postulated to be string flip interactions in the nucleus that splice the jet string into many short segments that fragment independently and leave the leading jet string with energy $\sim E - \kappa_0 R$ [12]. Because the energy loss of the leading string is limited to $\lesssim 10$ GeV by the finite size of nuclei, that mechanism also accounts naturally for the observed rapid reduction of quenching of higher energy $E \gtrsim 30$ GeV jets in nuclei.

The change in the energy loss mechanism at the deconfinement transition to one dominated by elastic collisions with partons could lead to a smaller energy loss than observed in nuclear matter if the effective coupling is as small as indicated by lattice studies.

Eventually at very high temperatures the collisional energy loss in the plasma certainly exceeds the string tension in the confined phase due to the T^2 growth, but the possibility remains that in the vicinity of T_c there may be a *reduction* of dE/dx . We note that this would also imply a reduced thermalization rate and increased viscosity coefficients near T_c .

We consider next the role of radiative energy loss in the plasma. Could the radiative energy loss induced by soft scattering processes be important for hard jets? We show next that the Landau-Pomeranchuk effect [14] strongly suppresses such radiation in the ultra-relativistic limit. This is most readily seen by considering the abelian problem of the radiation induced by one soft elastic rescattering following an initial hard scattering in a medium. Let the hard scattering occur at space-time point x_0 , whereby an electron at rest [$p_0 = (m, 0, 0, 0)$] and charge e is suddenly accelerated to very high energy [$p_1 \approx (\gamma m, \gamma \beta m, 0, 0)$ with $E = \gamma m \gg m$]. Assume that at position x_1 it suffers a soft elastic rescattering leading to $p_2 = p_1 + q$ with $q \approx (0, 0, \mathbf{q}_\perp)$. The invariant radiation distribution is then given by

$$\omega \frac{d^3 n}{d^3 k} = - \frac{\alpha}{4\pi^2} |(a_0 - a_1)^\mu \exp(ikx_0) + (a_1 - a_2)^\mu \exp(ikx_1)|^2, \quad (2)$$

where $a_i^\mu = p_i^\mu / k p_i$, and $|j^\mu|^2 \equiv j_\mu j^\mu < 0$ for the space-like current above. The integrated radiative energy loss can be decomposed into a hard and soft part as $\Delta E = \Delta E_0 + \Delta E_1$, where the hard radiative energy loss is given by

$$\Delta E_0 = - \frac{\alpha}{4\pi^2} \int d\omega d\Omega \omega^2 (a_0 - a_1)^2 \approx \frac{\alpha}{\pi} E \log\left(\frac{4E^2}{m^2}\right), \quad (3)$$

and the soft radiative energy loss including the interference between the soft and hard amplitudes is given by

$$\Delta E_1 = - \frac{\alpha}{4\pi^2} \int d\omega d\Omega \omega^2 (a_1 - a_2)_\mu \times [(a_1 - a_2)^\mu + 2f(k)(a_0 - a_1)^\mu]. \quad (4)$$

Note that the hard component is independent of the

medium, while the soft component is controlled by a form factor

$$f(k) = \text{Re} \langle \exp[ik(x_0 - x_1)] \rangle = \text{Re} \left(\frac{i\chi}{kp_1 + i\chi} \right). \quad (5)$$

The ensemble average over the coordinate difference yields the right most expression under the assumption that the distance, z , between hard and soft collision is distributed according to $\exp(-z/\lambda)/\lambda$, where λ is the mean free path for soft interactions. In eq. (5) $\chi = \gamma\beta m/\lambda$.

With the simple form of $f(k)$ above, the angular integration can be performed with the usual Feynman trick, yielding

$$\Delta E_1 = \frac{\alpha}{\pi} \int_0^E d\omega \text{Re} [I(r, z, \gamma) - I(0, z, \gamma)], \quad (6)$$

where the relevant dimensionless variables are $r = -q^2/m^2$ and $z = \gamma\beta/\omega\lambda$, and the integrand is

$$I(r, z, \gamma) = x \log \left(\frac{x+1}{x-1} \right), \quad (7)$$

with

$$x(r, z, \gamma) = \frac{2+r+2i\gamma z}{[(r+2i\gamma z)^2 + 4(r+z^2)]^{1/2}}. \quad (8)$$

Note that $\Delta E_1 = 0$ for $r=0$ as expected, and that in the $\lambda=0$ limit, $\Delta E_1=0$, as well. In the limit $\lambda \rightarrow \infty$ ($z \rightarrow 0$), the interference term drops out, and we recover the familiar result for an isolated soft scattering, $\Delta E_1/E \approx (\alpha/\pi) \frac{2}{3} r$ for $r \lesssim 1$. In the ultra-relativistic limit, on the other hand, where $\gamma \rightarrow \infty$, we find that

$$\Delta E_1 \approx \frac{\alpha}{3\pi} \frac{(m\lambda)^2}{E} (-q^2) \log \left(\frac{E}{m} \right) \rightarrow 0. \quad (9)$$

The vanishing of the soft radiative energy loss in this limit occurs because the time between collisions in the electron rest frame is too short to form the virtual photon field that is usually shaken off after a soft collision [14]. Comparing (9) to the $\lambda \rightarrow \infty$ limit, this suppression effect is seen to become important when the time dilated coherence length, γ/m , exceeds the mean free path λ .

In the non-abelian case, a similar effect is expected due to the finite formation time [14,11], $t_f \sim$

$\cosh(y)/k_\perp$, of gluons as well. However, because the exchanged gluon in the soft interaction can itself radiate gluons, leading to a uniform rapidity radiation distribution [15] rather than the strongly forward peaked radiation with $y \sim \log(E/m)$ familiar in the abelian case, a finite residual radiative energy loss may arise. Dimensional arguments together with gauge invariance requiring ΔE_1 to vanish when $q^2 = 0$ allow in principle a linear in λ contribution, $\Delta E_1 \sim (\alpha_s/\pi) (-q^2)\lambda$ modulo logarithmic factors. Since $\langle -q^2 \rangle \sim g^2 T^2$ in the plasma, the soft radiative energy loss is at most comparable to the collisional one in eq. (1) if the destructive interference effects leading to (9) are reduced in QCD to allow a linear term in λ . In any case, we conclude that the magnitude of dE/dx for jets in a plasma is controlled mainly by the magnitude of the effective coupling and could decrease significantly near T_c .

We turn next to the estimate of the sensitivity of jet quenching to such changes in dE/dx in central collisions of ultra-relativistic nuclei. The number of back-to-back jets with $y_1 = y_2 = 0$ in the CM with $p_{\perp 1} = p_{\perp 2} = p_T$ per central U+U collisions is shown in fig. 1 for various CM energies per incident nucleon ($= \frac{1}{2}\sqrt{s}$). The number of $\{ij\}$ type jets was calculated using [1,2]

$$\frac{dN_{ij}}{dy_1 dy_2 dp_T^2} = T_{AA}(0) \sigma_{ij}(p_T), \quad (10)$$

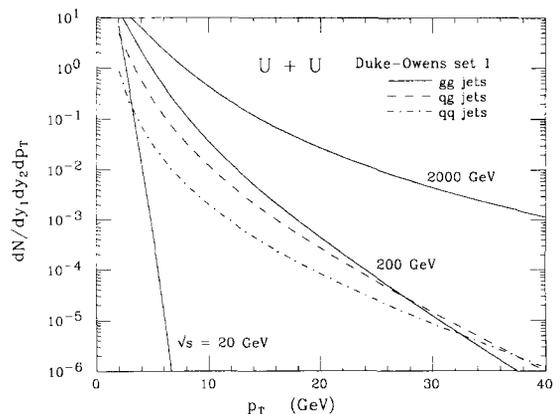


Fig. 1. The density of back-to-back jets at $y_1 = y_2 = 0$ in central U+U collisions at various CM energies per nucleon ($\frac{1}{2}\sqrt{s}$) as a function of the jet transverse momentum p_T . Solid, dashed and dash-dotted show contributions from gg, gq, and qq jets, respectively.

where the 90° CM parton differential cross section is

$$\sigma_{ij}(p_T) = \frac{\pi\alpha_s(p_T^2)}{p_T^4} C_{ij}(xf)_i(xf)_j K, \quad (11)$$

where $(xf)_i \equiv x_T f_i(x_T, p_T^2)$ with $x_T = 2p_T/\sqrt{s}$ and the structure functions, f_i , taken from the Duke-Owens [16] set 1. The factors $C_{gg} = 1.89$, $C_{qg} = 0.76$, and $C_{qq} = 0.16$ include both color and 90° kinematical factors. We took the correction factor $K = 2.5$ to fit measured jet cross sections [17]. The nuclear geometry for $b=0$ collisions of $A+A$ gives rise to the nuclear thickness factor [2], $T_{AA}(0) \approx A^2/\pi R^2$ with $R \approx 1.15A^{1/3}$ fm. Antiquark jets are included in the q terms. In the range $p_\perp \lesssim 25$ GeV, we see that gluon jets dominate at RHIC energies $\sqrt{s} = 200$ GeV. Also it is clear that present SPS energies $\sqrt{s} = 20$ GeV are too small for jet studies.

The energy loss in matter reduces the initial energy $E = p_\perp$ of the jet parton i by an amount

$$\Delta E_i(r, \phi) \approx C_i \int_0^{\tau(r, \phi)} dt \frac{dE(\tau)}{dx}, \quad (12)$$

where $r = |\mathbf{r}|$ is its initial transverse radius and ϕ is the jet angle relative to \mathbf{r} in the plane perpendicular to the beam axis. For a sharp cylindrical geometry with radius R , the energy loss ceases at the escape time, $\tau(r, \phi) = (R^2 - r^2 \sin^2 \phi)^{1/2} - r \cos \phi$. The energy loss depends of course on the space-time evolution of the system. For a constant $dE/dx = \kappa$, $\Delta E_i = C_i \kappa \tau_i$. However, due to longitudinal expansion of the system dE/dx will in general change with time along the jet trajectory. To take into account the expected difference between the energy loss of quarks and gluons [7], the color factor C_i is introduced above such that $C_q = 1$ and $C_g = \frac{9}{4}$.

The time variation of dE/dx is computed assuming the validity of scaling hydrodynamics with Bjorken initial conditions [18]. Longitudinal expansion implies then that the proper entropy density decreases as $\sigma(\tau) = \sigma_0(\tau_0/\tau)$, with $\tau_0 \sim 1$ fm/c. Taking a bag model equation of state [18], the proper energy density $\epsilon(T) = 12T^4 + B \gg \epsilon_Q \equiv \epsilon(T_c)$ in the deconfined plasma phase, $\epsilon \leq \epsilon_H$ in the confined hadron phase, and $\epsilon_H \leq \epsilon \leq \epsilon_Q$ in the mixed phase. The jets propagate in the plasma phase only while $\sigma(\tau) \geq \sigma_Q = \frac{4}{3}(\epsilon_Q - B)/T_c$. For $\sigma_Q \geq \sigma(\tau) \geq \sigma_H \equiv [\epsilon_H + \frac{1}{3}(\epsilon_Q - 4B)]/T_c$ they propagate in the mixed phase with a

fraction $f(\tau) = [\epsilon(\tau) - \epsilon_H]/(\epsilon_Q - \epsilon_H)$ of the matter in the plasma phase and $1 - f$ in the hadronic phase. The time τ_Q where $f(\tau_Q) = 1$ defines the upper bound on the time in the plasma phase, and τ_H where $f(\tau_H) = 0$ defines the lower bound on the time in the pure hadronic phase. The temperature in the plasma phase decreases as $T(\tau)/T_c = (\tau_Q/\tau)^{1/3}$, while it is constant $T = T_c$ between τ_Q and τ_H .

We calculate the integrated energy loss assuming that dE/dx is approximately constant $\kappa_H = 1$ GeV/fm in the hadronic phase. In the plasma phase, $dE/dx \approx \kappa_Q (\tau_Q/\tau)^{2/3}$ neglecting the slow time dependence of the logarithmic factor in eq. (1) and denoting the plasma energy loss at T_c by κ_Q . Assuming furthermore a uniform transverse distribution of the matter, the energy loss is then given by

$$\begin{aligned} \frac{dE(\tau)}{dx} &= \kappa_Q \left(\frac{\tau_Q}{\tau} \right)^{2/3} \theta(\tau_Q - \tau) + \kappa_H \theta(\tau - \tau_H) \\ &+ [(\kappa_Q - \kappa_H)f(\tau) + \kappa_H] \theta(t - \tau_Q) \theta(\tau_H - \tau). \end{aligned} \quad (13)$$

The initial transverse coordinate and jet orientation in this case only affect the time spent in the matter, $\tau_i(r, \phi)$. The values of τ_Q and τ_H depend of course on the assumed initial equilibrated energy density $\epsilon_0 = \epsilon(\tau_0)$ at the thermalization time τ_0 .

The final momenta of each jet is then given by $p_i = p_\perp - \Delta E_i$, where p_\perp is the initial transverse momentum. Since the final momenta of the two jets are sensitive to the unmeasurable initial coordinates (r, ϕ) , it is advantageous to consider the distribution of the summed energy $E_{tot} = p_1 + p_2 = 2p_\perp - (\Delta E_1 + \Delta E_2)$. That distribution is maximally sensitive to the nuclear size, and fluctuations of ΔE_i due to fluctuations of the time spent in the matter cancel out in the sum to a large extent. The distribution of the summed energy also minimizes the trigger bias inherent in single jet distributions. Measurements of a single jet will obviously be biased to those events where the hard scattering originates near the nuclear surface. In the summed distribution, at least one of the jets propagates through the dense matter. In addition, the A dependent Cronin enhancement of high p_\perp hadrons due to multiple initial state interactions is minimized by concentrating on correlated back-to-back jets [19]. Therefore, that summed distribution amplifies as much as possible the quenching effect.

In practice, E_{tot} , must be defined in terms of the sum of the energies of hadrons in a narrow forward-backward angular cone with $\Delta\theta \sim 10^\circ$ oriented 90° to the beam axis. That cone includes all fragmented hadrons with energy $\gtrsim 2$ GeV and transverse momentum $\lesssim 0.4$ GeV relative to the jet axis. It excludes the low energy hadrons with $E \lesssim 2$ GeV produced as a result of the jet energy loss. Therefore, the measured E_{tot} can be approximated by the initial jet energy minus the summed energy loss.

Averaging over the initial transverse coordinate and jet orientation assuming the collision of uniform spherical nuclei, the number of back-to-back jets with $y_1 = y_2 = 0$ and a given E_{tot} is reduced in central ($b=0$) $A+A$ collisions by

$$R_{AA}(E_{\text{tot}}) = \frac{1}{\sigma_0(E_{\text{tot}})} \int_0^{R_A^2} \frac{2 dr^2}{R_A^4} (R_A^2 - r^2) \times \int_0^{2\pi} \frac{d\phi}{2\pi} \sum_{ij} \sigma_{ij}(\frac{1}{2}(E_{\text{tot}} + \Delta E_i + \Delta E_j)), \quad (14)$$

where $\Delta E_i = \Delta E_i(r, \phi)$, $\Delta E_j = \Delta E_j(r, \pi - \phi)$, and $\sigma_0(E) = \sum_{ij} \sigma_{ij}(\frac{1}{2}E)$ with the sum over both quark and gluon jets. The calculated jet quenching factor in central U+U collisions at $\sqrt{s}=200$ GeV (RHIC energies) is shown in fig. 2 as a function of $E_{\text{tot}} = p_{T1} + p_{T2}$. The A dependence of jet quenching for fixed

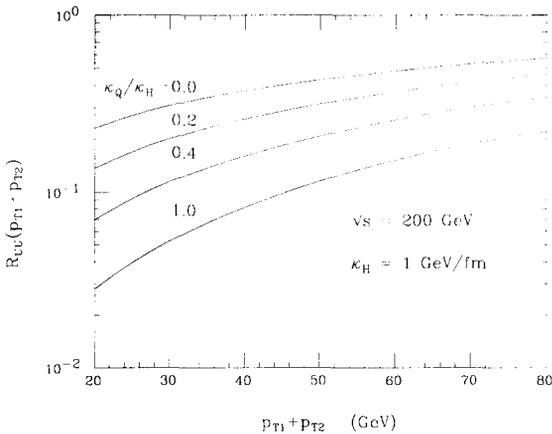


Fig. 2. Dijet reduction factor for central U+U collisions at 100 A GeV as a function of the dijet energy $E = p_{T1} + p_{T2}$. Different curves are labeled by the ratio κ_Q/κ_H of dE/dx just above and below T_c assuming $\kappa_H = 1$ GeV/fm. All curves are normalized such that $R_{pp}(E) = 1$.

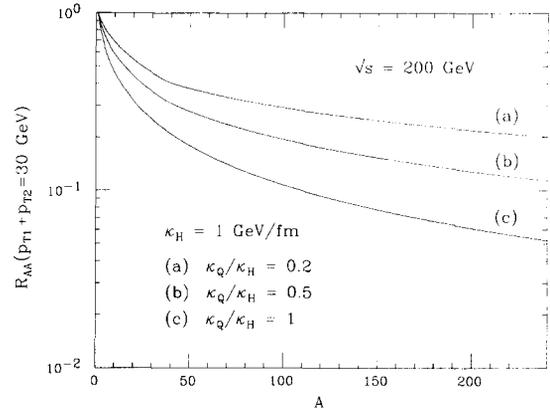


Fig. 3. Dijet reduction factor for central $A+A$ collisions at 100 A GeV as a function of A for dijet energy $E = 30$. Curves are labeled by the ratio κ_Q/κ_H as in fig. 2 and dE/dx just above and below T_c assuming $\kappa_H = 1$ GeV/fm. All curves are normalized such that $R_{pp}(30) = 1$.

$E_{\text{tot}} = 30$ GeV is shown in fig. 3. In both figures the normalized ratio R_{AA}/R_{pp} is shown where $R_{pp} \sim 0.85$ for the initial conditions below. The bag model parameters were chosen such that $T_c = 190$ MeV, $\epsilon_Q = 2.5$ GeV/fm³, $\epsilon_H = 0.5$ GeV/fm³, and $B = 0.5$ GeV/fm³. The initial conditions for these calculations were assumed to be $\tau_0 = 1$ fm/c and

$$\epsilon_0 = \epsilon_s A^{1/3} + \epsilon_h A^{2/3}, \quad (15)$$

where the energy density due to soft multiparticle production processes is $\epsilon_s \approx 0.5$ GeV/fm³ that scales with the number of interacting “wounded” nucleons per unit area $\propto A^{1/3}$, and the energy density at τ_0 due to semi-hard mini-jets is $\epsilon_h(\sqrt{s}=200) \approx 0.08$ GeV/fm³ that scales with the number of binary inelastic processes per unit area $\propto A^{2/3}$. The above estimates for ϵ_s and ϵ_h are taken from the results of ref. [2]. For U+U collisions at (RHIC) collider energies $\sqrt{s} = 200$ GeV, approximately one half the initial energy density is expected to arise from the mini-jet contribution. We note that $\epsilon_h(\sqrt{s})$ is expected to vary approximately linearly with \sqrt{s} in that domain.

We find that jet quenching in heavy nuclei is expected to reduce the rate of back-to-back jets with $E_{\text{tot}} \sim 30$ GeV by an order of magnitude. The different curves illustrate the sensitivity of the quenching to the ratio of dE/dx in the two phases. (For these calculations we neglect the weak time dependence of

dE/dx in the plasma phase since the initial temperature remains always close to T_c for these initial conditions.) A reduction of dE/dx just above T_c by $\kappa_Q/\kappa_H \lesssim \frac{1}{2}$ is seen to enhance jets relative to the $\kappa_Q = \kappa_H$ case by a factor $\gtrsim 2$. However, even for $\kappa_Q = 0$ the energy loss due to jet propagation through the mixed and hadronic phases results in a substantial amount of quenching. From fig. 2 it is also clear that the quenching factor is most sensitive to variations in κ_Q at lower E_{tot} because the ΔE_i are limited due to the finite size of nuclei. The optimal E_{tot} for study will depend of course on the lowest resolvable jet energies.

For $\kappa_Q/\kappa_H = 0.2$, we see in fig. 3 a small discontinuity of the slope of R_{AA} near $A \approx 40$, marking the critical A beyond which $\epsilon_0 > \epsilon_Q$ in our case. Smaller ϵ_U or larger ϵ_Q shift that point to higher A . In addition, variation of the incident energy will shift that point due to the \sqrt{s} dependence of the mini-jet component of ϵ_0 . Of course, for smaller changes in $\kappa_Q/\kappa_H \sim \frac{1}{2}$ the discontinuity in the slope is very small and is likely to be washed out by straggling and surface effects. On the other hand, a reduction of the latent heat, $\epsilon_Q - \epsilon_H$, at fixed ϵ_Q enhances the discontinuity. It is clear that a high precision A and \sqrt{s} dependence study of jet quenching would be necessary to look for onset of the deconfinement transition in this way.

Finally, we consider the sensitivity of the results to the assumed initial conditions taking into account the full time dependence of dE/dx in eq. (13). For a given ϵ_0 corresponding to an initial temperature $T_0 \approx [\frac{1}{12}(\epsilon_0 - B)]^{1/4}$ the uncertainty principle places a lower bound on the thermalization time $\tau_0 \gtrsim 1/T_0$. At this time, the interacting nuclei have only separated by the uncertainty in the localization of an average thermal parton along the beam axis. Taking the lower bound for the thermalization time, the dependence of the absolute quenching factor for central U+U collisions at $E_{\text{tot}} = 30$ GeV on ϵ_0 is shown in fig. 4. Note that for $\kappa_Q/\kappa_H < \frac{1}{2}$, the quenching decreases with increasing ϵ_0 until $dE/dx \approx \kappa_Q(T/T_c)^2$ exceeds κ_H in the confined phase. However, if Landau initial conditions [20] were to apply with $\epsilon_L \sim 2\gamma_{\text{cm}}^2 \epsilon_{\text{nuc}} \sim 3$ TeV/fm³!, then jets would in fact be much more attenuated than if the more probable $\epsilon_0 \lesssim 10$ GeV/fm³ initial conditions apply. Finally, we note that the quenching factor is sensitive to the assumed thermalization time τ_0 . The dashed curve shows the results for $\tau_0 = 1/3T_0$.

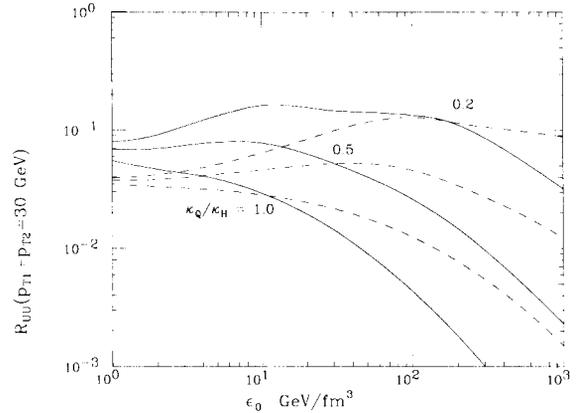


Fig. 4. Absolute dijet reduction factor for central U+U collisions at 100 A GeV for dijet energy $E = 30$ as a function of the initial energy density ϵ_0 assuming a thermalization time $\tau_0 = 1/T_0$ (solid curves) and $\tau_0 = 1/3T_0$ (dashed curves).

In summary, we have shown that energy loss in dense matter is expected to reduce the jet yields with $p_T \sim 20$ GeV/c in ultra-relativistic nuclear collisions by a factor ~ 10 . The detailed A and E_{tot} dependence of back-to-back jets was shown to be sensitive to changes in the energy loss mechanism near the deconfinement transition. Finally, we motivated and explored the consequences of a possible reduction of dE/dx near T_c and suggested how jet systematics may help in that case to identify quark-gluon plasma formation in such reactions.

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