

## Jet modification in three dimensional fluid dynamics at next-to-leading twist

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(Received 17 April 2007; revised manuscript received 29 August 2007; published 23 October 2007)

The modification of the single inclusive spectrum of high transverse momentum ( $p_T$ ) pions emanating from an ultrarelativistic heavy-ion collision is investigated. The deconfined sector is modeled using a full three-dimensional (3-D) ideal fluid dynamics simulation. Energy loss of high  $p_T$  partons and the ensuing modification of their fragmentation is calculated within perturbative QCD at next-to-leading twist, where the magnitude of the higher twist contribution is modulated by the entropy density extracted from the 3-D fluid dynamics simulation. The nuclear modification factor ( $R_{AA}$ ) for pions with a  $p_T \geq 8$  GeV as a function of centrality as well as with respect to the reaction plane is calculated. The magnitude of contributions to the differential  $R_{AA}$  within small angular ranges, from various depths in the dense matter, is extracted from the calculation and demonstrates the correlation of the length integrated density and the  $R_{AA}$  from a given depth. The significance of the mixed and hadronic phase to the overall magnitude of energy loss is explored.

DOI: [10.1103/PhysRevC.76.041902](https://doi.org/10.1103/PhysRevC.76.041902)

PACS number(s): 12.38.Mh, 11.10.Wx, 24.10.Nz, 25.75.Dw

Current experiments at the Relativistic Heavy Ion Collider (RHIC) have suggested two fundamental properties of the dense partonic matter produced in the collision of heavy ions [1]. Comparisons with ideal fluid dynamical simulations [2] indicate that the produced matter probably achieves local thermal equilibrium rapidly and may be ascribed a very low viscosity. Second, the propagation of high transverse momentum partons, produced in the early hard scatterings, is systematically impeded. The spectrum of high transverse momentum hadrons with  $p_T \geq 8$  GeV shows a reduction of almost a factor of 5 in central Au-Au collisions compared to that from the spectrum of such particles in elementary nucleon-nucleon encounters enhanced by the number of expected binary collisions [3].

In most estimates from *jet-quenching* calculations [4,5], there exists no real dynamical model of the medium [6]; in the best of circumstances, a time-dependent, three-dimensional, functional form  $\rho(\vec{r}, \tau) \sim \rho_0(\vec{r})(\tau_0/\tau)$ , with the initial density estimated from a wounded-nucleon model, is used. There have been extensions involving a one-dimensional hydrodynamical model [7]; however, such calculations cannot address the centrality dependence of the modification. The first attempt involving a realistic model of the medium [8] estimated the effects of three-dimensional (3-D) expansion on the uncharged pion spectrum and the  $R_{AA}$ . However, this approach treated the energy loss of jets in a rather simplified manner. Detailed 3-D simulations of the medium have recently become available in Ref. [9], where the various parameters of the simulation were fitted by comparison with experimental data in the soft sector. A study of the modification of jets in this medium, using the quenching weights of Ref. [5], was carried out in Ref. [10]. However, such calculations, along with others, have tended to reformulate the effect of a 3-D expanding medium in terms of an effective length. In the current effort, the medium modified fragmentation functions of high  $p_T$  pions will be computed in the higher-twist energy loss formalism [11]. The 3-D expanding plasma of Ref. [9] will furnish the space-time dependence of the two-gluon matrix element (related to

the well-known transport coefficient  $\hat{q}$ ), which parameterizes the sensitivity of the jet-modification equations to the properties of the medium. The calculation of the contributions that lead to the Landau Pomeranchuk Migdal (LPM) interference will be carried out directly in such a medium without recourse to an effective length. In the following, the salient features of the 3-D fluid dynamical simulations and the ensuing calculation of the nuclear modification factor at next-to-leading twist will be briefly reviewed.

Ideal Relativistic Fluid Dynamics (RFD, see, e.g., Ref. [12]) is the study of matter where the precise degrees of freedom are unimportant to the evolution of the extensive variables of the system. It is ideally suited for the high-density strongly coupled phase of heavy-ion reactions at RHIC, but breaks down in the later, dilute, stages of the reaction when the mean free paths of the hadrons become large and flavor degrees of freedom are important. RFD is based on the solution of the four equations of energy momentum conservation,

$$\partial_\mu T^{\mu\nu}(x) = 0, \quad (1)$$

where  $T^{\mu\nu}$  is the energy momentum tensor, given by

$$T^{\mu\nu} = (\epsilon + p)U^\mu U^\nu - pg^{\mu\nu}. \quad (2)$$

Here,  $\epsilon$ ,  $p$ ,  $U$ , and  $g^{\mu\nu}$  are the local energy density, pressure, four velocity, and metric tensor, respectively, at  $x$ . The relativistic hydrodynamic equation [Eq. (1)] is solved numerically using baryon number  $n_B$  conservation,

$$\partial_\mu (n_B(T, \mu)U^\mu) = 0, \quad (3)$$

as a constraint and closing the resulting set of partial differential equations by specifying an equation of state (EoS):  $\epsilon = \epsilon(p)$ . In the ideal fluid approximation (i.e., neglecting off-equilibrium effects), once the initial conditions for the calculation have been fixed, the EoS is the *only* input to the equations of motion and relates directly to the properties of the matter under consideration. As a result, even striking phenomena occurring within systems in local equilibrium, such as phase transitions, may be easily incorporated into the

bulk dynamics of the full system. Ideally, either the initial conditions or the EoS should be determined beforehand by an *ab-initio* calculation (e.g., the EoS may be obtained from a lattice-QCD calculation), in which case a fit to the data would allow for the determination of the other quantity.

We assume that hydrodynamic expansion starts at  $\tau_0 = 0.6$  fm. Initial energy density and baryon number density are parameterized by

$$\begin{aligned}\epsilon(x, y, \eta) &= \epsilon_{\max} W(x, y; b) H(\eta), \\ n_B(x, y, \eta) &= n_{B\max} W(x, y; b) H(\eta),\end{aligned}\quad (4)$$

where  $b$  and  $\epsilon_{\max}$  ( $n_{B\max}$ ) are the impact parameter and maximum value of energy density (baryon density), respectively.  $W(x, y; b)$  is given by a combination of wounded nuclear model and binary collision model [13] and  $H(\eta)$  is given by  $H(\eta) = \exp[-(|\eta| - \eta_0)^2 / 2\sigma_\eta^2 \cdot \theta(|\eta| - \eta_0)]$ . RFD has been very successful in describing soft matter properties at RHIC, especially collective flow effects and particle spectra [9,14]. All parameters of the RFD evolution [9] have been fixed by a fit to the soft sector (elliptic flow, pseudorapidity distributions, and low  $p_T$  single-particle spectra), thereby providing us with a fully determined medium evolution for the hard probes to propagate through. While different descriptions for the hadronic phase are presented in Ref. [9], in this calculation, the entire evolution of both the deconfined and hadronic phase will be assumed to be describable by the RFD simulation.

Assuming a factorization of initial and final state effects, the differential cross section for the production of a high  $p_T$  hadron at midrapidity from the impact of two nuclei  $A$  and  $B$  is obtained as a convolution of initial nuclear structure functions,  $G_a^A(x_a, Q^2)$ ,  $G_b^B(x_b, Q^2)$  (shadowing functions are taken from Ref. [15]), a hard cross section to produce a high  $p_T$  parton ( $d\hat{\sigma}/d\hat{t}$ ), and a medium modified fragmentation function. The entire effect of final state jet modification is encoded in the medium modified fragmentation function  $\tilde{D}(z, Q^2) = D + \Delta D$ , expressed as the sum of the leading twist vacuum fragmentation function [16] and a correction from rescattering of the struck quark (gluon) in the medium followed by the radiation of a gluon (quark) with transverse momentum  $l_\perp$  that leaves a momentum fraction  $y$  in the quark (gluon) denoted as parton  $i$  [ $P_{q \rightarrow i}(y)$  is the splitting function for this process], which then fragments leading to the detected hadron [11]. In the collinear limit, the modification is computed by isolating corrections, suppressed by powers of  $Q^2$ , which are enhanced by the length of the medium [17]. At next-to-leading twist, the correction for the fragmentation of a quark has the expression (generalized from deep-inelastic scattering (DIS) [11,18])

$$\begin{aligned}\Delta D(z, Q^2, \vec{r}) &= \frac{\alpha_s}{2\pi} \int \frac{dl_\perp^2}{l_\perp^2} \int \frac{dy}{y} P_{q \rightarrow i}(y) 2\pi \alpha_s C_A \\ &\times T^M(\vec{b}, \vec{r}, x_a, x_b, y, l_\perp) D_i(z/y, Q^2) \\ &\times \left[ \int d^2 b d^2 r t_A(\vec{r} + \vec{b}/2) t_B(\vec{r} - \vec{b}/2) \right. \\ &\times \left. l_\perp^2 N_c G_a^A(x_a) G_b^B(x_b) \frac{d\hat{\sigma}}{d\hat{t}} \right]^{-1} + \text{v.c.} \quad (5)\end{aligned}$$

In the above equation,  $t_A(\vec{r})$  and  $t_B(\vec{r})$  are the thickness functions of nuclei  $A$  and  $B$  at the transverse location  $\vec{r} \equiv (x, y)$ . The v.c. refers to virtual corrections. The factor  $T^M$ , which has its origin in higher twist matrix elements of the quark, may be expressed as

$$\begin{aligned}T^M &= \int d^2 b d^2 r t_A(\vec{r} + \vec{b}/2) t_B(\vec{r} - \vec{b}/2) G_a^A(x_a) G_b^B(x_b) \frac{d\hat{\sigma}}{d\hat{t}} \\ &\times \int_0^{\zeta_{\max}} d\zeta x_g \rho_g(x_g, \hat{n}\zeta + \vec{r}) (2 - 2\cos(\eta_L \zeta)). \quad (6)\end{aligned}$$

In the above equation, the factor  $\eta_L = l_\perp^2 / [2\hat{p}_T y(1-y)]$ , where  $\hat{p}_T$  represents the transverse momentum of the produced parent jet.

The jet direction ( $\hat{n}$ ) is chosen as the  $x$  axis in the transverse ( $x, y$ ) plane ( $z$  direction is set by the beam line); the angle of the reaction plane vector  $\vec{b}$  is measured with respect to this direction. The distance traveled by the jet in this direction prior to scattering off a gluon is denoted as  $\zeta$ . The gluon's forward (toward the jet) momentum fraction is denoted as  $x_g$ . In what follows, the nuclear modification of the single inclusive spectrum, both integrated and as a function of the angle with respect to the reaction plane, will be computed. The angle integrated modification, the  $R_{AA}$  [integrated over a small range of impact parameter ( $b_{\min}$  to  $b_{\max}$ )], is defined as

$$R_{AA} = \frac{\frac{d\sigma^{AA}(b_{\min}, b_{\max})}{dy d^2 p_T}}{T_{AA}(b_{\min}, b_{\max}) \frac{d\sigma^{pp}(p_T, y)}{dy d^2 p_T}}, \quad (7)$$

where  $T_{AA}(b_{\min}, b_{\max})$  is related to the mean number of binary nucleon-nucleon encounters in the range of impact parameter chosen ( $\langle N_{\text{bin}} \rangle$ ), the geometric  $A$ - $A$  cross section [ $\sigma_{\text{Geo}}^{AA}(b_{\min}, b_{\max})$ ], and the inelastic  $p$ - $p$  cross section ( $\sigma_{pp}^{\text{in}}$ ) as  $T_{AA}(b_{\min}, b_{\max}) \simeq \langle N_{\text{bin}} \rangle \sigma_{\text{Geo}}^{AA}(b_{\min}, b_{\max}) / \sigma_{pp}^{\text{in}}$ .

The phenomenological input that is required to understand the variation of the nuclear modification factor with centrality or  $p_T$  is the space-time dependence of the gluon density  $\rho(\vec{r}, \tau) = x_g \rho_g(x_g, \vec{r}, \tau)$ . This density is directly proportional to the transport coefficient  $\hat{q}_R$  (for a jet parton in representation  $R$ ) [19,20],

$$\hat{q}_R(\vec{r}, \tau) = \frac{4\pi^2 \alpha_s(Q^2) C_R}{N_c^2 - 1} \rho(\vec{r}, \tau), \quad (8)$$

where  $C_R$  is the representation dependent Casimir and  $N_c$  is the number of colors.

Note that the  $z$  coordinate is considered to be zero as the interest is strictly on midrapidity observables. The gluon density at a given location depends on the phase of matter at that location and is in general assumed to depend on the number of degrees of freedom prevalent in the excited matter at that location. The exact degrees of freedom in the deconfined phase are as yet unknown [21]. The partonic density as seen by the jet within each degree of freedom is also unclear. For this first attempt at higher twist energy loss in a hydrodynamically expanding medium, we invoke the very simple ansatz

$$\hat{q}_R(\vec{r}, \tau) = \frac{\hat{q}_0 T^3(\vec{r}, \tau) [R(\vec{r}, \tau) + c_{\text{HG}} \{1 - R(\vec{r}, \tau)\}]}{\sqrt{1 - U_\perp^2(\vec{r}, \tau) T_0^3}}, \quad (9)$$

where  $T(\vec{r}, \tau)$ ,  $U_{\perp}(\vec{r}, \tau)$ , and  $R(\vec{r}, \tau)$  represent the temperature, the flow transverse to the jet, and the volume fraction in the plasma phase at the space-time point  $\vec{r}, \tau$ . It is this information that is extracted from the RFD simulation. The factors  $\hat{q}_0$  and  $T_0$  represent the maximum  $\hat{q}$  and temperature achieved in the simulation; in this particular version of RFD,  $T_0 = 0.405$  GeV and  $\hat{q}_0$  is a fit parameter adjusted to fit one data point of the  $R_{AA}$  at one centrality. The factor  $c_{HG}$  accounts for the fact that the partonic density per degree of freedom in the hadron gas phase may be different from that in the QGP phase and plots with different choices of  $c_{HG}$  will be presented. It should be pointed out that contrary to the (pre-)hadronic scattering picture of Ref. [22], the energy loss of jets in the hadronic phase is considered here as a completely partonic process [11], where the jets fragment into hadrons outside the medium.

The above prescription of assuming  $\hat{q}$  to be a linear function of  $T^3$  may be referred to as the thermal prescription. In the energy density prescription used in Ref. [10], one assumes  $\hat{q}$  to be a linear function of  $\epsilon^{3/4}$ . Alternatively, one may also assume that  $\hat{q}$  is a linear function of the entropy density  $s$ . Relationships may be drawn between these various schemes with a knowledge of the EoS. For instance, in a plasma phase that is populated by massless partons, the entropy density  $s$  is a linear function of  $T^3$  and the two prescriptions are identical. This breaks down in the hadronic phase as the entropy density is no longer a simple linear function of  $T^3$ . Because of the variety of unknowns regarding the  $\hat{q}$  in an excited hadronic resonance and, as a result, its parametric dependence in the hadron resonance phase, we persist with the simple prescription of Eq. (9).

As mentioned above, the overall fit parameter  $\hat{q}_0$  is tuned to fit one experimental data point, at one centrality and  $p_T$ . For the current effort, the fit parameter is set by requiring that the  $R_{AA}$  at  $p_T = 10$  GeV in the most central event (0–5% centrality) is 0.2. For a  $c_{HG} = 1$  this corresponds to a  $\hat{q}_0 \simeq 1.3$  GeV<sup>2</sup>/fm for gluon jets. With the values of  $\hat{q}_0$  and  $c_{HG}$  fixed, the variations of  $R_{AA}$  as a function of  $p_T$  and centrality of the collision are predicted. These are presented in Fig. 1 for two centralities (0–5% and 20–30%) in comparison with data from PHENIX [23]. An important caveat to the estimates of  $\hat{q}$  are the estimates of the plasma temperature from the very early phase of the evolution. Hard jets are produced very early ( $t < 0.1$  fm/c) in the evolution of the system. However, fits to the spectrum of soft particles indicate that a locally equilibrated matter does not arise before 0.6 fm/c. This is the starting point of the RFD simulations used. In the regime between 0.1 to 0.6 fm/c, estimates of  $\hat{q}$  are open to speculation. In this first effort, the initial  $\hat{q}$  is assumed to be held constant from the time the jets are produced to 0.6 fm/c. In future efforts, variations of this form will be attempted.

In the RFD simulations used, an equation of state with a first order phase transition is included. As a result, such simulations include a prominent hadronic and mixed phase. In Fig. 1, the factor  $c_{HG} = 1$ , meaning that the  $\hat{q}$  per unit of  $T^3$  in the deconfined phase is similar to that in the hadron gas phase. From the equation of state used, one may discern the approximate and simple relation  $s_{QGP} \sim 5s_{HG}$ , between the entropy densities in the deconfined phase ( $s_{QGP}$ ) and the hadronic phase ( $s_{HG}$ ), in the vicinity of the phase transition.

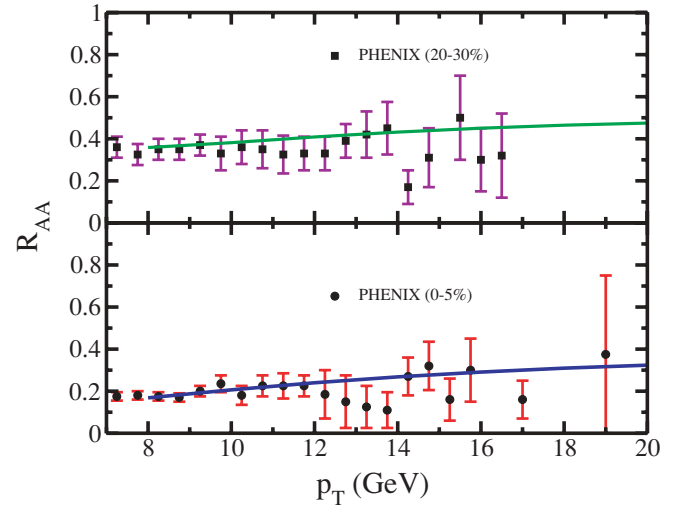


FIG. 1. (Color online)  $R_{AA}$  (with  $\hat{q}_0 \simeq 1.3$  GeV<sup>2</sup>/fm and  $c_{HG} = 1$ ) in Au-Au collisions at 0–5% (blue line) and 20–30% (green line) centrality compared to data from PHENIX [23].

This implies that a  $c_{HG} = 1$  places the ratio  $\hat{q}/s$ , i.e., the quenching strength per unit of entropy, in the hadronic phase to be five times larger than that in the QGP phase. Because of the error bars in the data, a range of values of  $c_{HG}$  is allowed; however, certain extremes may be ruled out.

As an illustration, the  $R_{AA}$  for the same two centralities as in Fig. 1 are plotted for two other cases of  $c_{HG} = 0$  and  $c_{HG} = 10$  in Fig. 2. In each case, the overall normalization factor  $\hat{q}_0$  is readjusted to the standard condition that the  $R_{AA}$  for a pion  $p_T = 10$  GeV and at a centrality of 0–5% be 0.2 (the

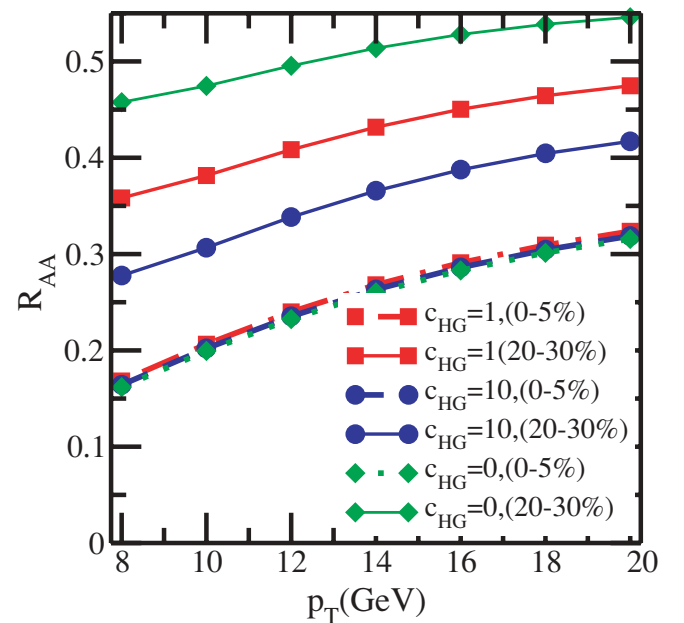


FIG. 2. (Color online)  $R_{AA}$  as a function of centrality for different choices of  $c_{HG} = \hat{q}_{HG}/T^3$  in the hadronic phase. The  $\hat{q}_0$  is readjusted in each case (for  $c_{HG} = 10$ ,  $\hat{q}_0 = 0.24$  GeV<sup>2</sup>/fm; for  $c_{HG} = 0$ ,  $\hat{q}_0 = 3.0$  GeV<sup>2</sup>/fm). See text for details.

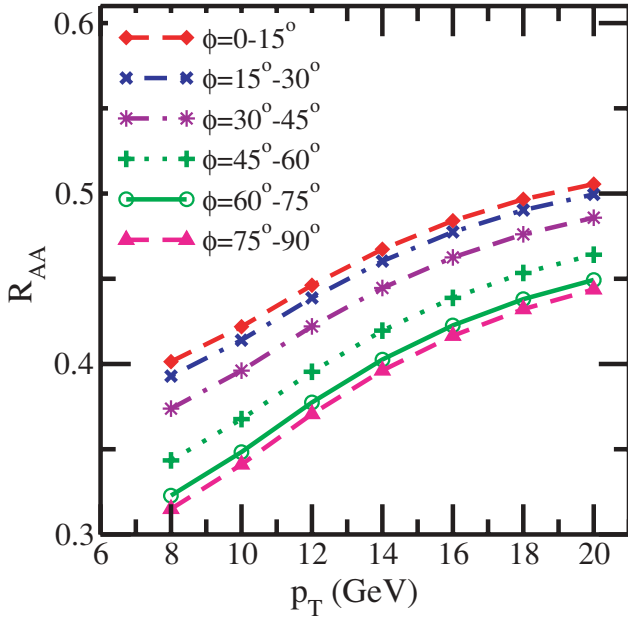


FIG. 3. (Color online)  $R_{AA}$  as a function of the reaction plane for Au-Au collisions at 20–30% centrality.

readjusted values are mentioned in the caption of Fig. 2). This is the reason that all three curves for different choices of  $c_{HG}$ , at 0–5% centrality, essentially lie on top of one another. However, the  $R_{AA}$  for the semicentral case (20–30%) shows noticeable variations. The red squares correspond to the case of Fig. 1. The case  $c_{HG} = 0$ , represented by the green diamonds, represents the case of no quenching in the hadronic phase. Because of the proportionately larger presence of the hadronic phase in more peripheral collisions, this case understandably demonstrates less modification.<sup>1</sup> The other extreme is with a  $c_{HG} = 10$ , where the  $\hat{q}$  in units of  $T^3$  is an order of magnitude larger in the hadronic phase compared to the deconfined phase. As would be expected, the results show an opposite trend: the centrality dependence of the  $R_{AA}$  seems to be somewhat underpredicted, though not completely excluded within experimental errors.

In semicentral events, the average path length traveled by jets in the reaction plane are different from the path lengths traveled out of plane. This should lead to the occurrence of an azimuthally dependent nuclear modification factor. Full three-dimensional RFD simulations capable of reproducing the observed azimuthal asymmetry in the low  $p_T$  spectrum require that the initial gluon density profile displays a considerable azimuthal eccentricity in the transverse plane. The hard scatterings, which lead to the production of back-to-back jets, tend to produce these isotropically. As a result, a sizable azimuthal asymmetry in the nuclear modification factor demonstrates the sensitivity of this observable to the early density profile of the QGP. As an illustration of this

<sup>1</sup>In the RFD simulations, no attempt is made to incorporate the fact that central collisions endure a longer lifetime and hence freeze-out later than peripheral collisions; simulations for all centralities are run up to 10 fm/c and then frozen out. Beyond this time, all simulations at these energies seem to demonstrate a very low density.

effect, the nuclear modification factor at an angle  $0^\circ \leq \phi \leq 90^\circ$  with respect to the reaction plane, in angular bins of  $15^\circ$  [19], is plotted in Fig. 3. One notes, that the  $R_{AA}$  in the plane shows the least modification, which grows as the angular bin is turned away from the reaction plane. As a result, it may be argued that the initial profile of the produced matter plays an important role in the determination of the final nuclear modification factor. The rate of change of the  $R_{AA}$  as a function of  $\phi$  shows the interesting pattern of being slow at the terminal points and maximal in the middle of the angular range (i.e.,  $\phi \sim 45^\circ$ ). The results obtained from this analysis (and especially the symmetric behavior of the modification as a function of angle with respect to the reaction plane) show a pattern identical to that of the experimental measurements of the  $R_{AA}$  versus the reaction plane at lower  $p_T$ , as reported in Ref. [24]. The shallow rise in the  $R_{AA}$  as a function of the  $p_T$  of the detected hadron is demonstrated by the nuclear modification factor in each angular bin: this property is thus independent of the density of the matter probed by the jet.

The existence of a calculation scheme for the modification of jets in a three-dimensional dynamically evolving medium allows for a more differential analysis of jet quenching in medium. As an example, we compute the  $R_{AA}$  as a function of location in the reaction plane of the 20–30% centrality collisions. The focus is on the  $R_{AA}$  in the angular bin of 0– $15^\circ$  with respect to the reaction plane, this corresponds to a  $30^\circ$  opening angle around the impact parameter vector  $\vec{b}$ . The transverse plane is then divided into ten longitudinal strips of width 1 fm perpendicular to  $\vec{b}$  (as illustrated in the inset in Fig. 4). The  $R_{AA}$  for jets originating in a given strip is then plotted (red dot-dashed line) as a function of the location of

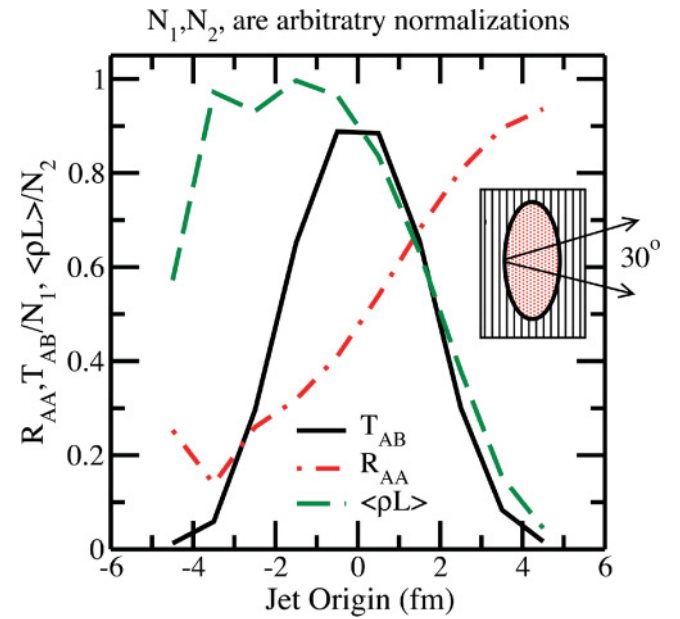


FIG. 4. (Color online)  $R_{AA}$  as a function of the location along the impact parameter from longitudinal strips perpendicular to  $\vec{b}$ . Also plotted is the number of initial binary nucleon-nucleon collisions and the length integral of the time-dependent density weighted with the number of binary collisions ( $\int dl \rho(\vec{r}, t)$ ) arbitrarily normalized to fit on the same plot.

the midpoint of the strip in the reaction plane. Also plotted is the initial number of binary nucleon-nucleon scatterings in each of the strips (black solid line). Plotted along with this is the averaged length integrated density  $\rho$  (which is directly proportional to  $\hat{q}$ , see Refs. [19,20]) as experienced by a jet originating in each strip (green dashed line). This is defined as the following quantity:

$$\langle \rho L \rangle = \frac{\int d^2b d^2r d\phi_{\hat{n}} d\zeta t_A(\vec{r} + \vec{b}/2) t_B(\vec{r} - \vec{b}/2) \rho(\vec{r} + \zeta \hat{n}, \zeta)}{[\int d^2b d^2r t_A(\vec{r} + \vec{b}/2) t_B(\vec{r} - \vec{b}/2)]}$$

In the above equation, the density  $\rho$  is evaluated at the location  $\vec{r} + \hat{n}\zeta$  and at time  $\zeta$ , where  $\vec{r}$  is the location of the jet vertex,  $\zeta$  is the time elapsed, and  $\hat{n}$  is the direction of propagation of the jet. The range of  $\phi_{\hat{n}}$  is from  $-15^\circ$  to  $15^\circ$ , while  $\vec{r}$  ranges within the specified strip. Thus,  $\langle \rho L \rangle$  represents the length integrated density as experienced by a jet, on average, in the given strip. As may be noted from Fig. 4, the  $R_{AA}$  in a given strip is anticorrelated with the mean integrated density experienced by the jet on its way out of the medium. The  $R_{AA}$  tends to show a small rise in the strip farthest from exit (note that this is correlated with a drop in  $\langle \rho L \rangle$ ); this is due to the fact that a considerable fraction of jets that originate on this strip assume tangential paths of exit that do not pass through the center of the dense system produced. As one moves toward the middle of the impact parameter, a large fraction of jets, because of the

angular cut, are forced to pass through the densest part of the matter and, as a result, the  $R_{AA}$  first drops and then rises.

In this rapid communication, higher twist calculations of energy loss have been extended to include realistic medium evolution as stipulated by a 3-D hydrodynamical simulation. The nuclear modification factor at different centralities as a function of the  $p_T$  of the detected hadron has been calculated and demonstrates very good agreement with experimental data. The variation of the  $R_{AA}$  with centrality shows a noticeable sensitivity to the presence of a hadronic phase. This may be used to place bounds on the ratio between the  $\hat{q}$  in the hadronic phase compared to the partonic phase. More differential observables, such as the  $R_{AA}$  versus the reaction plane, are calculated and show qualitative agreement with the data at a somewhat lower  $p_T$ . These results demonstrate the sensitivity of the  $R_{AA}$  at a given angle with respect to the reaction plane to the initial profile of the produced matter. The nuclear modification factor as a function of the depth in the medium is shown to be anticorrelated with the length integrated density as experienced by a jet on average. The magnitude of contributions to the  $R_{AA}$  from different depths demonstrates the diminished surface bias in semicentral collisions.

The authors thank B. Müller for insightful discussions. This work is supported in part by a grant from the U.S. Department of Energy (DE-FG02-03ER41239-0).

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