

Angular variation of hard back-to-back hadron suppression in heavy-ion collisions

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The basic idea of jet tomography is to infer information about the density evolution of the medium created in heavy-ion (A - A) collisions by studying the suppression of hard probes in an A - A environment as compared to the baseline process known from proton-proton (p - p) collisions. The suppression of back-to-back correlations in heavy-ion collisions allows, due to a different geometrical bias, a view into the medium that is qualitatively different from the one offered by single-hadron suppression. A control parameter for the suppression corresponding to a systematic variation of in-medium path lengths and densities can be obtained by studying collisions at finite impact parameter b . A systematic variation of path length can then be introduced by studying the suppression pattern as a function of the angle ϕ with the reaction plane. Using a three-dimensional hydrodynamic evolution model for the medium and a Monte Carlo model that has been shown to successfully reproduce the measured suppression of back-to-back correlations in central collisions of Au-Au at 200A GeV, we compute the suppression as a function of ϕ for b of 2.4, 4.5, 6.3, and 7.5 fm. Given that this involves variations in both control parameters b and ϕ , a comparison with data should eventually allow one to place strong constraints on the combination of the energy loss model and the medium evolution model.

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I. INTRODUCTION

The original aim of studying medium-induced modifications to hard processes in heavy-ion collisions has been to perform jet tomography on the medium [1–6]. The underlying idea is as follows. Hard processes take place simultaneously with soft processes responsible for the creation of soft bulk matter. This places partons emerging from hard vertices inside the medium, and interactions with the soft medium subsequently lead to an energy transfer from hard partons to the soft medium, resulting in a suppression of observed hard hadron yield. The strength of the interaction with the medium reflects the density of the medium, thus one should be able to gain information about the medium density from the observed strength of the suppression. The chief observable considered so far has been the nuclear suppression factor of single hadrons R_{AA} (cf., e.g., Refs. [7,8]) which is the measured yield in A - A collisions divided by the yield in proton-proton (p - p) collisions multiplied by the number of binary scatterings (i.e., the default expectation if soft processes were not forming a medium in an A - A collision).

Unfortunately, a number of problems occur when one tries to carry out this program in practice. First, R_{AA} is not a sensitive measure of the averaged energy loss probability distribution $\langle P(\Delta E, E) \rangle_{T_{AA}}$ [9,10]. Assuming that eikonal propagation is a good approximation for hard partons, this quantity is the complete momentum space information about energy loss—given that the vertex of origin of any observed hard hadron cannot be known, production point, direction, and hence the length of propagation through the medium must be averaged probabilistically, leaving a momentum space probability distribution that a parton with original

energy E lost the amount ΔE to the medium. The fact that $\langle P(\Delta E, E) \rangle_{T_{AA}}$ cannot be reliably determined from R_{AA} is rather unfortunate, as knowledge of the energy loss probability distribution in momentum space is necessary before any mapping into position space (and hence any determination of density profiles) can be attempted.

The second obstacle is that there is not even qualitative agreement as to what the precise nature of in-medium energy loss is. In addition to medium-induced radiation [1–6], a number of publications have advocated a sizable component of elastic energy loss [11–16] which has a qualitatively different path-length dependence (linear in a homogeneous medium, as opposed to quadratic for radiative energy loss). As energy loss models typically contain one adjustable parameter linking the medium density with the interaction strength, R_{AA} could be reproduced by a number of very different scenarios.

Some steps can be taken to improve the situation. First, instead of trying to determine an *a priori* unknown density evolution using hard probes, one may settle for the more modest goal of using a medium evolution model that is already constrained by soft hadronic observables and studying the variation of the hard probes with changes of a known control parameter such as centrality. In this way, one employs hard probes as a test of a specific combination of evolution and energy loss model. Using the same three-dimensional hydrodynamic evolution that has been successful in describing bulk matter evolution [17], researchers have shown that different energy loss models predict different suppression patterns as a function of reaction plane when going to off-central collisions, even when all are adjusted to describe the observed suppression for central collisions [18–20].

One may also turn to observables in which the medium geometry is probed in a qualitatively different way, such as hard back-to-back correlations. Such correlations have been measured by the STAR Collaboration [21]. Employing again

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the strategy to use a set of *a priori* constrained medium evolution models, the work in Refs. [22,23] could show that the path-length dependence of radiative energy loss is compatible with the strength of back-to-back suppression, whereas a linear path-length dependence characteristic of elastic energy loss is not [24], thus any contribution with linear path-length dependence must be a small correction.

In the present paper, we continue the work of Refs. [18, 23] in computing the suppression of back-to-back correlations for noncentral collisions as a function of the angle with the reaction plane. Since this is a very differential observable, it constitutes a rather sensitive test of the radiative energy loss model [2,4] in combination with the three-dimensional hydrodynamic evolution model [17].

II. THE MODEL

The main ingredients of the model we wish to use have been extensively described in Refs. [17,18,22,23,25]. We will therefore confine ourselves here to a summary and describe only the improvements made to the framework in more detail.

We calculate the correlation strength of hadrons back to back with a hard trigger in a Monte Carlo (MC) simulation. There are three important building blocks to this computation: (1) the primary hard parton production, (2) the propagation of the partons through the medium, and (3) the hadronization of the primary partons. Only step (2) probes properties of the medium, and hence it is here that we must specify details of the evolution of the medium and of the parton-medium interaction.

The strength of the parton-medium interaction contains one adjustable parameter K (see below). This parameter is fixed by the requirement that the model should describe the suppression of single inclusive pions for central collisions. The results for noncentral collisions, back-to-back correlations, different hadron species, and different orientation with respect to the reaction plane are (given the hydrodynamic calculation as input) obtained without additional free parameters.

A. Primary parton production

The production of two hard partons k, l in leading order (LO) perturbative quantum chromodynamics (pQCD) is described by

$$\frac{d\sigma^{AB \rightarrow kl+X}}{dp_T^2 dy_1 dy_2} = \sum_{ij} x_1 f_{i/A}(x_1, Q^2) x_2 f_{j/B}(x_2, Q^2) \frac{d\hat{\sigma}^{ij \rightarrow kl}}{d\hat{t}}, \quad (1)$$

where A and B stand for the colliding objects (protons or nuclei) and $y_{1(2)}$ is the rapidity of parton $k(l)$. The distribution function of a parton type i in A at a momentum fraction x_1 and a factorization scale $Q \sim p_T$ is $f_{i/A}(x_1, Q^2)$. The distribution functions are different for the free protons [26,27] and nucleons in nuclei [28,29]. The fractional momenta of the colliding partons i, j are given by $x_{1,2} = \frac{p_T}{\sqrt{s}} (\exp[\pm y_1] + \exp[\pm y_2])$.

Expressions for the pQCD subprocesses $\frac{d\hat{\sigma}^{ij \rightarrow kl}}{d\hat{t}}(\hat{s}, \hat{t}, \hat{u})$ as a function of the parton Mandelstam variables $\hat{s}, \hat{t}$, and $\hat{u}$

can be found, e.g., in Ref. [30]. By selecting pairs of k, l while summing over all allowed combinations of i, j , i.e., $gg, gq, g\bar{q}, qq, q\bar{q}$, and $\bar{q}\bar{q}$, where q stands for any of the quark flavors u, d, s , we find the relative strength of different combinations of outgoing partons as a function of p_T .

For the present investigation, we require $y_1 = y_2 = 0$, i.e., we consider only back-to-back correlations detected at midrapidity. In a first step, we sample Eq. (1) summed over all k, l to generate p_T for the event; in the second step, we perform a MC sampling of the decomposition of Eq. (1) according to all possible combinations of outgoing partons k, l at the p_T obtained in the first step. We thus end with a back-to-back parton pair with known parton types and flavors at transverse momentum p_T .

To account for various effects, including higher order pQCD radiation, transverse motion of partons in the nucleon (nuclear) wave function, and effectively also the fact that hadronization is not a collinear process, we fold into the distribution an intrinsic transverse momentum k_T with a Gaussian distribution, thus creating a momentum imbalance between the two partons as $\mathbf{p}_{T1} + \mathbf{p}_{T2} = \mathbf{k}_T$.

B. Parton propagation through the medium

The probability density $P(x_0, y_0)$ for finding a hard vertex at the transverse position $\mathbf{r}_0 = (x_0, y_0)$ and impact parameter b is in leading order given by the product of the nuclear profile functions as

$$P(x_0, y_0) = \frac{T_A(\mathbf{r}_0 + \mathbf{b}/2)T_A(\mathbf{r}_0 - \mathbf{b}/2)}{T_{AA}(\mathbf{b})}, \quad (2)$$

where the thickness function is given in terms of the Woods-Saxon nuclear density $\rho_A(r, z)$ as $T_A(r) = \int dz \rho_A(r, z)$. Note that Eq. (2) may receive (presumably) small corrections when going beyond a leading order calculation.

If we call the angle between outgoing parton and the reaction plane ϕ , the path of a given parton through the medium $\xi(\tau)$ is specified by (r_0, ϕ) and we can compute the energy loss probability $P(\Delta E)_{\text{path}}$ for this path. We do this by evaluating the line integrals

$$\omega_c(\mathbf{r}_0, \phi) = \int_0^\infty d\xi \xi \hat{q}(\xi) \quad \text{and} \quad \langle \hat{q}L \rangle(\mathbf{r}_0, \phi) = \int_0^\infty d\xi \hat{q}(\xi) \quad (3)$$

along the path, where we assume the relation

$$\hat{q}(\xi) = K 2\epsilon^{3/4}(\xi)(\cosh \rho - \sinh \rho \cos \alpha) \quad (4)$$

between the local transport coefficient $\hat{q}(\xi)$ (specifying the quenching power of the medium), the energy density ϵ , and the local flow rapidity ρ with angle α between the flow and the parton trajectory [31,32]. Energy density ϵ and local flow rapidity ρ are the input from the three-dimensional hydrodynamic simulation of the medium evolution [17]. ω_c is the characteristic gluon frequency, setting the scale of the energy loss probability distribution, and $\langle \hat{q}L \rangle$ is a measure of the path length weighted by the local quenching power. We view parameter K as a tool to account for the uncertainty in the selection of α_s and possible nonperturbative effects increasing

the quenching power of the medium (see discussion in Ref. [23]) and adjust it such that pionic R_{AA} for central Au-Au collisions is described. This leads to a value of $K = 3.6$ [18].

Using the numerical results of Ref. [33], we obtain $P(\Delta E; \omega_c, R)_{\text{path}}$ for ω_c and $R = 2\omega_c^2 / \langle \hat{q} L \rangle$ for given jet production vertex and angle ϕ . In the MC simulation, we first sample Eq. (2) to determine the vertex of origin. For a given choice of ϕ , we then propagate both partons through the medium evaluating Eq. (3) and use the output to determine $P(\Delta E; \omega_c, R)_{\text{path}}$ which we sample to determine the actual energy loss of both partons in the event.

C. Hadronization

Finally, we convert the simulated partons into hadrons. More precisely, to determine if there is a trigger hadron above a given threshold, given a parton k with momentum p_T , we need to sample $A_1^{k \rightarrow h}(z_1, p_T)$, i.e., the probability distribution to find a hadron h from the parton k where h is the most energetic hadron of the shower and carries the momentum $P_T = z_1 p_T$. In the following, we make the assumption that the hadronization process itself, at least for the leading hadrons of a shower, happens well outside the medium, and as a consequence we neglect any interaction of formed hadrons with the medium. The timescale for hadronization of a hadron h in its rest frame can be estimated by the inverse hadron mass, $\tau_h \sim 1/m_h$; boosting this expression to the laboratory frame, one finds $\tau_h \sim E_h/m_h^2$. Inserting a hard scale of 6 GeV or more for the hadron energy and the pion mass in the denominator (as pions constitute the bulk of hadron production), this assumption seems well justified.

In previous works [22,23], we approximated this by the normalized fragmentation function $D_{k \rightarrow h}(z, p_T)$, sampled with a lower cutoff $z_{\min}$ which is adjusted to the reference d -Au data. This procedure can be justified by noting that only one hadron with $z > 0.5$ can be produced in a shower, thus above $z = 0.5$, the $D_{k \rightarrow h}(z, p_T)$ and $A_1^{k \rightarrow h}(z_1, p_T)$ are (up to the scale evolution) identical; and only in the region of low z , where the fragmentation function describes the production of multiple hadrons, do they differ significantly.

We improve on these results by extracting $A_1(z_1, p_T)$ and the conditional probability to find the second most energetic hadron at momentum fraction z_2 given that the most energetic hadron was found with fraction z_1 $A_2(z_1, z_2, p_T)$ from a shower evolution code. The procedure is described in detail in Ref. [25], where we used PYTHIA [34] to simulate the shower; in the present paper, we employ HERWIG [35] instead, which provides a slightly better description of the d -Au baseline data. A comparison of both approaches with the d -Au data is shown in Fig. 1.

Our way of modeling hadronization corresponds to an expansion of the shower development in terms of a tower of conditional probability densities $A_N(z_1, \dots, z_n, \mu)$ with the probability to produce n hadrons with momentum fractions $z_1, \dots, z_n$ from a parton with momentum p_T being $\Pi_{i=1}^n A_i(z_1, \dots, z_i, p_T)$. Taking the first two terms of this expansion is justified as long as we are interested in sufficiently hard correlations.

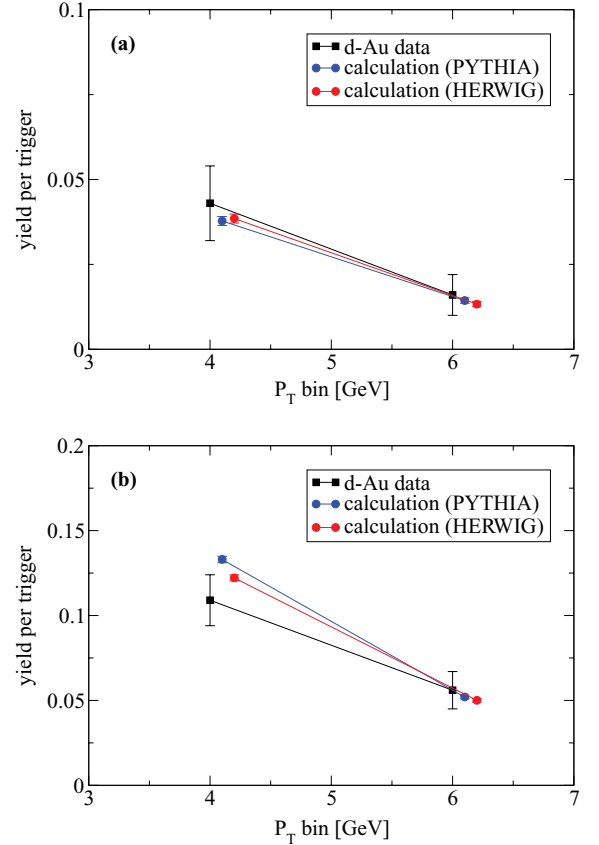


FIG. 1. (Color online) Comparison of the 200A GeV d -Au baseline measurement [21] for the (a) near-side and (b) away-side correlation strength with calculations utilizing PYTHIA [34] or HERWIG [35] codes for shower development and hadronization.

Sampling $A_1(z_1, p_T)$ for any parton that emerged with sufficient energy from the medium provides the energy of the two most energetic hadrons on both sides of the event. The harder of these two is by definition the trigger hadron and defines the near side. The hadron opposite to it is then the leading contribution to the away-side correlation.

The leading contribution to the correlation strength on the trigger side arises only in next-to-leading fragmentation, i.e., when sampling $A_2(z_1, z_2, p_T)$ for the (fixed) z_1 of the trigger to find the momentum of the second most energetic near-side hadron. This term provides a correction to the away-side correlation strength. This correction is P_T dependent and decreases from 25% in the lowest P_T bin in the present investigation (4–6 GeV) to 2.7% in the highest (10 + GeV) bin. Since we make a model comparison with the baseline on both the near and away sides, we compute both the leading term and the next-to-leading correction on the away side for consistency.

D. Hydrodynamic description of the medium

For the hydrodynamic description of the soft medium, we use the three-dimensional relativistic fluid dynamics model by Nonaka and Bass [17]. The starting point is the relativistic

hydrodynamic equation

$$\partial_\mu T^{\mu\nu} = 0, \quad (5)$$

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

$$T^{\mu\nu} = (\epsilon + p)U^\mu U^\nu - pg^{\mu\nu}. \quad (6)$$

Here ϵ , p , U , and $g^{\mu\nu}$ are energy density, pressure, four-velocity, and metric tensor, respectively. The relativistic hydrodynamic equation (5) is solved numerically using baryon number n_B conservation,

$$\partial_\mu (n_B(T, \mu)U^\mu) = 0, \quad (7)$$

as a constraint and closing the resulting set of partial differential equations by specifying an equation of state (EOS): $\epsilon = \epsilon(p)$. In Ref. [17], a bag model EOS is used for the quark-gluon plasma.

The code utilizes a Lagrangian mesh and light-cone coordinates (τ, x, y, η) ($\tau = \sqrt{t^2 - z^2}$) to optimize the model for the ultrarelativistic regime of heavy collisions at the BNL Relativistic Heavy Ion Collider (RHIC). It is assumed that hydrodynamic expansion starts at $\tau_0 = 0.6$ fm. Initial energy density and baryon number density are parametrized by

$$\begin{aligned} \epsilon(x, y, \eta) &= \epsilon_{\max} W(x, y; b) H(\eta), \\ n_B(x, y, \eta) &= n_{B\max} W(x, y; b) H(\eta), \end{aligned} \quad (8)$$

where b and $\epsilon_{\max}$ ($n_{B\max}$) are the impact parameter and the maximum value of energy density (baryon number density), respectively. $W(x, y; b)$ is given by a combination of wounded nuclear model and binary collision model, and $H(\eta)$ is given by

$$H(\eta) = \exp \left[\frac{-(|\eta| - \eta_0)^2}{2\sigma_\eta^2 \theta(|\eta| - \eta_0)} \right].$$

All parameters of the hydrodynamic evolution [17] have been fixed by a fit to the soft sector (elliptic flow, pseudorapidity distributions, and low- p_T single-particle spectra), thereby providing a fully determined medium evolution for hard probes to propagate through.

Note that the fluid description in principle also contains hard hadrons from the high- P_T tails of the thermal distributions. However, these contributions are small; and while the soft medium may influence hadron production between 4 and 6 GeV due to recombination processes, hadron production above 6 GeV is clearly dominated by hard processes [36–38]. Consequently, we neglect mechanisms of hadron production other than hard processes in the following.

III. RESULTS

We compute the suppression of the away-side correlation strength for a momentum range of 12–20 GeV for the trigger hadron. This choice is motivated by the results of Ref. [23], where we found that information about the medium density becomes apparent only if the spread between trigger momentum and lowest associate momentum bin is large enough to not only allow particles with no energy loss to be registered but also include contributions from partons after

finite energy loss. On the other hand, we found that the assumption that high- P_T hadron production can be modeled by jet fragmentation only holds for 6 GeV and above. Below 6 GeV, coalescence processes [36–38] play a significant role in hadron production in heavy-ion collisions.

This requirement essentially determines the choice of the associate momentum windows. We include the window from 4–6 GeV for comparison, stressing that we expect that the model underpredicts the yield in this region. The remaining bins are chosen as 6–8, 8–10, and 10 + GeV. To cancel systematic errors in the calculation of the baseline reaction without a medium, we present our results in the form of I_{AA} , i.e., the per-trigger yield in Au-Au collisions (for given b and ϕ) divided by the per-trigger yield in p-p collisions.

Since we do not find any significant modifications on the near side (cf. the discussion in Ref. [23]), we confine ourselves in the following to studying the modifications seen on the away side.

A. I_{AA} as a function of P_T , b , and ϕ

We show the resulting I_{AA} as a function of associate hadron momentum bin in Fig. 2 for four different values of impact parameter b for both in-plane ($\phi = 0$) and out-of-plane ($\phi = \pi/2$) emission. The first feature observed is that there is no strong P_T dependence in I_{AA} for any impact parameter, albeit a small rise may be present. The strongest trend seen is a rise in I_{AA} with b , corresponding to the fact that both entropy production (and hence medium density) and average path length are reduced for peripheral collisions as compared to central collisions. A growing split between in-plane and out-of-plane emissions as a function of b shows that the spatial asymmetry in the initial state of the hydrodynamic calculation is indeed translated into a ϕ -dependent suppression of the correlation strength.

We show the P_T dependence of the difference between in-plane and out-of-plane emissions in Fig. 3. There is some indication that the separation between the different angular

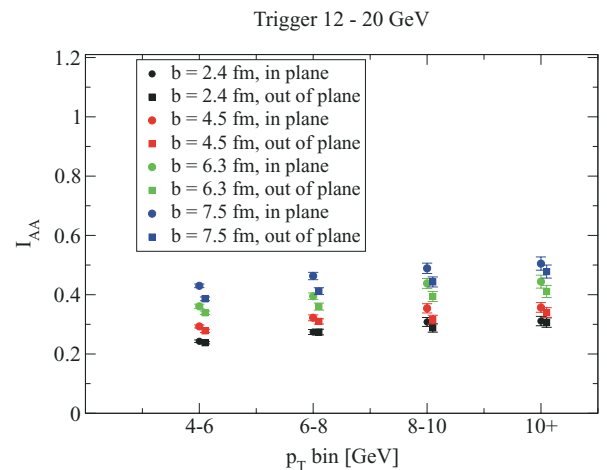


FIG. 2. (Color online) P_T dependence of charged hadron I_{AA} in 200A GeV Au-Au collisions for a 12–20 GeV trigger as a function of impact parameter b , shown for in-plane and out-of-plane emission.

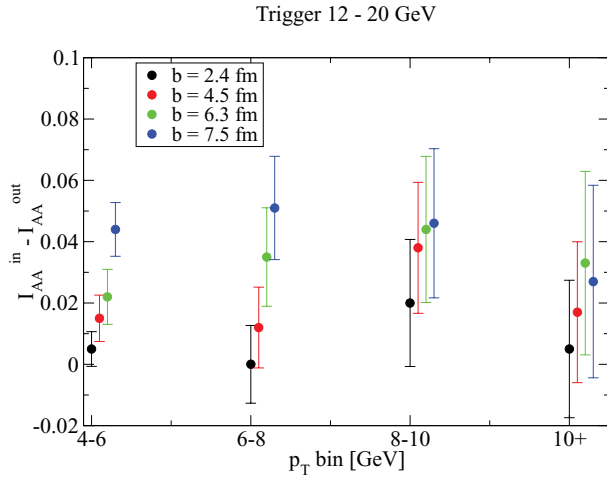


FIG. 3. (Color online) P_T dependence of the spread between in-plane and out-of-plane emissions in 200A GeV Au-Au collisions for a 12–20 GeV trigger as a function of impact parameter b .

orientations is more pronounced at lower P_T , although the significant reduction of MC statistics in the high- P_T bins as compared to the low- P_T bins poses difficulties to a solid statement.

Surprisingly, the magnitude of both I_{AA} and the split between in-plane and out-of-plane emissions as a function of b is rather similar to the values found for single-hadron suppression in terms of R_{AA} found within the same formalism [18]. However, R_{AA} and I_{AA} are in fact not the same; the similarity holds only on the level of about 30%. For this particular set of trigger and associate momenta, the magnitude of I_{AA} is larger than R_{AA} , whereas the spread between in-plane and out-of-plane emissions is consistently smaller than the spread found in R_{AA} .

B. Geometry of single-hadron and dihadron suppression

We show the probability density of vertices in the (x, y) plane leading to a near-side trigger hadron between 12 and 20 GeV in Fig. 4 for impact parameters $b = 2.4$ and 7.5 fm for both in-plane and out-of-plane emissions.

For the most central impact parameter $b = 2.4$ fm, the results of the simulation in the three-dimensional hydrodynamic medium evolution are very similar to the results obtained in Ref. [23] for central Au-Au collisions in a two-dimensional hydrodynamics simulation [39]. There is a clear surface effect visible: the most likely point of origin of a triggered hadron is not in the medium center but displaced about 4 fm and hence close to the surface. The detailed position of this maximum arises from a balance between the distribution of primary vertices coming from the nuclear overlap Eq. (2) which peaks in the medium center and from the fact that partons originating in a high-density region are more likely to be quenched. On the other hand, close to the medium edge where the quenching is small, the probability to find a vertex is much reduced.

For $b = 2.4$ fm, there is no large difference visible between in-plane and out-of-plane, corresponding to the fact that the spatial asymmetry of the medium is small. This is different

for the simulation with $b = 7.5$ fm. Here, clear changes in geometry when going from in-plane to out-of plane emission are visible. In particular, the surface effect is stronger for out-of-plane emission because of the greater amount of medium in the path of outgoing partons for this configuration.

Note that the distribution shown here is qualitatively (i.e., up to small changes with P_T) the same for any hard single-hadron observable, in particular R_{AA} (averaged over the trigger range) and γ -hadron correlations (where the photon escapes into the $+x$ direction).

We show the probability density of vertices in the (x, y) plane leading to a near-side trigger hadron between 12 and 20 GeV and an associate away-side hadron with $4 < p_T < 6$ GeV in Fig. 5. It is immediately obvious that the distribution is very different from the distribution of vertices for single-hadron observation shown in Fig. 4. First, the dihadron distributions are much wider in the $\pm y$ direction, indicating the importance of the periphery where both near- and away-side partons have a short in-medium path (or the halo where the production vertex lies outside the medium). Second, the distribution is almost symmetric with respect to the y axis, indicating the events are favored in which the near- and away-side path lengths are about equal. Both features are consistent with the distributions observed in Ref. [23] for different density evolution models. In particular, note the absence of tangential emission, which in the framework presented here does occur, albeit at much stronger quenching power of the medium [23].

It is evident from the underlying geometry that the physics (in terms of average path length and relevant medium density) of single-hadron and dihadron suppression is rather different in the calculation. Geometry alone clearly cannot account for the observed similarity between R_{AA} and I_{AA} .

IV. DISCUSSION

One may be tempted to conclude from the similarity of R_{AA} and I_{AA} that there is no additional information contained in I_{AA} beyond that which can be obtained from single-hadron suppression. However, this is clearly not the case. Let us illustrate in a qualitative way the meaning of this observation.

In a purely geometrical suppression picture, a parton is absorbed when the hard vertex of its origin is in a dense region but unmodified when the vertex falls into the dilute halo region. In this case, the whole back-to-back event is either absorbed or unmodified. This in turn implies that the number of triggered events is reduced, but once an event is triggered, there is no modification of the near- or away-side parton. From this, one may immediately deduce that in this scenario R_{AA} may be below unity, but I_{AA} will always be (modulo small effects by the nuclear parton distribution functions) equal to unity.

Thus, in any situation in which quenching is chiefly driven by the density at the vertex of origin and the medium can be modeled as an opaque and a dilute region, I_{AA} is expected to be substantially larger than R_{AA} . Within this model, this trend has been demonstrated as a rise in the per-trigger yield in the black core scenario described in Ref. [23].

Let us next consider the momentum space. By definition, the trigger hadron is the hardest hadron of the event. Consequently,

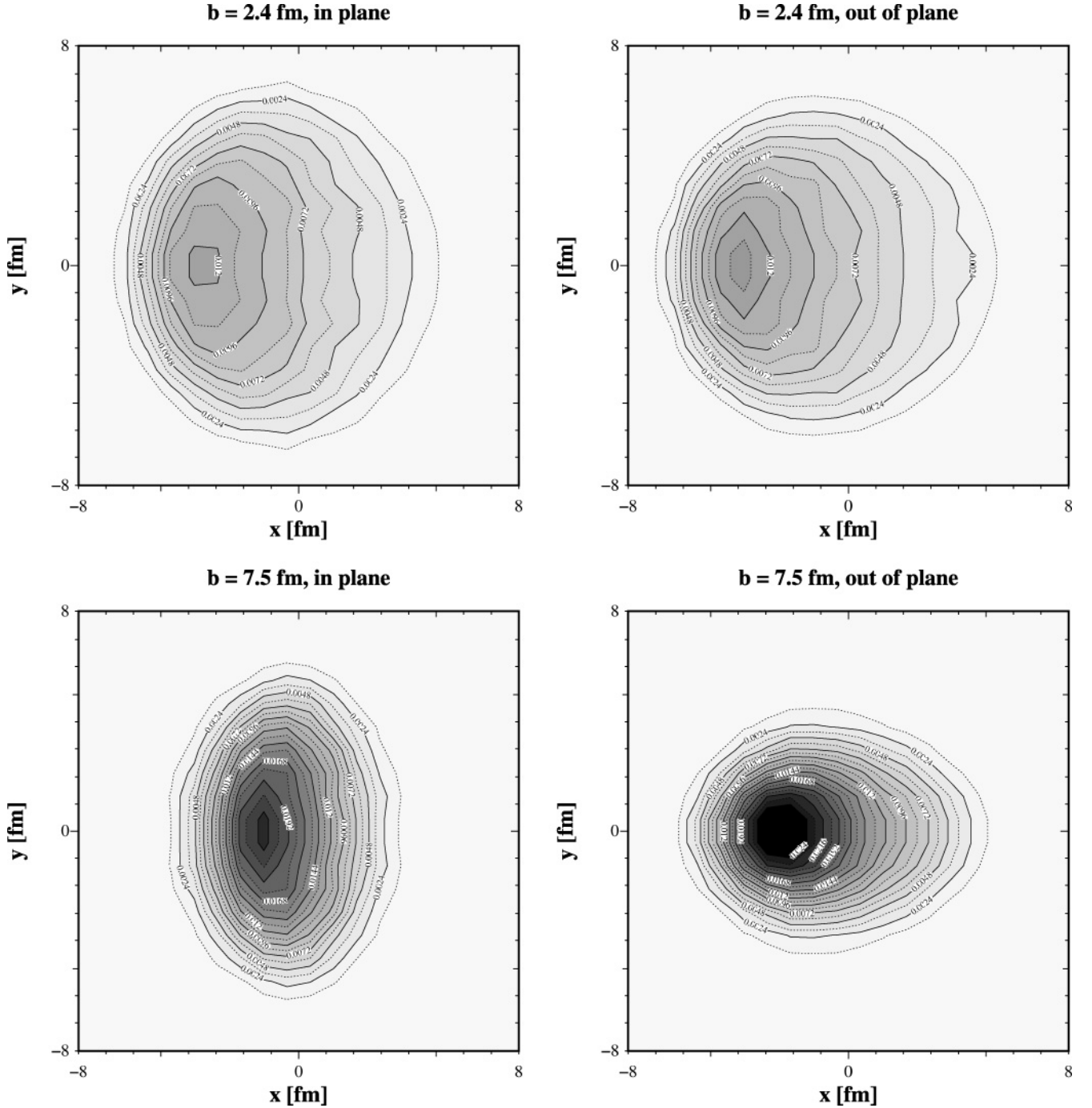


FIG. 4. Probability density for an inclusive hard hadron measurement, i.e., for finding a parton production vertex at (x, y) given a triggered event with $12 < p_T < 20$ GeV for different b and ϕ . In all cases, the near-side (triggered) hadron propagates to the $-x$ direction. Contours are at linear intervals.

there is the strongest bias on the trigger parton to experience no or only a small amount of energy loss. In contrast, the bias *given that there is a hard hadron on the near side* on the away-side hadron is considerably reduced, allowing average energy losses of several GeV on the away side [40]. If the geometry-averaged energy loss probability on the near and away sides were equal (i.e., if there were not a bias toward surface emission), this would imply again that I_{AA} should be

substantially larger than R_{AA} (for the measured R_{AA} of about 0.2, this would amount to almost a factor of 2 difference).

However, there is an opposing bias in position space. The fact that the hardest hadron tends to be produced close to the surface means that the away-side parton on average has a longer path length. Here, the weighting of the energy loss with path length is crucial, as the medium density drops as a result of the longitudinal and radial expansion of the medium while the

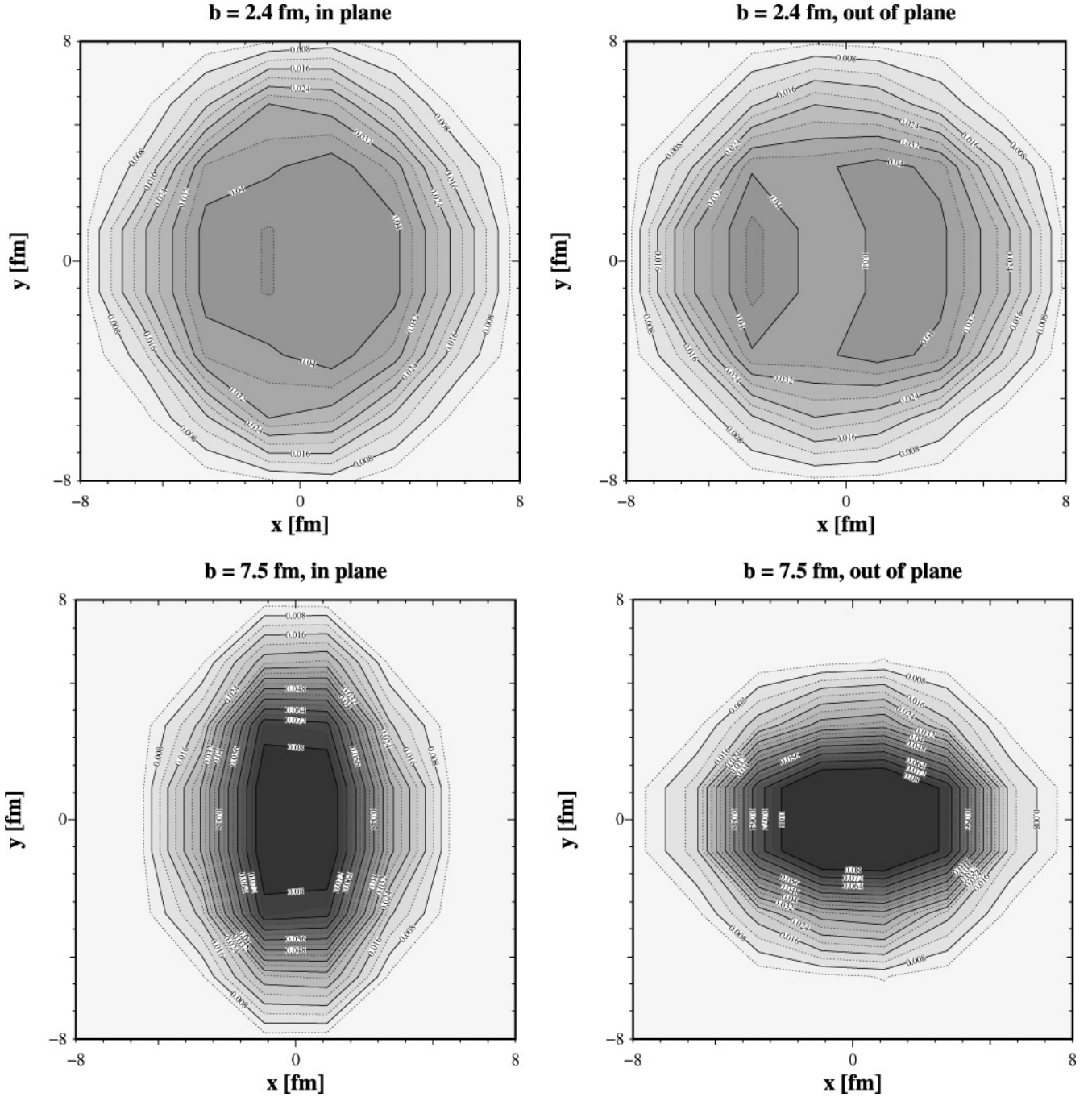


FIG. 5. Probability density for a coincidence hard hadron measurement, i.e., for finding a vertex at (x, y) leading to a triggered event with $12 < p_T < 20$ GeV and an away-side hadron with $4 < p_T < 6$ GeV for different b and ϕ . In all cases, the near-side hadron propagates to the $-x$ direction. Contours are at linear intervals.

parton traverses the medium. An energy loss that grows linear in path length L in a constant medium (as characteristic of elastic energy loss), for example, translates into a logarithmic path-length dependence $\sim \ln(L/L_0)$ in a medium undergoing longitudinal Bjorken expansion where the density drops like $1/\tau$. In this situation, path-length differences on the near and away sides do not matter much, and the bias in momentum space is expected to win out. This was observed in Ref. [24] within the model described here.

On the other hand, for an energy loss with quadratic path-length dependence $\sim L^2$ in a constant medium, the energy loss gets proportional to $\sim L$ in a medium undergoing longitudinal Bjorken expansion. If the away-side path is on average twice as long as the near-side path, this effect would be expected to cancel the reduced bias in momentum space and make R_{AA} and I_{AA} about equal. This argument strongly relies on modeling the expansion of the medium correctly. Note that in the static medium studied in Ref. [41]

in a conceptually similar framework using the same radiative energy loss picture, back-to-back events are almost completely suppressed, corresponding to an energy loss dependence with path length that is too strong.

These arguments may help us understand why the fact that I_{AA} and R_{AA} are of similar magnitude is in fact highly nontrivial and in itself probes properties of the medium and its expansion along with properties of the parton-medium interaction.

V. CONCLUSIONS

We have computed the variation of the suppression of back-to-back correlations for noncentral 200A GeV Au-Au collisions as a function of impact parameter b for both in-plane and out-of-plane emissions. Since the parameter determining the quenching power of the medium has been adjusted such that single-hadron R_{AA} for central collisions is described, the extension to back-to-back correlations for noncentral collisions and the dependence on the angle with the reaction plane are computed without additional free parameters, given the underlying hydrodynamic evolution (which is in turn constrained by the need to describe bulk matter observables such as momentum spectra and elliptic flow v_2). Thus, the results represent a rather constrained prediction which can be used to test the validity of the combination of hydrodynamic evolution [17] with the radiative energy loss model [2,33].

The suppression pattern we find is not unexpected given previous results [18,23]. As a function of associate hadron momentum, I_{AA} is approximately flat but may exhibit a small rise. I_{AA} increases with impact parameter, reflecting the reduced soft matter entropy production and a shortening of average in-medium path length. At the same time, a split between in-plane and out-of-plane emission grows, reflecting the spatial asymmetry in the initial state. There is an interesting similarity between the numerical values of I_{AA} (and the magnitude of the split) and the values of R_{AA} and the split between in-plane and out-of-plane emission observed there. There is no single physics reason that would dictate this outcome, rather one is looking at a cancellation of various different effects. An investigation in the geometry reveals that the physics underlying single-hadron and dihadron observables in the model is rather different in terms of dominant point of origin, typical density being probed, and average path length. A detailed comparison with data and computations for a different system size, e.g., Cu-Cu, may be helpful to investigating this further.

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