Discovery of Jet Quenching and Beyond

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Recent observation of high-p_T hadron spectra suppression and mono-jet production in central $Au + Au$ collisions and their absence in $d + Au$ collisions at RHIC have confirmed the long predicted phenomenon of jet quenching in high-energy heavy-ion collisions. Detailed analyses of the experimental data also show parton energy loss as the mechanism for the discovered jet quenching. Using a pQCD parton model that incorporates medium modified parton fragmentation functions and comparing to experimental data from deeply inelastic scattering off nuclei, one can conclude that the initial gluon (energy) density of the hot matter produced in central $Au+Au$ collisions that causes jet quenching at RHIC is about 30 (100) times higher than in a cold Au nucleus. Combined with data on bulk and collective properties of the hot matter, the observed jet quenching provides strong evidence for the formation of a strongly interacting quark gluon plasma in central $Au + Au$ collisions at RHIC.

I. INTRODUCTION

In the last three years, since the physics operation of the Relativistic Heavy-ion Collider (RHIC) at Brookhaven National Laboratory (BNL), one has witnessed tremendous progress in heavy-ion experimental physics and many milestones in the search for the elusive quark gluon plasma (QGP) that is expected to be formed in high-energy heavy-ion collisions. One of the most exciting phenomena observed at RHIC is jet quenching, or suppression of leading hadrons from fragmentation of hard partons due to their strong interaction with the dense medium. Such a phenomenon was long predicted by a pQCD-based model calculation [1], but was not seen in high-energy heavy-ion collisions until the first measurements in central $Au+Au$ collisions at $\sqrt{s} = 130$ GeV at RHIC [2,3]. The azimuthal distribution of high p_T hadrons was also found to display large anisotropy with respect to the reaction planes [4] of non-central $Au + Au$ collisions which was expected to be caused by the path-length dependence of jet quenching [5,6].

The measurements were confirmed and further extended to a larger p_T range in $Au + Au$ collisions at $\sqrt{s} = 200$ GeV, and the suppression was found to have a weak p_T-dependence at $p_T > 6$ GeV/c [7,8], independent of hadron species. At the same time, the back-side high-p_T two-hadron correlation in azimuthal angle, characteristic of high-p_T back-to-back jets in $p + p$ collisions, was found to disappear in central $Au+Au$ collisions [9], confirming the predicted mono-jet phenomenon of jet quenching [10].

The final milestone in the experimental discovery of jet quenching was achieved during the third year of the RHIC physics program, when both the single-hadron spectra [11–13] and the disappearance of away-side two-hadron correlation [12] were found to be absent in the same central rapidity region of $d + Au$ collisions at $\sqrt{s} = 200$ GeV. These $d + Au$ results prove that the observed high-p_T suppression patterns in $Au+Au$ collisions are not initial state effects encoded in the wavefunction of a beam nucleus, but are jet quenching caused by final state interaction of hard partons with the produced dense medium.

Following these three major experimental results, additional data—such as the dependence of away-side suppression on the azimuthal angle relative to the reaction plane [14], modification of the hadron distributions (fragmentation functions) both along and opposite the direction of the triggered high-p_T hadron [15], and absence of suppression in the direct photon spectra [16] in $Au + Au$ collisions—have now solidified the conclusion that the observed jet quenching is caused by parton multiple scattering and energy loss in the hot and dense medium. Furthermore, the splitting of baryon and meson spectra suppression and azimuthal anisotropy in the intermediate $p_T < 6$ GeV/c region [17–19] also point to nontrivial medium modification of hadronization that is indicative of dense partonic matter. Such a wealth of data affords one a quantitative phenomenological analysis of jet quenching and a tomographical picture of the hot and dense matter formed in heavy-ion collisions at RHIC.

In this report, we will first briefly review the recent progress in pQCD study of parton multiple scattering and induced radiative energy loss in a dense medium. Combining this with the pQCD parton model of high-p_T jet and hadron production, one can analyze the observed jet quenching phenomena to extract properties of the dense matter produced. One can combine the deduced properties with other information from analyses of bulk particle spectra such as collective flow and total multiplicity and energy production to present a collection of compelling evidence for the existence of a strongly interacting quark gluon plasma in $Au + Au$ collisions at RHIC. However, we will focus in this report only on the analysis of jet quenching data and the properties of the dense matter one can extract from these data.

II. PARTON ENERGY LOSS

One important step in establishing evidence of QGP formation is to characterize the properties of the produced dense medium, for example, the parton and energy density and current-current correlations, among many other characteristics. Traditionally, one can study the properties of a medium via scattering experiments with particles. In deeply inelastic scattering (DIS) experiments, for example, leptons scatter off the nucleon medium via photon exchange. The response function or the correlation function of the electromagnetic currents $j_\mu^{em}(x) = \sum_q e_q \bar{\psi}_q(x)\gamma_\mu\psi_q(x)$,

$$W_{\mu\nu}(q) = \frac{1}{4\pi}\int d^4x e^{iq\cdot x}\langle A \mid j_\mu^{em}(0) j_\nu^{em}(x) \mid A\rangle, \quad (1)$$

is a direct measurement of the quark distributions in a nucleon or nucleus. For a dynamic system in heavy-ion collisions, one can no longer use the technique of scattering with an external beam of particles in the conventional sense because of the transient nature of the matter. The lifetime of the system is very short, on the order of a few fm/c. The initial spatial size is only the size of the colliding nuclei, about 6 fm in diameter in the transverse dimension for the heaviest nuclei. The system expands very rapidly both in the longitudinal and the transverse direction. These characteristics make it impossible to use external probes to study the properties of the produced dense matter in high-energy heavy-ion collisions. One has to resort to internal probes such as the electromagnetic emission from the dynamic medium, whose rate is given by the thermal average of the above correlation function. The emission rate depends mainly on the local temperature or the parton density while the total yield also depends on the whole evolution history of the system. Therefore, a strongly interacting system can reveal its properties and dynamics through photon and dilepton emission. The same medium interaction should also cause attenuation of fast and energetic partons propagating through the medium. Such an effect is the underlying physics of the jet quenching [20,21] phenomenon and jet tomography technique for studying properties of dense matter in high-energy heavy-ion collisions.

Jet quenching as a probe of the dense matter in heavy-ion collisions takes advantage of the hard processes of jet production in high-energy heavy-ion collisions. Because large-p_T partons are produced very early in heavy-ion collisions, they can probe the early stage of the formed dense medium. Similar to the technology of computed tomography (CT), the study of these energetic particles, their initial production and interaction with the dense medium can yield critical information about the properties of the medium that is otherwise difficult to access through soft hadrons from the hadronization of the bulk matter. Though relatively rare with small cross sections, the jet production rate can be calculated perturbatively in QCD and agrees well with experimental measurements in high-energy $p+p(\bar{p})$ collisions. A critical component of the jet tomography is then to understand the jet's attenuation through dense matter as it propagates through the medium.

There have been many theoretical studies [22–24] of jet attenuation in a hot medium in recent years. The first attempt was by Bjorken [20] to calculate elastic energy loss of a parton via elastic scattering in the hot medium. A simple estimate can be given by the thermal averaged energy transfer $\nu_{\rm el} \approx q_\perp^2/2\omega$ of the jet parton to a thermal parton, with energy ω, and $q_\perp$ the transverse momentum transfer of the elastic scattering. The resultant elastic energy loss [25]

$$\frac{dE_{\rm el}}{dx} = C_2 \frac{3\pi\alpha_{\rm s}^2}{2} T^2 \ln\left(\frac{3ET}{2\mu^2}\right) \quad (2)$$

is sensitive to the temperature of the thermal medium but is small compared to radiative energy loss. Here, μ is the Debye screening mass and C_2 is the Casimir coefficient of the propagating parton in its fundamental presentation. The elastic energy loss can also be calculated within finite temperature field theory of QCD [26] with a more careful and consistent treatment of the Debye screening effect.

Though there had been estimates of the radiative parton energy loss using the uncertainty principle [27], the first theoretical study of QCD radiative parton energy loss incorporating Landau-Pomeranchuk-Migdal interference effect [28] is by Gyulassy and Wang [29] where multiple parton scattering is modeled by a screened Coulomb potential model. Baier *et al.* (BDMPS) [30] later considered the effect of gluon rescattering which turned out to be very important for gluon radiation induced by multiple scattering in a dense medium. These two studies have ushered in many recent works on the subject, including a path integral approach to the problem [31], twist [32] or opacity expansion framework [33,34] which is more suitable for multiple parton scattering in a finite system. The radiative parton energy loss in the leading order of the twist or opacity expansion was found to have a simple form [33]

$$\frac{dE_{\rm rad}}{dx} \approx C_2 \frac{\alpha_{\rm s}\mu^2}{4}\frac{L}{\lambda}\ln\left(\frac{2E}{\mu^2 L}\right), \quad (3)$$

for an energetic parton ($E \gg \mu$) in a static medium, where λ is the gluon's mean free path in the medium. The unique L-dependence of the parton energy loss, which was first discovered by BDMPS [30] is a consequence of the non-Abelian LMP interference effect in a QCD medium. In a dynamic system the total energy loss was found [6,35,36] to be proportional to a line integral of the gluon density along the parton propagation path.

Since gluons are bosons, there should also be stimulated gluon emission and absorption by the propagating parton in the presence of thermal gluons in the hot medium. Such detailed balance is crucial for parton thermalization and should also be important for calculating the energy loss of an energetic parton in a hot medium. Taking into account such detailed balance in gluon emission, Wang and Wang [37] obtained an asymptotic behavior of the effective energy loss in the opacity expansion framework ,

$$\frac{\Delta E}{E} \approx \frac{\alpha_s C_2 \mu^2 L^2}{4\lambda E}\left[\ln\frac{2E}{\mu^2 L} - 0.048\right] - \frac{\pi\alpha_s C_F}{3}\frac{LT^2}{\lambda E^2}\left[\ln\frac{\mu^2 L}{T} - 1 + \gamma_{\rm E} - \frac{6\zeta'(2)}{\pi^2}\right], \quad (4)$$

where the first term is from the induced bremsstrahlung and the second is due to gluon absorption in detailed balance, which effectively reduces the total parton energy loss in the medium. Though the effect of detailed balance is small at large energy ($E \gg \mu$), it changes the effective energy dependence of the energy loss in the intermediate energy region.

III. MODIFIED FRAGMENTATION FUNCTION

Because of color confinement in the vacuum, one can never separate hadrons fragmenting from the leading parton and particles materializing from the radiated gluons. The total energy in the conventionally defined jet cone in principle should not change due to induced radiation, assuming that most of the energy carried by radiative gluons remains inside the jet cone [38]. Additional rescattering of the emitted gluon with the medium could broaden the jet cone significantly, thus reducing the energy in a fixed cone. However, fluctuation of the underlying background in high-energy heavy-ion collisions makes it very difficult, if not impossible, to determine the energy of a jet on an event-by-event base with sufficient precision to discern a finite energy loss of the order of 10 GeV. Since high-p_T hadrons in hadronic and nuclear collisions come from fragmentation of high-p_T jets, energy loss naturally leads to suppression of high-p_T hadron spectra. This was why Gyulassy and Wang proposed [1] to measure the suppression of high-p_T hadrons to study parton energy loss in heavy-ion collisions.

Since parton energy loss effectively slows down the leading parton in a jet, a direct manifestation of jet quenching is the modification of the jet fragmentation function, $D_{a\to h}(z,\mu^2)$, which can be measured directly in events in which one can identify the jet via a companion particle like a direct photon [39] or a triggered high p_T hadron. This modification can be directly translated into the energy loss of the leading parton. Since inclusive hadron spectra are a convolution of the jet production cross section and the jet fragmentation function in pQCD, the suppression of inclusive high-p_T hadron spectra is a direct consequence of the medium modification of the jet fragmentation function caused by parton energy loss.

Since a jet parton is always produced via a hard process involving a large momentum scale, it should also have final state radiation with and without rescattering, leading to the DGLAP evolution of fragmentation functions. Such final state radiation effectively acts as a self-quenching mechanism, softening the leading hadron distribution. This process is quite similar to the induced gluon radiation and the two should have a strong interference effect. It is therefore natural to study jet quenching and modified fragmentation functions in the framework of modified DGLAP evolution equations in a medium [32].

To demonstrate medium modified fragmentation function and parton energy loss, we review here the twist expansion approach in the study of deeply inelastic scattering (DIS) eA [32,40]. While the results can be readily extended to the case of a hot medium, they also provide a baseline for comparison to parton energy loss in cold nuclei.

In the parton model with the collinear factorization approximation, the leading-twist contribution to the semi-inclusive cross section can be factorized into a product of parton distributions, parton fragmentation functions and the hard partonic cross section. Including all leading log radiative corrections, the lowest order contribution from a single hard $\gamma^* + q$ scattering can be written as

$$\frac{dW^S_{\mu\nu}}{dz_h} = \sum_q e_q^2 \int dx f_q^A(x,\mu_I^2) H^{(0)}_{\mu\nu}(x,p,q) D_{q\to h}(z_h,\mu^2)\,,$$

where $H^{(0)}_{\mu\nu}(x,p,q)$ is the hard part of the process in leading order, ℓ_h the observed hadron momentum, $p = [p^+, 0, \mathbf{0}_\perp]$ the momentum per nucleon in the nucleus, $q = [-Q^2/2q^-, q^-, \mathbf{0}_\perp]$ the momentum transfer carried by the virtual photon. In the chosen frame, q^- is the quark momentum transferred from the virtual photon. The momentum fraction carried by the hadron is defined as $z_h = \ell_h^-/q^-$ and $x = x_B = Q^2/2p^+q^-$ is the Bjorken variable. μ_I^2 and μ^2 are the factorization scales for the initial quark distributions $f_q^A(x,\mu_I^2)$ in a nucleus and the fragmentation functions $D_{q\to h}(z_h,\mu^2)$, respectively.

In a nuclear medium, the propagating quark in DIS will experience additional scattering with other partons from the nucleus. The rescatterings may induce additional gluon radiation and cause the leading quark to lose energy. Such induced gluon radiations will effectively give rise to additional terms in a DGLAP-like evolution equation leading to the modification of the fragmentation functions in a medium. These are the so-called higher-twist corrections since they involve higher-twist parton matrix elements and are power-suppressed. The leading contributions involve two-parton correlations from two different nucleons inside the nucleus.

One can apply the generalized factorization [41] to these multiple scattering processes. In this approximation, the double scattering contribution to radiative correction can be calculated and the effective modified fragmentation function is [32]

$$\widetilde{D}_{q\to h}(z_h,\mu^2) \equiv D_{q\to h}(z_h,\mu^2) + \int_0^{\mu^2} \frac{d\ell_T^2}{\ell_T^2}\frac{\alpha_s}{2\pi}\int_{z_h}^1 \frac{dz}{z} \Big[\Delta\gamma_{q\to qg}(z,x,x_L,\ell_T^2)\, D_{q\to h}(z_h/z) + \Delta\gamma_{q\to gq}(z,x,x_L,\ell_T^2) D_{g\to h}(z_h/z)\Big] \ , \quad (5)$$

in much the same form as the DGLAP correction in vacuum, where $D_{q\to h}(z_h,\mu^2)$ and $D_{g\to h}(z_h,\mu^2)$ are the leading-twist quark and gluon fragmentation functions. The modified splitting functions are given as

$$\Delta\gamma_{q\to qg}(z,x,x_L,\ell_T^2) = \left[\frac{1+z^2}{(1-z)_+}T^A_{qg}(x,x_L) + \delta(1-z)\Delta T^A_{qg}(x,\ell_T^2)\right] \frac{2\pi\alpha_s C_A}{\ell_T^2 N_c f_q^A(x,\mu_I^2)} \ , \quad (6)$$

$$\Delta\gamma_{q\to gq}(z,x,x_L,\ell_T^2) = \Delta\gamma_{q\to qg}(1-z,x,x_L,\ell_T^2). \quad (7)$$

Here, the fractional momentum $x = x_B$ is carried by the initial quark and $x_L = \ell_T^2/2p^+q^-z(1-z)$ is the additional momentum fraction carried either by the quark or gluon in the secondary scattering that is required for induced gluon radiation. The twist-four parton matrix element of the nucleus,

$$T^A_{qg}(x,x_L) = \int \frac{dy^-}{2\pi}\, dy_1^- dy_2^- e^{i(x+x_L)p^+y^-}(1-e^{-ix_Lp^+y_2^-}) \times(1-e^{-ix_Lp^+(y^- - y_1^-)})\theta(-y_2^-)\theta(y^- - y_1^-) \times\frac{1}{2}\langle A|\bar{\psi}_q(0)\,\gamma^+\, F_\sigma^{\ +}(y_2^-)\, F^{+\sigma}(y_1^-)\,\psi_q(y^-)|A\rangle \ , \quad (8)$$

has a dipole-like structure which is a result of LPM interference in gluon bremsstrahlung.

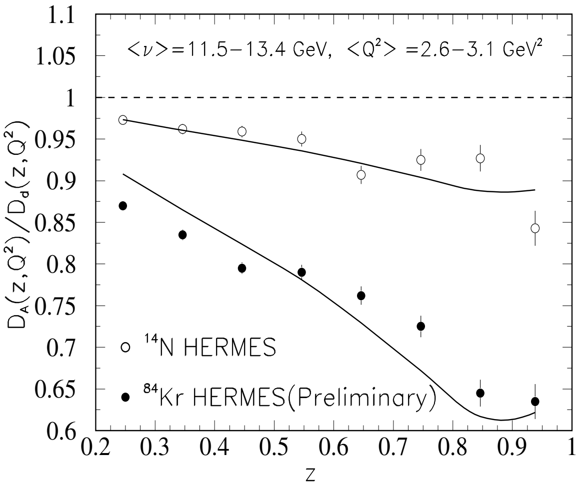

FIG. 1. Predicted nuclear modification of the jet fragmentation functions compared to the HERMES data [43] on ratios of hadron distributions between A and d targets in DIS.

Averaged over a Gaussian nuclear distribution, the interference will produce a factor $1 - e^{-x_L^2/x_A^2}$ with $x_A = 1/MR_A$. Here R_A is the nuclear size and M is the nucleon mass. Using the factorization approximation [32,41,42], one can relate the twist-four parton matrix elements of the nucleus to the twist-two parton distributions of nucleons and the nucleus,

$$T^A_{qg}(x,x_L) \approx \frac{\widetilde{C}}{x_A}(1-e^{-x_L^2/x_A^2})f_q^A(x), \quad (9)$$

where $\widetilde{C} \equiv 2Cx_T f_g^N(x_T)$ is considered a constant. One can identify $1/x_Lp^+ = 2q^-z(1-z)/\ell_T^2$ as the formation time of the emitted gluons. In the limit of collinear radiation ($x_L \to 0$) or when the formation time of the gluon radiation is much larger than the nuclear size, the above matrix element vanishes, demonstrating a typical LPM interference effect.

Since the LPM interference suppresses gluon radiation whose formation time ($\tau_f \sim Q^2/\ell_T^2p^+$) is larger than the nuclear size MR_A/p^+ (here in the infinite momentum frame), ℓ_T^2 should then have a minimum value of $\ell_T^2 \sim Q^2/MR_A \sim Q^2/A^{1/3}$. Therefore, the leading higher-twist contribution is proportional to $\alpha_s R_A/\ell_T^2 \sim \alpha_s R_A^2/Q^2$ due to double scattering and depends quadratically on the nuclear size R_A.

With the assumption of the factorized form of the twist-4 nuclear parton matrices, there is only one free parameter $\widetilde{C}(Q^2)$ which represents quark-gluon correlation strength inside nuclei. Once it is fixed, one can predict the z, energy and nuclear dependence of the medium modification of the fragmentation function. Shown in Fig. 1 are the calculated [35] nuclear modification factor of the fragmentation functions for ^{14}N and ^{84}Kr targets as compared to the recent HERMES data [43]. The predicted shape of the z-dependence agrees well with the experimental data. A remarkable feature of the prediction is the quadratic $A^{2/3}$ nuclear size dependence, which is verified for the first time by an experiment. By fitting the overall suppression for one nuclear target, one obtains the only parameter in the calculation, $\widetilde{C}(Q^2) = 0.0060$ GeV^2 with $\alpha_s(Q^2) = 0.33$ at $Q^2 \approx 3$ GeV^2. The predicted energy (ν) dependence also agrees with the experimental data as shown in Fig. 2. Note that the suppression goes away as the quark energy increases. This is in sharp contrast to the jet quenching in heavy-ion collisions.

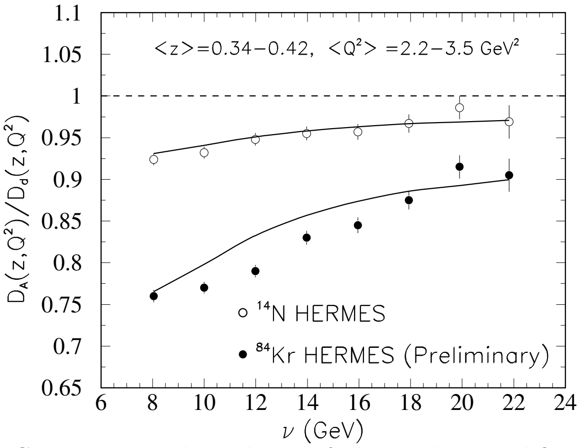

FIG. 2. Energy dependence of the nuclear modification compared with the HERMES data [43].

One can quantify the modification of the fragmentation by the quark energy loss which is defined as the momentum fraction carried by the radiated gluon,

$$\begin{aligned}\langle\Delta z_g\rangle &= \int_0^{\mu^2}\frac{d\ell_T^2}{\ell_T^2}\int_0^1 dz\frac{\alpha_s}{2\pi}z\,\Delta\gamma_{q\to gq}(z,x_B,x_L,\ell_T^2)\\ &=\widetilde{C}\frac{C_A\alpha_s^2}{N_c}\frac{x_B}{x_AQ^2}\int_0^1 dz\frac{1+(1-z)^2}{z(1-z)}\\ &\times\int_0^{x_\mu}\frac{dx_L}{x_L^2}(1-e^{-x_L^2/x_A^2}),\end{aligned} \tag{10}$$

where $x_\mu=\mu^2/2p^+q^-z(1-z)=x_B/z(1-z)$ if the factorization scale is set as $\mu^2=Q^2$. When $x_A\ll x_B\ll 1$ the leading quark energy loss is approximately

$$\langle\Delta z_g\rangle(x_B,\mu^2)\approx\widetilde{C}\frac{C_A\alpha_s^2}{N_c}\frac{x_B}{Q^2x_A^2}6\sqrt{\pi}\ln\frac{1}{2x_B}\,. \tag{11}$$

Since $x_A=1/MR_A$, the energy loss $\langle\Delta z_g\rangle$ thus depends quadratically on the nuclear size.

In the rest frame of the nucleus, $p^+=M$, $q^-=\nu$, and $x_B\equiv Q^2/2p^+q^-=Q^2/2M\nu$. One can get the averaged total energy loss as $\Delta E=\nu\langle\Delta z_g\rangle\approx\widetilde{C}(Q^2)\alpha_s^2(Q^2)MR_A^2(C_A/N_c)3\ln(1/2x_B)$. With the determined value of $\widetilde{C}$, $\langle x_B\rangle\approx 0.124$ in the HERMES experiment [43] and the average distance $\langle L_A\rangle=R_A\sqrt{2/\pi}$ for the assumed Gaussian nuclear distribution, one gets the quark energy loss $dE/dL\approx 0.5$ GeV/fm inside a Au nucleus.

IV. JET QUENCHING IN HOT MEDIUM

To extend the study of modified fragmentation functions to jets in heavy-ion collisions, one can assume $\langle k_T^2\rangle\approx\mu^2$ (the Debye screening mass) and a gluon density profile $\rho(y)=(\tau_0/\tau)\theta(R_A-y)\rho_0$ for a 1-dimensional expanding system. Since the initial jet production rate is independent of the final gluon density, which can be related to the parton-gluon scattering cross section [44] $[\alpha_s x_T G(x_T)\sim\mu^2\sigma_g]$, one has then

$$\frac{\alpha_s T_{qg}^A(x_B,x_L)}{f_q^A(x_B)}\sim\mu^2\int dy\sigma_g\rho(y)[1-\cos(y/\tau_f)], \tag{12}$$

where $\tau_f=2Ez(1-z)/\ell_T^2$ is the gluon formation time. The averaged fractional energy loss is then,

$$\begin{aligned}\langle\Delta z_g\rangle=\frac{C_A\alpha_s}{\pi}\int_0^1 dz\int_0^{Q^2/\mu^2}du\frac{1+(1-z)^2}{u(1+u)}\\ \times\int_{\tau_0}^{R_A}d\tau\sigma_g\rho(\tau)\left[1-\cos\left(\frac{(\tau-\tau_0)\,u\,\mu^2}{2Ez(1-z)}\right)\right].\end{aligned} \tag{13}$$

Keeping only the dominant contribution and assuming $\sigma_g\approx C_a 2\pi\alpha_s^2/\mu^2$ (C_a=1 for qg and 9/4 for gg scattering), one obtains the total energy loss,

$$\langle\Delta E\rangle\approx\frac{\pi C_a C_A\alpha_s^3}{R_A}\int_{\tau_0}^{R_A}d\tau\rho(\tau)(\tau-\tau_0)\ln\frac{2E}{\tau\mu^2}. \tag{14}$$

which has also be obtained in the opacity expansion approach [6,36] in a thin plasma.

In a static medium, the gluon density is independent of time and one can recover the quadratic lenghth dependence in Eq. (3). Neglecting the logarithmic dependence on τ, the averaged energy loss can be expressed as

$$\langle\frac{dE}{dL}\rangle_{1d}\approx(dE_0/dL)(2\tau_0/R_A) \tag{15}$$

in a 1-dimensional expanding system, $\rho(\tau)=\rho_0\tau_0/\tau$. Here $dE_0/dL\propto\rho_0R_A$ is the energy loss in a static medium with the same gluon density ρ_0 as in the 1-d expanding system at time τ_0. Because of the expansion, the averaged energy loss $\langle dE/dL\rangle_{1d}$ is suppressed as compared to the static case and does not depend linearly on the system size at a fixed value of initial gluon density.

One should emphasize now that the above simple logarithmic energy dependence of the parton energy loss in Eqs. (11) and (14) is only an asymptotic behavior. However, kinematic limits on induced gluon radiation from a parton with finite energy can result in a stronger energy dependence [33]. The effect of detailed balance with thermal absorption further increases the energy dependence [37], while reducing the effective parton energy loss for intermediate values of parton energy. Such a detailed balance effect sets the energy dependence of the energy loss in cold nuclei in DIS apart from the hot medium in heavy-ion collisions. If one parameterizes the energy dependence of the energy loss including the full kinematic limits and thermal absorption, one would get

$$\langle\frac{dE}{dL}\rangle_{1d}=\epsilon_0(E/\mu-1.6)^{1.2}/(7.5+E/\mu). \tag{16}$$

The threshold is the consequence of gluon absorption that competes with radiation and effectively shuts off the energy loss. The parameter μ is set to be 1.5 GeV in the following discussions. Such a detailed balance effect can also explain why the total hadron multiplicity that is dominated by soft hadrons does have significant enhancement over the $p+p$ collisions as result of induced gluon emission.

A. Single Spectra

To calculate the modified high-p_T spectra in $A+A$ collisions, we use a LO pQCD model [45],

$$\begin{aligned}\frac{d\sigma_{AA}^h}{dyd^2p_T}&=K\sum_{abcd}\int d^2bd^2rdx_adx_bd^2k_{aT}d^2k_{bT}\\ &\times t_A(r)t_A(|\mathbf{b}-\mathbf{r}|)g_A(k_{aT},r)g_A(k_{bT},|\mathbf{b}-\mathbf{r}|)\\ &\times f_{a/A}(x_a,Q^2,r)f_{b/A}(x_b,Q^2,|\mathbf{b}-\mathbf{r}|)\\ &\times\frac{D'_{h/c}(z_c,Q^2,\Delta E_c)}{\pi z_c}\frac{d\sigma}{d\hat{t}}(ab\to cd),\end{aligned} \tag{17}$$

with medium modified fragmentation functions $D'_{h/c}$. Here, $z_c = p_T/p_{Tc}$, $y = y_c$, $\sigma(ab \to cd)$ are elementary parton scattering cross sections and $t_A(b)$ is the nuclear thickness function normalized to $\int d^2b t_A(b) = A$. The $K \approx 1.5$–2 factor is used to account for higher order pQCD corrections.

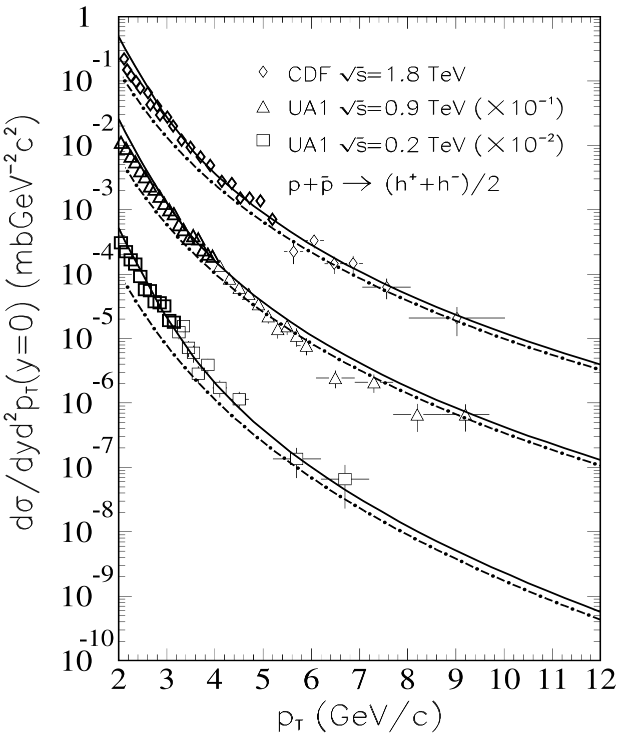

FIG. 3. Single-inclusive spectra of charged hadrons in $p+\bar{p}$ collisions at $\sqrt{s} = 200, 900, 1800$ GeV. The solid (dot-dashed) lines are parton model calculations with (without) intrinsic k_T. Experimental data are from Refs. [50,51].

To simplify the incorporation of the medium modification of the fragmentation function in the parton model, an effective model [39] can be used:

$$D'_{h/c}(z_c, Q^2, \Delta E_c) = e^{-\langle\frac{\Delta L}{\lambda}\rangle} D^0_{h/c}(z_c, Q^2) + (1 - e^{-\langle\frac{\Delta L}{\lambda}\rangle}) \times \left[\frac{z'_c}{z_c} D^0_{h/c}(z'_c, Q^2) + \langle\frac{\Delta L}{\lambda}\rangle \frac{z'_g}{z_c} D^0_{h/g}(z'_g, Q^2)\right] \quad (18)$$

where z'_c, z'_g are the rescaled momentum fractions after parton energy loss. This effective model is found [46] to reproduce the pQCD result from Eq.(5) very well, but only when $\Delta z = \Delta E_c/E$ is set to be $\Delta z \approx 0.6\langle z_g\rangle$. Therefore the actual averaged parton energy loss should be $\Delta E/E = 1.6\Delta z$, with Δz extracted from the effective model. The factor 1.6 is mainly caused by the unitarity correction effect in the pQCD calculation. The fragmentation functions in free space $D^0_{h/c}(z_c, Q^2)$ will be given by the BBK parameterization [47].

The parton distributions per nucleon $f_{a/A}(x_a, Q^2, r)$ inside the nucleus are assumed to be factorizable into the parton distributions in a free nucleon given by the MRSD$-'$ parameterization [48] and the impact-parameter dependent nuclear modification factor given by the new HIJING parameterization [49]. The initial transverse momentum distribution $g_A(k_T, Q^2, b)$ is assumed to have a Gaussian form with a width that includes both an intrinsic part in a nucleon and nuclear broadening. This parton model can describe high-p_T hadron spectra in $p + p(\bar{p})$ well, as shown in Fig. 3 that compares the parton model calculation [45] with the experimental data on the inclusive charged hadron spectra at different collider energies.

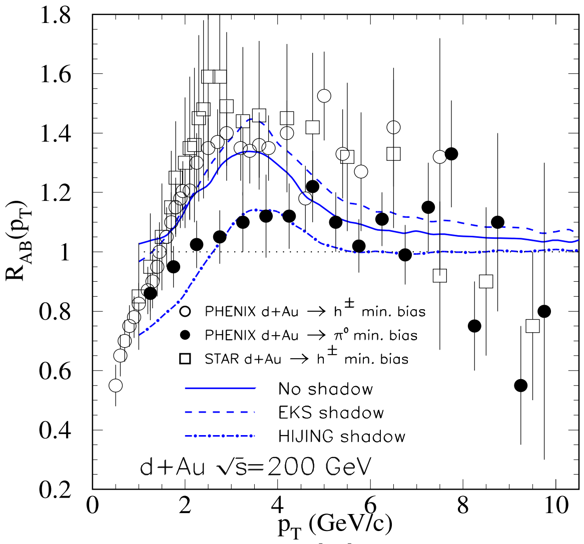

FIG. 4. The first prediction [45] of the Cronin effect in $p + Au$ collisions at $\sqrt{s} = 200$ GeV compared to the recent RHIC data from PHENIX [11] and STAR [12]. The dashed and dot-dashed lines correspond to HIJING [49] and EKS [52] parameterizations of nuclear parton distributions.

In $p + A$ collisions, this parton model incorporates both nuclear modification of the parton distributions and broadening of the initial transverse momentum. The initial momentum broadening leads to an enhancement, known as the Cronin effect, of the hadron spectra in the intermediate p_T region. The enhancement disappears at large p_T. The parameters have been fitted to the nuclear modification of the p_T spectra in $p + A$ collisions at up to the Fermilab energy $\sqrt{s} = 40$ GeV. Shown in Fig. 4 is the first prediction made in 1998 [45] of the Cronin effect at RHIC for $p + Au$ collisions at $\sqrt{s} = 200$ GeV as compared to the RHIC data on $d + Au$ collisions. As one can see, the initial multiple scattering in nuclei can give some moderate Cronin enhancement of the high-p_T spectra, though the details depend on the parameterization of the nuclear modification of parton distributions. Therefore, any observed suppression of the high-p_T spectra in $Au + Au$ collisions has to be caused by jet quenching.

In $A + A$ collisions, one has to consider medium modification of the fragmentation functions due to parton energy loss according to Eq. (18). Assuming a 1-dimensional expanding medium with a gluon density

$\rho_g(\tau, r)$ that is proportional to the transverse profile of participant nucleons, one can calculate the impact-parameter dependence of the energy loss,

$$\Delta E(b, r, \phi) \approx \langle \frac{dE}{dL} \rangle_{1d} \int_{\tau_0}^{\Delta L} d\tau \frac{\tau - \tau_0}{\tau_0 \rho_0} \rho_g(\tau, b, \vec{r} + \vec{n}\tau), \tag{19}$$

according to Eq. (14), where $\Delta L(b, \vec{r}, \phi)$ is the distance a jet, produced at $\vec{r}$, has to travel along $\vec{n}$ at an azimuthal angle ϕ relative to the reaction plane in a collision with impact-parameter b. Here, ρ_0 is the averaged initial gluon density at initial time τ_0 in a central collision and $\langle dE/dL \rangle_{1d}$ is the average parton energy loss over a distance R_A in a 1-d expanding medium with an initial uniform gluon density ρ_0. The corresponding energy loss in a static medium with a uniform gluon density ρ_0 over a distance R_A is [35] $dE_0/dL = (R_A/2\tau_0)\langle dE/dL \rangle_{1d}$. We will use the parameterization in Eq. (16) for the effective energy dependence of the quark energy loss.

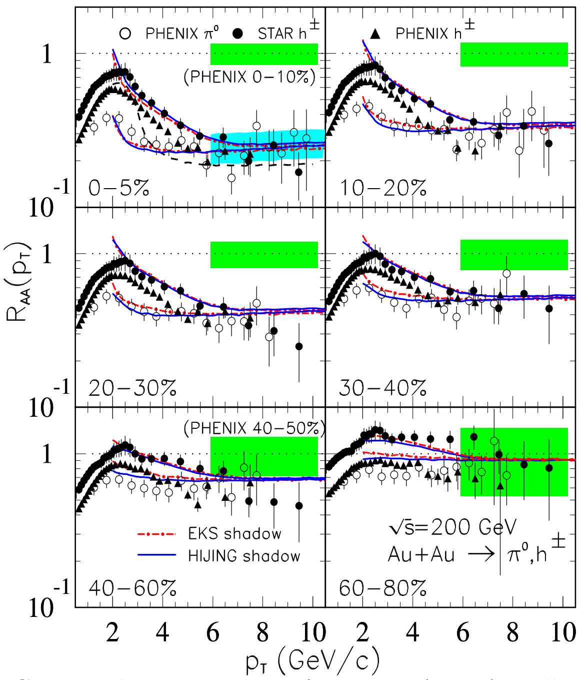

FIG. 5. Hadron suppression factors in $Au + Au$ collisions as compared to data from STAR [8] and PHENIX [7]. See text for a detailed explanation.

Shown in Fig. 5 are the calculated nuclear modification factors $R_{AB}(p_T) = d\sigma^h_{AB}/\langle N_{\rm binary}\rangle d\sigma^h_{pp}$ for hadron spectra ($|y| < 0.5$) in $Au + Au$ collisions at $\sqrt{s} = 200$ GeV, as compared to experimental data [8,7]. Here, $\langle N_{\rm binary}\rangle = \int d^2b d^2r t_A(r) t_A(|\vec{b} - \vec{r}|)$. To fit the observed π^0 suppression (solid lines) in the most central collisions, we have used $\mu = 1.5$ GeV, $\epsilon_0 = 1.07$ GeV/fm and $\lambda_0 = 1/(\sigma\rho_0) = 0.3$ fm in Eqs. (18) and (16). The hatched area (also in other figures in this report) indicates a variation of $\epsilon_0 = \pm 0.3$ GeV/fm. The hatched boxes around $R_{AB} = 1$ represent experimental errors of STAR data in overall normalization, including a constant factor of about 18% from the determination of total $p+p$ inelastic cross section. Nuclear k_T broadening and parton shadowing together give a slight enhancement of hadron spectra at intermediate $p_T = 2-4$ GeV/c without parton energy loss.

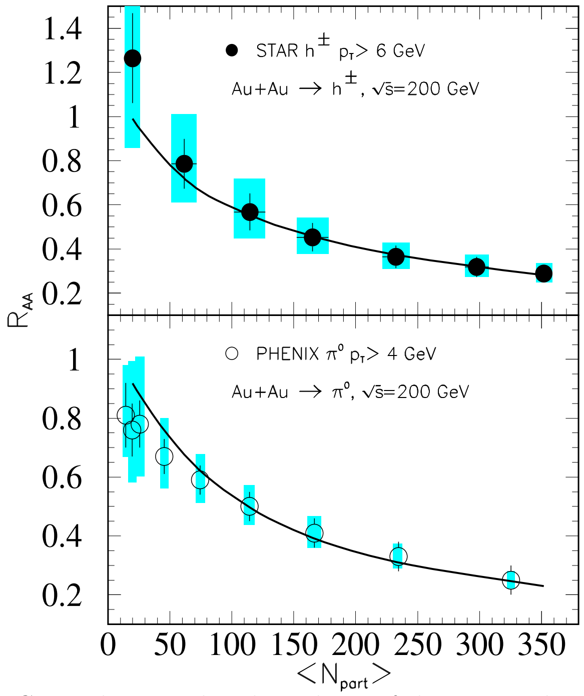

FIG. 6. The centrality dependence of the measured single inclusive hadron suppression [7,8] at high-p_T as compared to theoretical calculation with parton energy loss

The flat p_T dependence of the π^0 suppression is a consequence of the strong energy dependence of the parton energy loss as also noted in Refs. [53,54]. This is in sharp contrast to the energy dependence of the suppression in DIS as shown in Fig. 2 and should be a strong indication of the detailed balance effect in a thermal medium. Such an effect also causes the slight rise of R_{AB} at $p_T < 4$ GeV/c in the calculation. In this region, one expects the fragmentation picture to gradually lose its validity and to be taken over by other non-perturbative effects, especially for kaons and baryons. As a consequence, the $(K + p)/\pi$ ratio in central $Au + Au$ collisions is significantly larger than in peripheral $Au+Au$ or $p+p$ collisions [17]. This has been attributed to parton coalescence in the fragmentation of the leading parton [55–57]. The slight flavor dependence of the Cronin effect in $d + Au$ collisions as shown in Fig. 4 could also be attributed to the coalescence effect [59,58]. Such an effect, while providing interesting insight into the hadronization mechanism in the quark gluon plasma, will complicate the picture of jet quenching and introduce large uncertainties in the physics extracted from jet tomography analysis of the experimental data. Fortunately, the effect has been shown to disappear at large $p_T > 5$ GeV/c both in the coalescence models and the experimental data. The suppression ratio of charged hadrons and π^0 and of Λ and K all converge [19] at large p_T. In the calculation shown in Fig. 5, a nuclear-dependent (proportional to $\langle N_{\rm binary}\rangle$) soft component was added to kaon and baryon fragmentation functions to take into account the coalescence effect, so that $(K+p)/\pi \approx 2$ at $p_T \sim 3$ GeV/c in the most central $Au+Au$ collisions and it approaches its $p+p$ value at $p_T > 5$ GeV/c. To demonstrate the sensitivity to the parameterized parton energy loss in the intermediate p_T region, we also show $R_{AA}^{h^\pm}$ in 0-5% centrality (dashed line) for $\mu = 2.0$ GeV and $\epsilon_0 = 2.04$ GeV/fm without the soft component.

The measured centrality dependence of the single hadron suppression in $Au+Au$ collisions, shown in Fig. 6, agrees very well with parton model with energy loss. This is the consequence of the centrality dependence of the averaged total energy loss in Eq. (14), which leads to an effective surface emission of the surviving jets. Jets produced around the core of the overlapped region are strongly suppressed, since they lose the largest amount of energy. The centrality dependence of the suppression is found to be more dominated by the geometry of the produced dense matter than the length dependence of the parton energy loss.

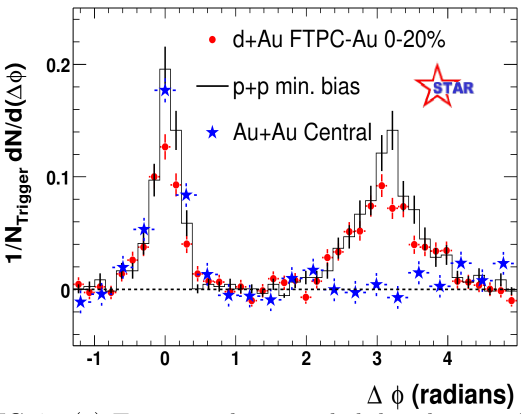

FIG. 7. (a) Two-particle azimuthal distributions for minimum bias and central $d+Au$ collisions, and for $p+p$ collisions [12]. (b) Comparison of two-particle azimuthal distributions for central d+Au collisions to those seen in p+p and central Au+Au collisions.

B. Di-hadron Spectra

Since jets are always produced in pairs in LO pQCD, two-hadron correlations in $p+p$ collisions should have unique characteristics pertaining to the back-to-back jet structure of the initial hard parton-parton scattering. Jet quenching should also modify this di-hadron correlation of the back-to-back jets. Shown in Fig. 7 are the two-hadron correlations in azimuthal angle as measured by STAR experiment in $p+p$, $d+Au$ and $Au+Au$ collisions at RHIC [12]. While the same-side correlation remains the same, the away-side correlation due to the back-side jet is strongly suppressed in central $Au+Au$ collisions. The energy loss that causes the suppression of the single inclusive hadron spectra should also be able to explain the disappearance of away-side correlation.

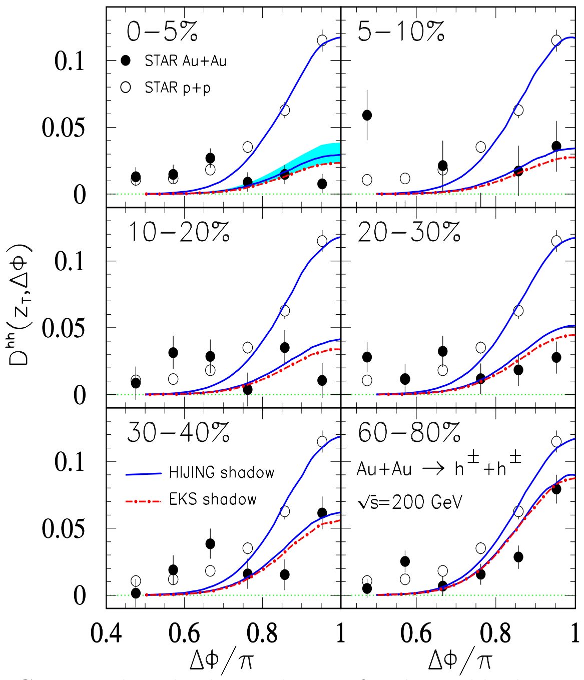

FIG. 8. Back-to-back correlations for charged hadrons with $p_T^{\rm trig} > p_T > 2$ GeV/c, $p_T^{\rm trig} = 4-6$ GeV/c and $|y| < 0.7$ in $Au+Au$ (lower curves) and $p+p$ (upper curves) collisions as compared to the STAR [9] data.

In the same LO pQCD parton model, one can calculate the spectra [60],

$$
\begin{aligned}
E_1E_2\frac{d\sigma_{AA}^{h_1h_2}}{d^3p_1d^3p_2} &= \frac{K}{2}\sum_{abcd}\int d^2bd^2rdx_adx_bd^2k_{aT}d^2k_{bT} \\
&\times t_A(r)t_A(|\mathbf{b}-\mathbf{r}|)g_A(k_{aT},r)g_A(k_{bT},|\mathbf{b}-\mathbf{r}|) \\
&\times f_{a/A}(x_a,Q^2,r)f_{b/A}(x_b,Q^2,|\mathbf{b}-\mathbf{r}|) \\
&\times D_{h/c}(z_c,Q^2,\Delta E_c)D_{h/d}(z_d,Q^2,\Delta E_d)\frac{\hat{s}}{2\pi z_c^2z_d^2} \\
&\times \frac{d\sigma}{d\hat{t}}(ab\to cd)\delta^4(p_a+p_b-p_c-p_d), \qquad (20)
\end{aligned}
$$

of two back-to-back hadrons from independent fragmentation of the back-to-back jets. Let us assume hadron h_1 is a triggered hadron with $p_{T1} = p_T^{\rm trig}$. One can define a hadron-triggered fragmentation function (FF) as the back-to-back correlation with respect to the triggered hadron:

$$
D^{h_1h_2}(z_T,\phi,p_T^{\rm trig}) = p_T^{\rm trig}\frac{d\sigma_{AA}^{h_1h_2}/d^2p_T^{\rm trig}dp_Td\phi}{d\sigma_{AA}^{h_1}/d^2p_T^{\rm trig}}, \qquad (21)
$$

similarly to the direct-photon triggered FF [39] in γ-jet events. Here, $z_T = p_T/p_T^{\rm trig}$ and integration over

$|y_{1,2}| < \Delta y$ is implied. In a simple parton model, two jets should be exactly back-to-back. The initial parton transverse momentum distribution in our model will give rise to a Gaussian-like angular distribution. In addition, we also take into account transverse momentum smearing within a jet using a Gaussian distribution with a width of $\langle k_\perp \rangle = 0.6$ GeV/c.

Shown in Fig. 8 are the calculated back-to-back correlations for charged hadrons in $Au + Au$ collisions as compared to STAR data [9]. The same energy loss that is used to calculate single hadron suppression can also describe well the observed away-side hadron suppression and its centrality dependence.

With cross sections integrated over ϕ, one obtains a hadron-triggered FF $D^{h_1 h_2}(z_T, p_T^{\rm trig})$. The suppression factor $I_{AA}(z_T, p_T^{\rm trig}) \equiv D_{AA}^{h_1 h_2}/D_{pp}^{h_1 h_2}$ defined by the STAR experiment [9] is just the medium modification factor of the hadron-triggered FF. The shape of the modification is predicted [60] to be very similar to that of the γ-triggered fragmentation functions [39].

C. High p_T Azimuthal Anisotropy

Another aspect of jet quenching is the azimuthal anisotropy of the spectra caused by parton energy loss which depends on the azimuthal angle according to Eq. (19). In non-central collisions, the average path length of parton propagation will vary with the azimuthal angle relative to the reaction plane. This leads to an azimuthal dependence of the total parton energy loss and therefore azimuthal asymmetry of high-p_T hadron spectra [5,6]. Such an asymmetry is another consequence of parton energy loss and yet it is not sensitive to effects of initial state interactions.

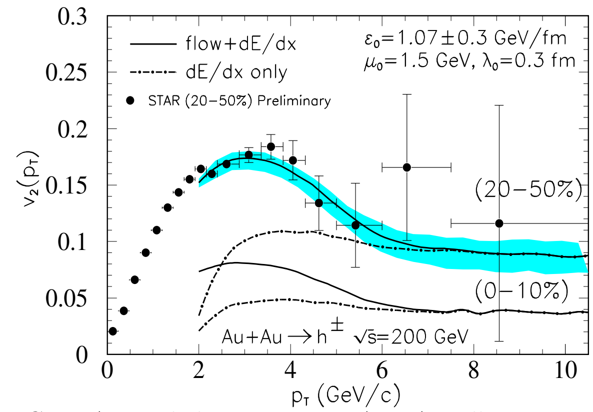

FIG. 9. Azimuthal anisotropy in $Au+Au$ collisions as compared to the STAR [61] 4-particle cumulant result.

In the parton model, the azimuthal angle dependence of the parton energy loss will be given by Eq. (14) for non-central heavy-ion collisions. Using this anisotropic energy loss in the effective modified fragmentation functions in the pQCD parton model of high-p_T hadron spectra, one will obtain azimuthal anisotropic spectra. Shown in Fig. 9 is $v_2(p_T)$ (second Fourier coefficient of the azimuthal angle distribution) of charged hadrons generated from parton energy loss (dot-dashed) as compared to preliminary STAR data [61], using the 4-particle cumulant moments method [62] which is supposed to reduce non-geometrical effects such as inherent two-particle correlations from di-jet production [63]. The energy loss extracted from high-p_T hadron spectra suppression can also account for the observed azimuthal anisotropy at large p_T. At intermediate p_T, the observed v_2 is larger than the simple parton model calculation. Such discrepancy can also be attributed to effects of parton coalescence. If the difference is made up by kaons and baryons from the coalescence contribution, one finds that they must have $v_2 \approx 0.23$ (0.11) for 20-50% (0-10%) collisions. The total $v_2(p_T)$ is shown by the solid lines.

V. JET QUENCHING AND QGP

From both single and di-hadron spectra and their azimuthal anisotropy $v_2(p_T)$, the extracted average energy loss in the parton model calculation for a 10 GeV quark in the expanding medium is $\langle dE/dL \rangle_{1d} \approx 0.85 \pm 0.24$ GeV/fm, which is equivalent to $dE_0/dL \approx 13.8 \pm 3.9$ GeV/fm in a static and uniform medium over a distance $R_A = 6.5$ fm at an initial time $\tau_0 = 0.2$ fm. Compared to the energy loss extracted from HERMES data on DIS, this value is about 30 times higher than the quark energy loss in cold nuclei. Since the parton energy loss in the thin plasma limit is proportional to the gluon number density, one can conclude that the initial gluon density reached in the central $Au + Au$ collisions at 200 GeV should be about 30 times higher than the gluon density in a cold Au nucleus. This number is consistent with the estimate from the measured rapidity density of charged hadrons [64] using the Bjorken scenario [65] and assuming duality between the number of initial gluons and final charged hadrons. Given the measured total transverse energy $dE_T/d\eta \approx 540$ GeV or about 0.8 GeV per charged hadron [66] in central $Au + Au$ collisiona at $\sqrt{s} = 130$ GeV, the initial energy density in the formed dense matter is at least 100 times higher than in cold nuclear matter.

The above analyses of RHIC data on jet quenching are based on an assumption that the dominant mechanism of the jet quenching is parton energy loss before hadronization. It is reasonable to ask whether leading hadrons from the jet fragmentation could have strong interaction with the medium and whether hadron absorption could be the main cause for the observed jet quenching [67,68]. A detailed analysis of jet quenching data concludes [69] that the data are not compatible with such a scenario of hadronic absorption and that the observed patterns of jet quenching in heavy-ion collisions at RHIC are the consequences of parton energy loss.

In addition to the long hadron formation time (it could be about 30-70 fm/c for a 10 GeV pion) estimated from uncertainty principle, there are many patterns of the suppression that can rule out the scenario of hadron absorption. Foremost among them, the large v_2 at high-p_T and the same-side di-hadron correlation cannot be reconciled with the hadron absorption scenario.

Since the spectra azimuthal anisotropy is caused by the eccentricity of the dense medium, which decreases rapidly with time [70] due to transverse expansion, it has to happen at a very early time. By the time of a few fm/c before the formation of any hadron, the eccentricity is already reduced to a non-significant value. Any further interaction cannot produce much spectral anisotropy.

As shown in Fig. 7, while the away-side jet is suppressed in central $Au+Au$ collisions, the same-side di-hadron correlation remains almost the same as in $p+p$ and $d+Au$ collisions. This is clear evidence that jet hadronization takes place outside the dense medium with a reduced parton energy. A recent study [71] of the di-hadron fragmentation functions shows that the conditional (or triggered) di-hadron distribution at large z within a jet is quite stable against radiative evolution (or energy loss). On the other hand, a hadron absorption mechanism will suppress both the leading and secondary hadron in the same way and should lead to the suppression of the same-side correlation. One may argue that the jet whose leading hadron is the trigger hadron could be emitted from the surface and suffer no energy loss and therefore no suppression of the secondary hadron. This is not true. Even though the same-side jet is produced close to the surface due to trigger bias, the parton jet on the average still loses about 2 GeV energy [69], which should be carried by soft hadrons in the direction of the triggered hadron. If one collects all the energy from jet fragmentation for a fixed value of trigger hadron p_T, this energy in central $Au+Au$ collisions should be larger than in $p+p$ collisions by about 2 GeV. This has indeed been confirmed by preliminary data from STAR [15], as shown in Fig. 10, where the scalar sum of the p_T on the same side of the triggered hadron is indeed about 2 GeV larger in central $Au+Au$ than in $p+p$ collisions. The analysis [15] also sees modification of the fragmentation function both along and opposite the trigger hadron direction.

Finally, if hadron absorption can suppress high-p_T hadrons and jets, it would most likely happen in heavy-ion collisions at the SPS energy. Hadron spectra at this energy are very steep at high-p_T and are very sensitive to initial transverse momentum broadening and parton energy loss [45]. However, the measured π spectrum in central $Pb+Pb$ collisions only shows the expected Cronin enhancement [72,73] with no sign of significant suppression. More recent analysis of the $Pb+Pb$ data at the SPS energy also shows [74] that both same-side and back-side jet-like correlations are not suppressed, though the back-side distribution is broadened.

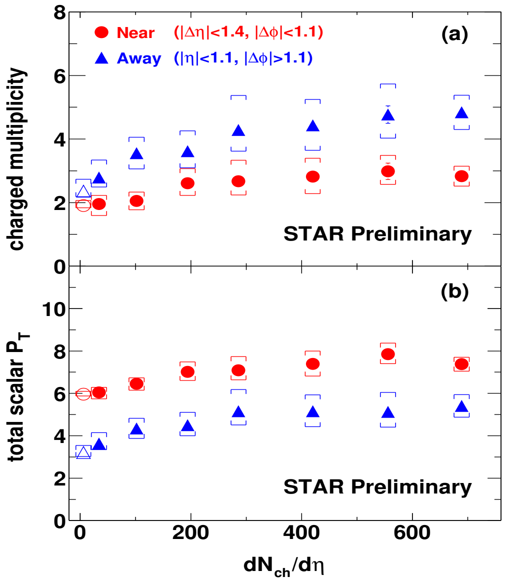

FIG. 10. Charged hadron multiplicity (a) and total scalar p_T (b) of the near and away side jet with $4 < p_T^{trig} < 6$ GeV/c and $0.15 < p_T < 4$ GeV/c in $p+p$ (open symbols) and $Au+Au$ collisions from STAR experiment [15].

VI. SUMMARY AND OUTLOOK

In summary, with the latest measurements of high-p_T hadron spectra and jet correlations in $p+p$, $d+Au$ and $Au+Au$ collisions at RHIC, the observed jet quenching in $Au+Au$ collisions has been established as a consequence of final-state interaction between jets and the produced dense medium. The collective body of data, suppression of high-p_T spectra and back-to-back jet correlation, high-p_T anisotropy and centrality dependence of the observables, points to parton energy loss as the culprit of the observed jet quenching. A simultaneous phenomenological study within a LO pQCD parton model incorporating the parton energy loss describes the experimental data of $Au+Au$ collisions very well. The extracted average energy loss for a 10 GeV quark in the expanding medium is equivalent to an energy loss in a static and uniform medium that is about 30 times larger than in a cold nucleus. This leads us to conclude that the initial gluon (energy) density is about 30 (100) times of that in a cold nuclear matter. It is inconceivable to imagine within our current understanding of QCD that any form of matter other than quark gluon plasma could exist at such a high energy density.

A vast collection of data on bulk properties of particle production in heavy-ion collisions provides further evidence of QGP formation at RHIC. Most striking among these data is the elliptic flow measurement [75] that was found to saturate the hydrodynamic limit and exhibit all the signature behaviors of a collective hydro flow, including the splitting in $v_2(p_T)$ of hadrons with different masses [18,19,76]. Together with jet quenching, they provide a collection of compelling evidence for the formation of a strongly interacting quark gluon plasma at RHIC.

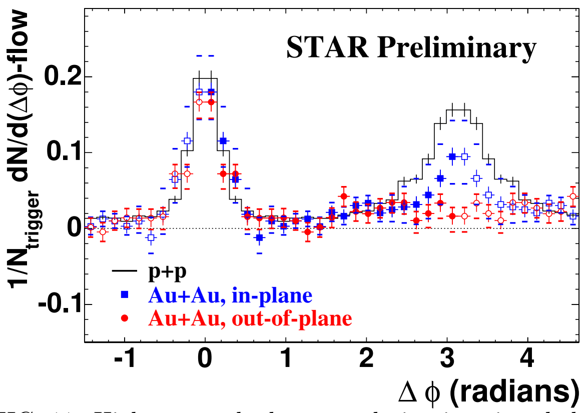

FIG. 11. High-p_T two-hadron correlation in azimuthal angle for different orientation of the trigger hadron with respect to the reaction [14].

The discovery of jet quenching and its central role in characterizing the strongly interacting quark gluon plasma will usher in another new era of studying the properties of the dense matter. With higher-precision data extending to higher p_T, one can find out the missing energy of the suppressed jet [77] and map out in detail the modification of the fragmentation functions and the energy dependence of the energy loss. The preliminary results from STAR [15] on modified fragmentation functions in $Au + Au$ collisions have already demonstrated the potential of the study and could even shed light on parton thermalization and hadronization. The ultimate measurement of a modified FF in heavy-ion collisions will be the direct-γ-triggered FF. Recent preliminary results from PHENIX [16] indicate that such a measurement will be within reach in the next few years.

One can also study the di-hadron correlations with respect to the reaction plane. Shown in Fig. 11 are the preliminary results from STAR [14] on di-hadron correlation in azimuthal angle for different orientations of the trigger hadron with respect to the reaction plane in $Au + Au$ collisions. This could eventually measure the length dependence of the parton energy loss. Measurements of charmed mesons can provide additional tests of the parton energy loss scenario of jet quenching, since recent theoretical studies [78–81] show many unique features of heavy quark energy loss that are different from those of light quarks and gluons.

ACKNOWLEDGEMENT

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