LBL-31002
Preprint

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Submitted to Nuclear Physics A

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July 1991

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Prepared for the U.S. Department of Energy under Contract Number DE-AC03-76SF00098

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High P_T Probes of Nuclear Collisions *

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Abstract

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*This work was supported by the Director, Office of Energy Research, Division of Nuclear Physics of the Office of High Energy and Nuclear Physics of the U.S. Department of Energy under Contract No. DE-AC03-76SF00098.

[1]Invited talk at the 4th Conference on the Intersections between Particle and Nuclear Physics, Tucson, Arizona, May 24-29, 1991, and at the 4th International Conference on Nucleus-Nucleus Collisions, Kanazawa, Japan, June 10-14, 1991

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ABSTRACT

Jet quenching in high energy nuclear collisions is shown to provide new information on the energy loss mechanisms in dense matter near the quark-gluon plasma transition. In addition it is shown that A dependence of the inclusive p_T spectra in $p + A$ provides novel constraints on gluon shadowing in nuclei.

INTRODUCTION

One of the interesting predictions[1, 2] of QCD is the transition from a confined, chirally broken phase of matter composed of hadrons to a deconfined, chirally symmetric state of quasifree quarks and gluons called the quark-gluon plasma (QGP). The order of the transition is still unclear as present QCD lattice simulations[2] with $2\frac{1}{2}$ flavors of dynamical quarks are restricted to small, coarse lattices ($16^3 \times 4$ with $\Delta x \sim 1/4$ fm). However, thermodynamic quantities such as the entropy density are found to very rapidly in the vicinity of the critical temperature $T_c \sim 200$ MeV. Future teraflop computers will certainly map out the behavior of thermodynamic quantities and correlation lengths in much more detail. To study this new phase of matter in the laboratory, experiments have been proposed using ultrarelativistic nuclear collisions. In such collisions extreme conditions necessary to form a QGP with $\epsilon \sim 20$ GeV/fm^3 may be created over relatively large space-time volumes $\sim 10^3\ fm^4 \gg \Lambda_{QCD}^{-4}$.

The present experimental program is restricted to the study of light ion collisions at energies 10-200 AGeV. However, experiments on $Au + Au$ at $\sqrt{s} \sim 200$ AGeV will become possible around 1998 at RHIC/BNL. During the past decade many different signatures of the QGP transition have been proposed[1]. One class of signatures involve the study of the low p_T single particle inclusive cross sections, $d\sigma_h/dyd^2p_T$, which are expected to reveal information on the entropy density, collective hydrodynamic flow

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phenomena, and flavor composition of the produced dense matter. Correlations between identical particles, e.g. $\pi^-\pi^-$, can be used to probe the space-time geometry of the final stages of such collisions. Antibaryon yields provide a measure of the baryon densities achieved in such collisions. Penetrating probes such as dileptons, direct photons, and heavy quarkonium production are expected to reveal information on the quasiparticle spectrum, temperature, and correlation lengths in that system.

In this talk we consider another class of observables and signatures that only become accessible at collider energies at RHIC and beyond (LHC). The other observables listed above have already received much theoretical and experimental attention[1]. Indeed, striking A dependence of those observables have already been seen in present light ion reactions, and the interpretation of those data is a major focus of present work on heavy ion physics[3]. The high p_T observables that we discuss here have on the other hand received less attention[4, 5, 6]. Our intent is to point out some of the new physics that may be extracted from such observables. In particular, high p_T processes essentially dominate the initial conditions and the event structure and provide a unique and complementary probe of the dense matter produced in nuclear reactions. In addition these processes are of broader interest since they naturally belong to a class of problems at the interface of particle and nuclear physics and overlap with the experimental programs on high p_T observables in lepton-nucleus and hadron-nucleus reactions.

WHY HIGH p_T?

The main advantage of high p_T is that the initial rate of such processes is calculable via perturbative QCD (pQCD) due to the factorization theorem. Given the parton structure functions, $f_{a/A}(x, Q^2)$, and perturbative cross sections, $d\sigma_{ab}$ for parton-parton scattering, jet cross sections can be calculated[7] via

$$\frac{d\sigma_{jet}}{dP_T^2 dy_1 dy_2} = K \sum_{a,b} x_1 f_a(x_1, P_T^2) x_2 f_b(x_2, P_T^2) d\sigma^{ab}(\hat{s}, \hat{t}, \hat{u})/d\hat{t}. \quad (1)$$

Jet data[8] can be well reproduced using even the lowest order parton cross sections with a factor $K \approx 2$ to account for higher orders. Another important aspect of high p_T processes is that they occur on much shorter time scales, $\tau \sim \hbar/p_T << 1$ fm/c, than the far more abundant low $p_T \sim \Lambda_{QCD}$ processes which determine the properties of the dense matter. Thus, all high p_T partons are formed while the two thin nuclear discs pass through each other on a time scale $R_A/\gamma_{cm} \sim 0.1$ fm/c at RHIC before the equilibrated QGP has formed. These high p_T partons can therefore serve as effective "external" probes of the produced dense matter. The essential point is that the *initial flux* of those high p_T partons is calculable and experimentally testable in pp collisions.

Another reason for considering high p_T physics is that semi-hard reactions such as $gg \to gg$ occurring on a scale $p_T \gtrsim p_0 \sim 2$ GeV are expected to be so abundant[9]

that they may well dominate the evolution of the dynamics in AA collisions at collider energies. While not resolvable as distinct jets, the rate of such "minijets" is calculable via pQCD. For $\sqrt{s} = 200$ (2000) GeV the inclusive minijet cross section is large, $\sigma_{jet}(p_0 = 2) \approx 10(70)$ mb (see [10, 11]). The total number of partons with $p_T > p_0$ in nuclear collisions can therefore be estimated via

$$N_j = 2T_{AA}(b)\sigma_{jet}(p_0) \ , \tag{2}$$

where $T_{AA}(b)$ is the nuclear overlap function[10] that counts the number of binary pp collisions per unit area at an impact parameter b. For $b = 0$, $T_{AA}(0) \approx A^2/\pi R_A^2$, resulting in $T_{AuAu}(0) \approx 30/\text{mb}$. Therefore we expect ~ 600 partons with $p_T > 2$ GeV in central $Au + Au$ collisions at RHIC. There will be then more minijets than valence diquark-quark (beam) jets in such collisions.

It has been estimated[10] that more than half of the final transverse energy per unit rapidity could be produced by minijets. This can be seen via

$$dE_T^h/dy \sim \langle p_T \rangle T_{AA}(0)\sigma_{jet}(p_0)/\Delta Y \sim 300 \text{ GeV},$$

where $\langle p_T \rangle \approx 3$ GeV for minijets above 2 GeV. In contrast, the contribution due to fragmentation of beam jets is only $dE_T^s/dy \sim A \times 1$ GeV ~ 200 GeV approximately independent of beam energy. In fact, as we show below, detailed Monte Carlo simulations[11] indicate that the minijet contribution could be twice as large as the above simple estimate due to initial and final state radiation effects but that gluon shadowing may effectively compensate for that increase.

The large transverse energy due to minijets is an important source for increasing the energy density of matter produced at midrapidity. Due to longitudinal expansion, the energy density in AA collisions decreases with time as $\epsilon(\tau) \approx \epsilon_0\tau_0/\tau$, where the initial energy density can be estimated via[1]

$$\epsilon_0 \approx dE_T/dy(\tau_0 \pi R_A^2)^{-1} \ . \tag{3}$$

By the uncertainty principle minijets are produced on a time scale $\tau_0 \gtrsim \hbar/p_0 \sim 0.1$ fm/c. The initial energy density due to minijet production in $Au + Au$ can therefore reach $\epsilon_0 \sim 20$ GeV/fm^3 at RHIC energies. In contrast, the fragmentation of soft beam jets occurs on a longer time scale $\tau_s \sim 1$ fm/c and leads to only an extra ~ 2 GeV/fm^3 at τ_s. The contribution from minijets is therefore essential to enhance the time spent in the QGP phase.

How do we know if the minijet contribution really exists if they are not resolvable as distinct jets? There has been much work on this problem[9]. We show in Fig.1 new results based on the recently developed Monte Carlo event generator HIJING[11] demonstrating the present level of quantitative understanding of multiparticle production in

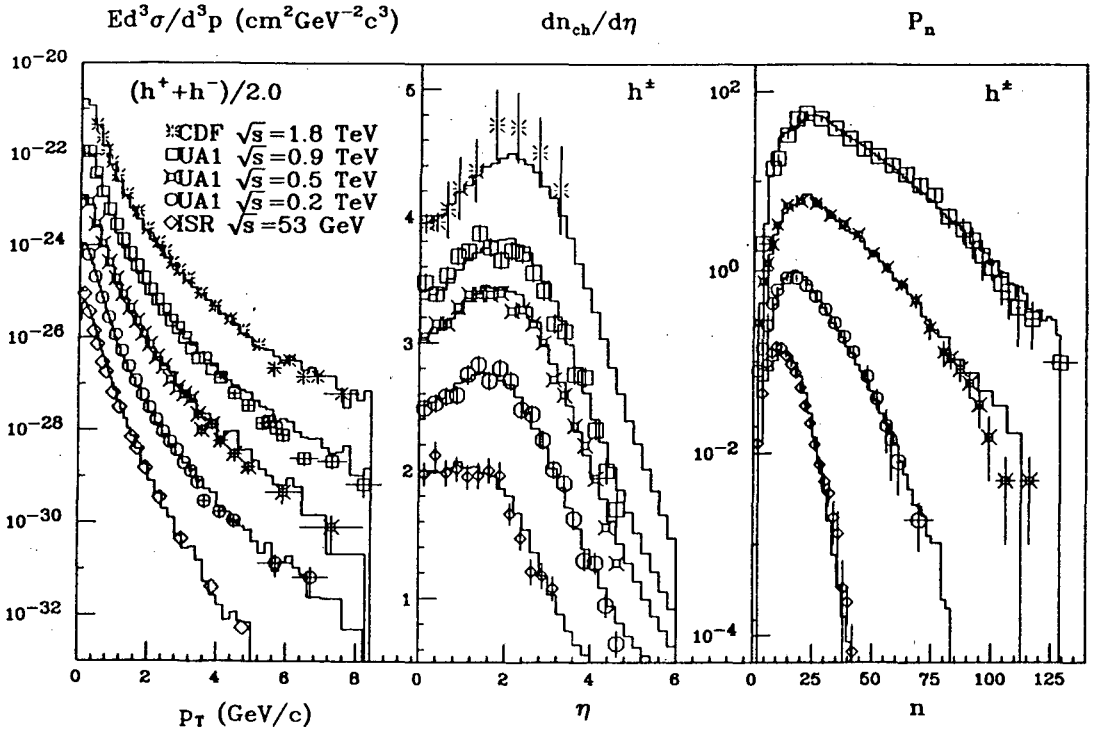

Figure 1: Data on charged particle inclusive p_T spectra, pseudorapidity distributions, and multiplicity fluctuations in pp and $p\bar{p}$ reactions compared to HIJING[11] simulations (histograms) as a function of $\sqrt{s}$. The transverse momentum spectra and multiplicity distributions have been displaced by extra factors of 10 relative to the absolutely normalized $\sqrt{s} = 53$ GeV data.

pp over a wide energy range, $\sqrt{s} = 50 - 2000$ GeV.

We see that, as $\sqrt{s}$ increases, the conspicuous power law tail characteristic of pQCD develops for the transverse momenta spectra. Furthermore, its magnitude and energy dependence can be quantitatively understood in terms of minijets. HIJING[11] combines multiple pQCD minijet physics[12] with the soft beam jet LUND phenomenology[13]. Hadronization is treated according to LUND string phenomenology[14] in which a gluon minijet is represented by a kink in one of the beam $qq - q$ strings and the JETSET7.2 jet fragmentation routine with default parameters is used.

From Fig.1 we see that the model accounts well for the energy dependence of not only the transverse momentum distributions but also the charged particle pseudorapidity distributions as well as the multiplicity fluctuation in pp and $p\bar{p}$ reactions. This overall quantitative understanding of multiparticle observables in pp inelastic reactions supports strongly the importance of minijet physics at collider energies.

WHAT HIGH p_T PARTONS PROBE IN A+A?

QCD factorization[15] allows one to calculate inclusive high p_T hadron cross sections for $A + B \to h + X$ as a convolution

$$d\sigma_{AB\to h} = [f_{a/A}(x_1, Q^2) f_{b/B}(x_2, Q^2)] \otimes [d\sigma_{ab\to c}] \otimes [F_{h/c}(z, Q^2)] \ , \tag{4}$$

where as in eq.(1), $f_{a/A}$ are the parton structure functions, $d\sigma_{ab\to c}$ are the pQCD cross sections, and where $F_{h/c}(z, Q^2)$ is the partonic fragmentation function into hadrons carrying a fraction z of the partonic momentum. Eq.(4) holds strictly only for $p_T \sim zQ \to \infty$. For finite Q^2, factorization breaks down at low z because the hadronization of parton c cannot by treated independently from the fragmentation of the beam jets due to interference effects (the string effect), and it breaks down for fixed z and Q^2 when the system becomes so large ($A \to \infty$) that final state interactions with the comoving partons cannot be neglected. It is the nuclear modification of the *effective* parton fragmentation functions that we want to use in order to probe the dense matter produced in nuclear collisions.

To illustrate the physics, consider the fragmentation of a gluon produced initially with a high p_{T0} at $y = 0$. As that gluon traverses the evolving dense comoving matter, it suffers energy loss, $\kappa \equiv dE/dx$, and random momentum kicks, $\delta\vec{p}$, transverse to its motion. The partonic final state interactions involve elastic scattering and induced gluon radiation (bremsstrahlung). The random momentum kicks will lead to an enhanced acoplanarity with respect to its partner high p_T parton moving in the opposite direction[5]. The energy loss will reduce the energy of the gluon[4]. When the gluon exits the dense cylindrical matter of radius, $\sim R_A$, it emerges into the QCD vacuum with a reduced average final transverse momentum

$$p_{Tf} \approx p_{T0} - \Delta E \ , \tag{5}$$

where $\Delta E = \int d\tau \kappa(\vec{x}_0 + \vec{v}_0\tau) \sim \kappa R_A$ is the net energy loss in the medium that depends of course on its initial production coordinate, $\vec{x}_0$, and velocity $\vec{v}_0$. This energy loss leads to jet quenching in AA collisions[4, 6]. If the pQCD initial inclusive jet distribution (in the absence of final state interactions) is parametrized as N_0/p_{T0}^n, then the ratio of the number of final jets to the number of initial jets with a given p_T is given by

$$R(p_T) \approx (1 + \Delta E/p_T)^{-n} \ . \tag{6}$$

The rate of jets with a given p_T is thus reduced relative to the pQCD rate by the factor $R(p_T)$. The inclusive high p_T hadron yields will then also be suppressed since

the effective fragmentation function will be

$$F^{eff}_{h/c}(z = p_T/p_{T0}) \approx F^{eff}_{h/c}(zp_{T0}/p_{Tf})p_{T0}/p_{Tf} \quad . \tag{7}$$

By studying the A dependence of the ratio high p_T jets or hardons, we can extract information on dE/dx as a function of the initial energy density of the comoving matter.

In ref.[16] we analyzed $\ell + A \to h + X$ data for evidence of jet quenching in cold nuclear matter. We found that the energy dependence of the suppression of high z hadrons could be reproduced assuming $dE/dx \sim 1$ GeV/fm in nuclei. New preliminary data from E665/FNAL[25] seems to confirm the trends in those earlier data and will eventually extend the data base considerably by mapping out the A dependence of the quark fragmentation function over a large kinematic domain, including the $x_F < 0$ domain. It will be especially interesting to compare jet quenching in ℓA and AA to look for differences in dE/dx in nuclear and QGP matter.

In addition to jet quenching, the A dependence of moderate high $p_T \sim 3 - 10$ GeV inclusive spectra can be used to extract information on the phenomenon of gluon shadowing in nuclei. As is well known from deep inelastic reactions[17], the quark distribution in nuclei is not additive, i.e. $f_{q/A}(x) \neq Af_{q/N}(x)$. Currently, much interest has focused on the physics of nuclear shadowing[15] that occurs for $x \ll 1/p_0R_A \sim 0.1$. For nuclear collisions at RHIC the number of minijets on a scale $p_0 \sim 2$ GeV depends on the number of gluons in a nucleus with $x \sim 2p_0/\sqrt{s} \sim 0.02$, which is well inside the shadowed domain. The magnitude of shadowing of quarks in heavy nuclei is quite large[17], $f_{q/Au}(0.02)/Af_{q/N}(0.02) \sim 0.6$. Of course, nothing is known experimentally about shadowing of gluons. However, because of the abundance of minijets in AA, the inclusive moderate p_T distributions of hadrons, especially in $p+A$ at RHIC, will provide a new tool to probe gluon shadowing.

dE/dx In DENSE MATTER

The first estimate of dE/dx of quarks in dense QCD matter was made by Bjorken[4]. The energy loss due to elastic scattering with partons in the medium is

$$dE_{el}/dx = \int_{\mu^2}^{s/4} d\sigma_{el}\rho\nu \approx 2\pi\alpha_s^2\langle\rho/2k\rangle \log s/4\mu^2 \quad , \tag{8}$$

where in a QGP

$$\langle\rho/2k\rangle = \langle(\rho_g + 4/9\rho_q)/2k\rangle \approx 0.7T^2 \quad ,$$

is the effective density of thermal partons that can scatter with the incident quark divided by their average transverse momentum, and $\nu \sim t/2k$ is the energy loss per scattering. The familiar Coulomb logrithmic divergence is cut off from above by the kinematics of the scattering with $s \approx 6ET$, where E is the energy of the jet. The

lower cutoff depends on the infrared color screening cutoff scale in the medium. In a perturbative QGP, $\mu \approx \mu_{Debye} \approx gT$.

For $E = 30$ GeV, $\alpha_s \approx 0.25$, and $T \approx 200$ MeV, $dE_{el}/dx \sim 0.2 - 0.3$ GeV/fm. This value is remarkably smaller than the dynamic $dE_d/dx \sim 1 - 2$ GeV/fm associated with the jet fragmentation into hadrons[19]. That dynamic (string) tension can be thought of as the dE/dx of quarks in the nonperturbative QCD vacuum. The elastic energy loss is also considerably smaller than the value deduced in [16] from jet quenching in cold nuclei via eA. Recent detailed calculations[20] using high temperature QCD resummation techniques have confirmed the smallness of the elastic contribution to energy loss.

In ref.[6], it was suggested that a small value of dE/dx in a QGP could lead to a relative reduction of jet quenching and may provide one of the signals for the QGP transitions in nuclear collisions. However, we recently found that at least deep in the QGP phase, the induced radiative energy loss could be quite large[21]. The spectrum of soft gluons radiated when two partons scatter with momentum transfer, $q_\perp$, in pQCD has been calculated in ref.[18] as

$$\frac{d\sigma_g}{d^2q_\perp dy d^2k_\perp} \approx \frac{d\sigma_{el}}{d^2q_\perp}\frac{dn_g}{dyd^2k_\perp} \approx \frac{C_i 2\pi\alpha_s^2}{q_\perp^4}\left\{\frac{C_A\alpha_s}{\pi^2}\frac{q_\perp^2}{k_\perp^2(\vec{k}_\perp - \vec{q}_\perp)^2}\right\} , \tag{9}$$

where the Rutherford color factor $C_i = 4/9, 1, 9/4$ for qq, qg, gg scattering and $C_A = 3$. Note that unlike the usual Altarelli-Parisi splitting functions, the *induced* radiation comes with the color factor C_A for both quarks and gluons. Also in contrast to QED, the radiated gluons are distributed approximately uniform in rapidity and limited in the transverse direction by the mean transverse momentum scale, q_T^2, of final state interactions in the medium. That scale in turn is controlled by the infrared screening scale, μ, as

$$\langle q_\perp^2\rangle \sim (\int_{\mu^2}^{s/4} d\sigma_{el}\, t)/(\int_{\mu^2}^{s/4} d\sigma_{el}) \sim \mu^2 \log(s/4\mu^2) \ . \tag{10}$$

The radiated energy loss can be estimated by integrating the energy weighted gluon distribution. However, it is essential at this point to take into account that the distance between interactions in the medium is relatively small:

$$\Delta x \sim \lambda_s = (\sigma_{el}\rho)^{-1} \sim 1/(\alpha_s T) \ . \tag{11}$$

In that interval, the induced radiation is limited to only those gluons that can be formed in a time less than Δx. The formation time is simply the time required by the uncertainty principle for a gluon with a transverse size $1/k_\perp$ to separate from its parent parton:

$$\tau_f(k_\perp, y) \sim \cosh(y)/k_\perp \ . \tag{12}$$

Because of this (Landau-Pomeranchuk-Feinberg) effect[22], the mean radiated energy

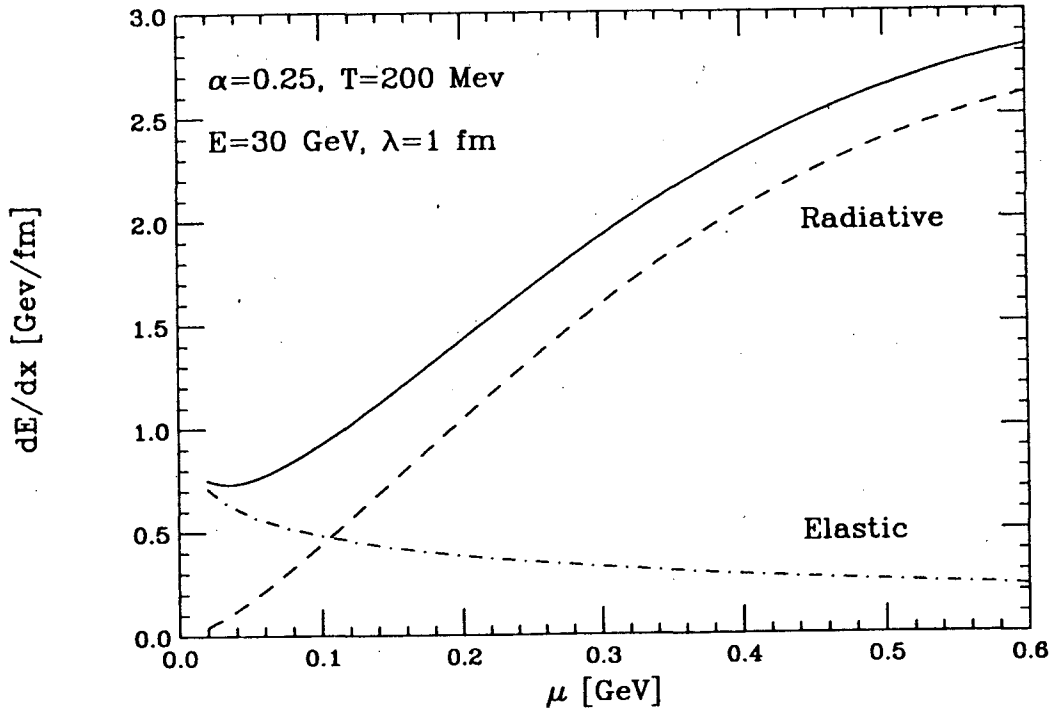

Figure 2: Comparison of estimated radiative (dashed) and elastic (dash-dot) energy loss as a function of the infrared screening scale μ.

loss over a distance λ is limited to

$$\Delta E_g(\lambda) \sim \langle \int dn_g \, k_\perp \cosh y \, \theta(x - \tau_f) \rangle \approx \lambda dE_{rad}/dx \quad , \tag{13}$$

where the $\langle \cdots \rangle$ denotes the average over $q_p erp$ via $d\sigma_{el}$, and the integration domain over the final gluon momenta is limited by kinematics to the range $k_\perp^2 < s/4$ and $k_\perp \cosh(y) < E$ with $s \approx 6ET$. The approximate linear dependence on λ holds only in the $\lambda \ll E/\mu^2$ limit. In the opposite limit, $\lambda \gg E/\mu^2$, the radiated energy loss saturates with $\Delta E_g(\infty) \sim \alpha_s E$.

In the small λ of high E limit, the following estimate of the induced radiative energy loss was derived in [21]:

$$dE_{rad}/dx = \frac{C_A \alpha_s}{\pi} \mu^2 L(E/\lambda\mu^2, s/4\mu^2) \quad , \tag{14}$$

where $L(x,y) \approx log(x)(log(x) - 1 + 1/x) + log(y)(1 - x/y)$ is a double logrithmic function in the limit $1 \ll x < y$. The above approximate form of the integrals was derived by regulating the propagators in eq.(9) with a constant self energy insertion, e.g., $q_\perp^2 \to q_\perp^2 + \mu^2$.

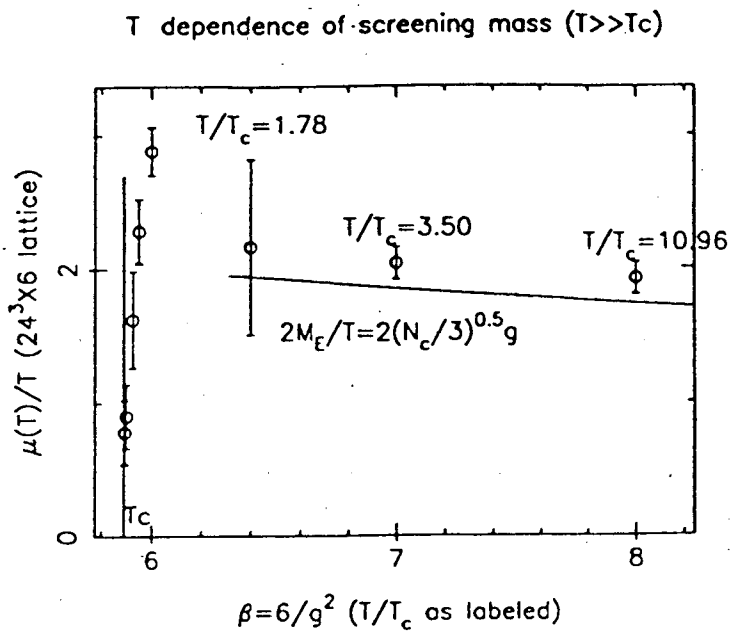

Figure 3: Lattice QCD ($N_f = 0$) estimates[23] of the Wilson line screening scale, $\mu(T)$, as a function of temperature. The solid line corresponds to the expected perturbative Debye scale at high T.

For $\mu = 0.2 - 0.4$ GeV, $\alpha_s = 0.25$, $E = 30$ GeV, $T = 0.2$ GeV, the induced radiative energy loss can thus be rather large $\sim 1 - 2$ GeV/fm as shown in Fig. 2. Of course, (14) is only an order of magnitude estimate because a sharp cutoff distribution for the formation time distribution was assumed and the regulation of the infrared singularities should be treated more carefully[20]. However, it is sufficient to estimate that radiative energy loss is likely to dominate deep inside a QGP, where $\mu \sim gT$ is expected to be large due to Debye screening. If dE/dx remains a few GeV/fm across the QGP transition, then, as shown in ref.[6], back to back jets with $p_T \sim 20 - 30$ GeV would be suppressed by an order of magnitude.

On the other hand, the systematics of jet quenching as a function of A may lead to surprises if there is a significant variation of the infrared cutoff scale, $\mu(T)$, near T_c. While little is yet known about $\mu(T)$ for physical QCD, lattice simulations[23] for $N_f = 0$ indicate a rapid variations of $\mu(T)$ is indeed possible as shown in Fig. 3. It is important to recall[1] that during the first few fm/c, the dense matter evolves mostly in the mixed phase with $T \approx T_c$. Jet quenching in AA is thus especially sensitive to dE/dx near T_c. A sharp reduction of μ as in Fig.3 would effectively suppress the induced radiation as Fig.2 indicates. This would lead to a less rapid increase of jet quenching as

a function of the incident A than expected if dE/dx were a constant and R_A increases. The observation of such a variation in conjunction with other signatures could in this way help to identify the QGP transition threshold[6].

Numerical Simulations

To test more quantitatively how jet quenching could affect the final hadron spectra in $A + A$, we turn next to numerical simulations based on the HIJING Monte Carlo event generator[11]. As noted above HIJING combines multiple high p_T jet and soft beam jet physics and reproduces well the multiparticle observables in pp reactions. For AA collisions it utilizes the Glauber multiple scattering formalism as in [13] and was tested on pA and OA data at $E_{lab} = 200$ AGeV (see Ref.[11] for a detailed discussion). For the present discussion we note two unique features in HIJING which are relevant for high p_T physics. (1) It incorporates a model of the impact parameter dependence of the nuclear modification of the structure functions. (2) It also includes a simple gluon splitting mechanism to model induced radiative energy loss in dense matter.

To test the sensitivity of the results to nuclear modifications of the structure functions, we assumed that both quarks and gluons are shadowed by the same amount in the region $x \lesssim 0.1$

$$f_{a/A}(x) = A f_{a/N}(x)[1 - a(x)(A^{1/3} - 1)\frac{4}{3}(1 - b^2/R^2)^{1/2}] \ , \tag{15}$$

where $a(x) = 0.1 \exp(-x^2/0.01)$. This parametrization is motivated by geometrical considerations[24] and the A dependence of the EMC data[17]. For the full range of x a more complex parametrization is used in HIJING[11].

Jet quenching in HIJING is implemented by a simple gluon splitting scheme. The transverse coordinates of interaction points are first determined using cylindrical geometry assuming a constant mean free path, λ_s. If Δx is the distance between two interaction points, the total radiated energy is taken to be $\kappa \Delta x$, where κ is the assumed dE/dx. This radiated energy is distributed into collinear gluons with a uniform rapidity distribution. For hadronization purposes the split off gluons are added as new kinks to the beam jet $qq - q$ string passing through the interaction point. In effect this model involves the exchange of gluon kinks between neighboring valence $qq - q$ strings. In this way, the energy of an initial hard gluon kink on one string is eventually shared between several strings separated in the transverse direction by a distance on the average, $\sim \lambda_s$. Every splitting introduces a new kink on a neighboring string with transverse energy $\sim \lambda_s \kappa$. With this model we study the effects of jet quenching by varying the two model parameters, κ and λ_s.

Shown in Fig.4 are results obtained via HIJING for central $Au + Au$ at $\sqrt{s} = 200$ AGeV. Note that without minijets (dotted curve), the $2A$ valence $qq - q$ beam jets lead

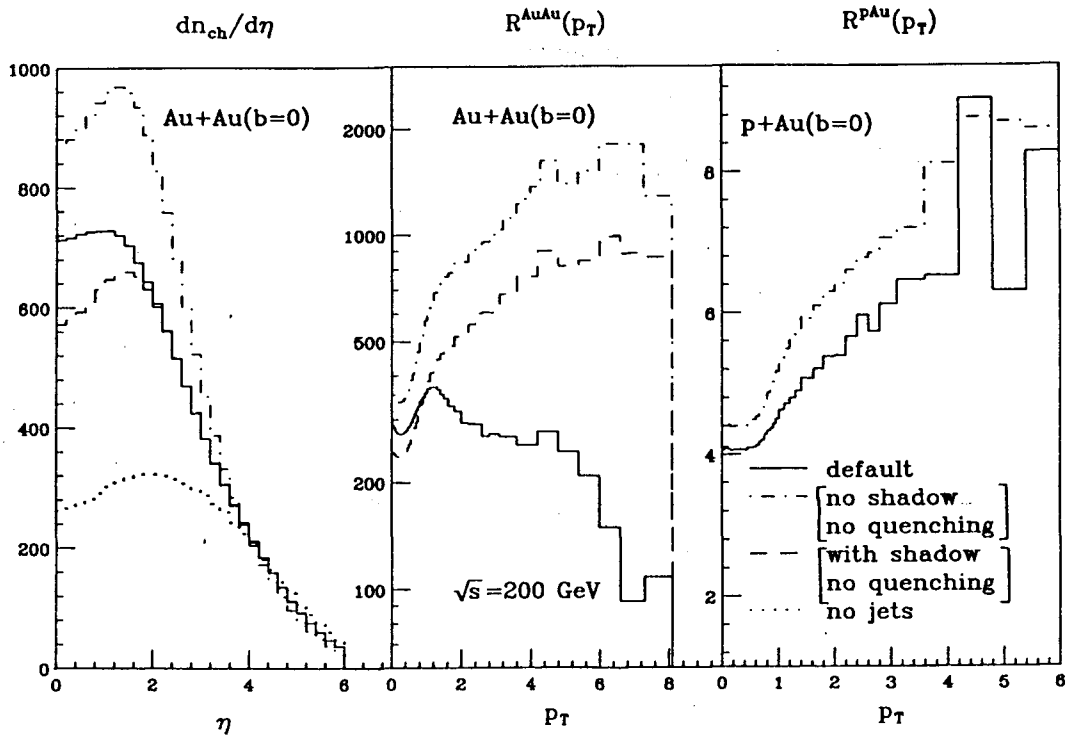

Figure 4: Dependence of the inclusive charged hadron spectra in central $Au + Au$ and $p + Au$ on minijet production (dashed-dot), gluon shadowing (dashed), and jet quenching (solid) assuming that gluon shadowing is identical to that of quarks and $\kappa = 2$ GeV/fm with $\lambda_s = 1$ fm. All curves are obtained using HIJING[11]. $R^{AB}(p_T)$ is the ratio of the number of charged hadrons with p_T produced in $A + B$ reactions to the number in $p + p$.

to $dN_{AA}/d\eta \approx AdN_{pp}/d\eta$, as expected. Without nuclear shadowing, $a(x) = 0$, (dashed-dot) the minijets are found to almost triple the rapidity density. However, if gluon shadowing is of the same magnitude as quark shadowing, then the minijet contribution is cut in half (dashed). The full line shows the small effect of high p_T jet quenching with $\kappa = 2$ GeV/fm and $\lambda_s = 1$ fm.

The large effects due to quenching only show up at high p_T. In particular, the ratio, $R^{AA}(p_T) = dN^{AA}(p_T)/dN^{pp}(p_T)$ in the middle panel shows that for $p_T > 3$ GeV both shadowing and quenching reduce dramatically the single hadron yields. In the absence of shadowing and quenching that ratio (dash-dot) approaches the number of binary collisions in the reaction, $T_{AA}(0)\sigma_{in} \approx 1200$, because of eq.(2). Shadowing alone (dashed) suppresses moderate p_T hadrons by a factor of 2. Inclusion of jet quenching with $\kappa = 2$ GeV/fm reduce that yield by a larger factor of $\sim 3 - 5$. It is remarkable that the *single* inclusive hadron spectrum is already so sensitive to these novel nuclear physics effects. Studies of higher order correlations, e.g. high p_T back-to-back hadron

pairs and exclusive jet pair analyses are found to be of course even more sensitive[11].

On the other hand, Fig.4 shows that AA data alone would not be sufficient to disentangle the two effects. To disentangle shadowing from quenching, obviously pA reactions must be studied at the same energy. In this case the density of comoving matter is so low that jet quenching should be negligible and the suppression of moderate p_T hadrons would be mainly due to gluon shadowing. The right panel in Fig. 4 shows the ratio of the single hadron p_T yield in central $p + Au$ to $p + p$ as a function of p_T. The form of the dash-dot curve is easily understood. In central $p + Au$ collisions about 8 nucleons from the Au participate. The number of binary collisions is then 8, which is the limit reached for $p_T \gtrsim 4$ GeV. The low $p_T < 1$ GeV ratio is, on the other hand, controlled by the number of pairs of beam $qq - q$ jets, which in this case is 4. We see from Fig. 4 that there is considerable sensitivity to gluon shadowing in $p + A$ and thus a systematic study of the inclusive hadron spectra should help constrain the magnitude of that interesting phenomenon even if the conventional direct photon techniques are not feasible at RHIC. With gluon shadowing so determined, we can return to the AA data and subtract theoretically the contribution due to shadowing and study systematically jet quenching as a function of the initial density of the comoving matter.

In summary, we have shown that high p_T probes can be expected to reveal interesting new physics in nuclear collisions at collider energies and complement the systematic study of other (low p_T) observables in the search for the QGP transition.

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