

Jet suppression and azimuthal anisotropy from RHIC to LHC

Yacine Mehtar-Tani*

*Physics Department, Brookhaven National Laboratory, Upton, New York 11973, USA*Daniel Pablos[†]*Instituto Galego de Física de Altas Enerxías IGFAE,
Universidade de Santiago de Compostela, E-15782 Galicia-Spain;
Departamento de Física, Universidad de Oviedo, Avda. Federico García Lorca 18, 33007 Oviedo, Spain
and Instituto Universitario de Ciencias y Tecnologías Espaciales de Asturias (ICTEA),
Calle de la Independencia 13, 33004 Oviedo, Spain*Konrad Tywoniuk[‡]*Department of Physics and Technology, University of Bergen, Allégaten 55, 5007 Bergen, Norway* (Received 15 April 2024; accepted 7 June 2024; published 8 July 2024)

Azimuthal anisotropies of high- p_T particles produced in heavy-ion collisions are understood as an effect of a geometrical selection bias. Particles oriented in the direction in which the QCD medium formed in these collisions is shorter suffer less energy loss, and thus, are over-represented in the final ensemble compared to those oriented in the direction in which the medium is longer. In this work we present the first semianalytical predictions, including propagation through a realistic, hydrodynamical background, of the elliptic azimuthal anisotropy for jets, obtaining a quantitative agreement with available experimental data as a function of the jet p_T , its cone size R , and the collisions centrality. Jets are multipartonic, extended objects and their energy loss is sensitive to substructure fluctuations. This sensitivity is determined by the physics of color coherence that relates to the ability of the medium to resolve those partonic fluctuations. Specifically, color dipoles with an angular separation smaller than a critical angle, θ_c , are not resolved by the medium and they effectively act as a coherent source of energy loss. We find that elliptic jet azimuthal anisotropy has a specially strong dependence on coherence physics due to the marked length dependence of θ_c . By combining our predictions for the collision systems and center-of-mass energies studied at RHIC and the LHC, covering a wide range of typical values of θ_c , we show that the relative size of elliptic jet azimuthal anisotropies for jets with different cone sizes R follows a universal trend that indicates a transition from a coherent regime of jet quenching to a decoherent regime. These results suggest a way forward to reveal the role played by the physics of jet color decoherence in probing deconfined QCD matter.

DOI: [10.1103/PhysRevD.110.014009](https://doi.org/10.1103/PhysRevD.110.014009)**I. INTRODUCTION**

Ultrarelativistic heavy-ion collisions taking place at present-day colliders provide insights into the dynamics of deconfined nuclear matter created in the aftermath of these collisions [1]. The remarkable correlation between the geometry of the initial collision state and the final state

momentum anisotropies of the observed particles has inspired the conjecture that this new state of matter, referred to as the quark-gluon plasma (QGP), behaves like an expanding droplet of liquid deconfined nuclear matter [2,3]. The surprising fact that this apparent hydrodynamic behavior is also observed in proton-nucleus and proton-proton collisions [4] has posed further challenges to our understanding of the applicability of hydrodynamics for smaller colliding systems, where spatial gradients are expected to be large. These challenges were put into a new perspective with the finding of hydrodynamic attractor solutions, present in a variety of microscopic descriptions, on which a longitudinally boosted system will fall provided that the nonhydrodynamic modes of the theory decay sufficiently fast [5–7] (for a recent review see Ref. [8]). While these developments support an effective description

*mehtartani@bnl.gov

†daniel.pablos@usc.es

‡konrad.tywoniuk@uib.no

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of deconfined Quantum Chromodynamics (QCD) matter using relativistic hydrodynamics, the details of the microscopic description of the system remain to be fully elucidated.

High-energy jets are rare events that are produced alongside the QGP with which they strongly interact while they travel through, offering valuable information about the properties of the medium. Due to their multiscale nature, spanning from the jet transverse momentum $\sim p_T$ to the nonperturbative scales of the vacuum $\sim \Lambda_{\text{QCD}}$, they probe the QGP at different length scales. Through the study of the modifications imprinted in this inmedium developing multipartonic system, we can infer the nature of the interaction between the energetic partons and the QGP. Moreover, jets that undergo significant modification or quenching will experience the absorption of a sizable portion of their energy and momentum by the medium. This occurs via dissipative processes, the study of which provides insights into the approach to thermal equilibrium as well as broader aspects of nonequilibrium QCD dynamics.

A direct connection between jet modifications and the fundamental interactions between the QGP and fast partons crucially relies on our understanding of jet fragmentation in these systems. From the point of view of a high-energy jet, the modest energy scales that characterize a droplet of QGP are typically small compared to the largest jet energy scales. This implies that medium modifications set in at intermediate length scales, preceded by a nearly vacuum-like evolution [9,10]. Nevertheless, as confirmed in experimental data, the impact of medium modifications is striking [11]. Jet yield suppression, typically quantified by the nuclear modification factor R_{AA} , stands as a paradigmatic illustration of jet quenching. This phenomenon is interpreted as arising from the interaction between energetic colored objects and deconfined QCD matter.

A complete description of the dynamics of deconfined QCD matter will however require establishing a consistent picture between the microscopic physics of the bulk of the system and that of its interaction with the jets that traverse it. Recent theoretical progress on our understanding of the intricate interplay between the jet and the medium scales is now allowing for first-principles phenomenological applications which can be meaningfully confronted to the high-precision data expected from upcoming runs at RHIC and LHC.

In a previous publication [12] we provided the first analytical description of jet yield suppression as a function of the jet size R . In short, this description accounts both for the elastic and radiative processes induced by the interaction of a fundamental colored object, a parton, with the medium, as well as the process of resolving which partons participate in the medium interactions during the jet fragmentation. Our analytical calculations were convolved with a realistic model for the hydrodynamical background and compared to experimental data at LHC energies.

We performed a thorough study of the relative sizes of the different sources of theoretical uncertainties and established that these are dominated by the computation of the resolved phase space of the collinear invacuum parton cascade up to relatively large angles $R \sim 0.6$, which is systematically improvable within perturbative QCD (pQCD). In the present work, we extend this framework to the study of jet suppression at RHIC energies. This extension serves to further validate the fundamental working assumptions based on factorizing jet production, fragmentation, and subsequent medium modification.

We also address another paradigmatic observable that is jet azimuthal anisotropy. Roughly speaking, while R_{AA} is sensitive to the average path length traveled by the jet, this observable is sensitive to path-length differences in jet suppression due to the relative orientation in the transverse plane of the jet direction ϕ with respect to the event plane of the collision.

Analogously to how azimuthal anisotropies are quantified for the particles belonging to the bulk of the collision, jet azimuthal anisotropies are defined via the p_T -dependent Fourier coefficients of their azimuthal distribution, the $v_n(p_T)$, as

$$\frac{dN}{d\phi dp_T} = \frac{dN/dp_T}{2\pi} \left[1 + 2 \sum_n v_n(p_T) \cos(n(\phi - \Psi_n)) \right], \quad (1)$$

where dN/dp_T is the multiplicity in the given p_T -bin and Ψ_n defines the n th order event plane. The behavior of p_T -inclusive v_n , dominated by soft particle production from the bulk of the system, are understood in terms of the preferred directions determined by the initial geometry of the collision, which translate into momentum anisotropies due to the conversion of pressure gradients into momentum flow via the final state hydrodynamic evolution [13]. In contrast, high- p_T v_n are caused by a selection bias effect that results into larger yields in the directions in which the high- p_T object has had to traverse less amount of medium, i.e., in the directions with least quenching [14–16].

Due to the dominant elliptical, almondlike shape characteristic of the collision of two fairly round objects at finite impact parameter b , the largest v_n is typically the elliptical flow coefficient, v_2 . Nevertheless, fluctuations in the positions of the incoming nucleons inside the nuclei [17], and subnucleonic fluctuations in the color field configurations inside each of the nucleons at the moment of the collision [18], will give rise to finite values for the higher harmonics v_3 , v_4 , and so on. For the case of quenching-induced flow coefficients, also fluctuations in the process of energy loss will contribute to the value of all flow harmonics [19].

There is extensive recent literature on the study of flow coefficients for single energetic particles, or high- p_T hadrons (see, e.g., [19–26]). However, much less work has been devoted to the study of the flow coefficients for

jets, and most of what is present is obtained from Monte Carlo jet quenching models [27–29]. Given that jets are extended objects, consisting of a collection of hadrons within a given cone R , they are affected by a number of physical mechanisms that are not present in the description of single-hadron observables. Our analytical theoretical framework consistently combines the most dominant of such mechanisms, rendering it appropriate to address the more challenging scenario involving jets.

Even within the medium, high-energy jets experience part of their evolution as if they were approximately in vacuum [10]. This is due to the presence of a wide-scale separation between the typical time of high-virtuality parton splitting, as determined by the Dokshitzer–Gribov–Lipatov–Altarelli–Parisi (DGLAP) evolution equations, and the typical time to produce a medium-induced emission via multiple soft scatterings, as determined by the BDMPS-Z equations [9,10,30]. The DGLAP evolution resums large logarithms of the cone angle R , i.e., $\sim \log 1/R$, and is responsible for jet “energy loss” in vacuum from emissions off an initial hard parton with angles larger than R . Hence, the jet spectrum becomes a function of the cone angle R and transverse momentum p_T [31–33]. We incorporate this QCD evolution in the present work as well.

Vacuumlike radiation taking place within the cone R will no longer contribute to energy loss in vacuum at leading-logarithmic (LL) accuracy. However, in the presence of a medium they will determine the amount of color charges that can source out-of-cone radiation via interactions with that medium. The enhancement of multiple color charge energy loss by a Sudakov double logarithm is attributed to the distinction in energy loss between real and virtual contributions. Consequently, a critical aspect in the understanding of jet energy loss lies in determining the size of the vacuum-generated in-cone phase space resolved by the medium, which we refer to as the “quenched” phase space.

The first observations of the close relation between the size of the “quenched” phase space and the amount of energy loss were obtained using phenomenological models. These were able to show that the origin of the jet core narrowing observed in data is due to a selection bias towards those jets that experienced a narrower fragmentation during its vacuumlike evolution [34,35]. Even though each individual jet does get broadened in the process of energy loss, the final ensemble is biased towards those jets that lost the least energy, the narrower ones, due to the steepness of jet production spectra. This same phenomenon has been shown to account for a large part of the dijet asymmetry modification of lead-lead (PbPb) compared to proton-proton (pp) [36], even in the absence of path-length differences between the leading and subleading jet.

On the theory side, the role of hard radiation that occurs immediately after the hard partonic collision in sourcing further (medium-induced) emissions was elucidated in a series of papers analyzing the so-called *antenna*

setup [37–41]. These calculations revealed the existence of the so-called decoherence time, associated with the typical time needed by medium interactions before being able to resolve the individual colors of the created pair of partons. The case when this time is comparable to the medium length determines a critical angle θ_c , corresponding to a minimal angle below which jet splittings act coherently as a single color charge and which profoundly affects the quenching of jets in the medium [42]. For a homogeneous plasma, the critical angle is defined as

$$\theta_c \sim \frac{1}{\sqrt{\hat{q}L^3}}, \quad (2)$$

where L is the medium length and \hat{q} is the jet transport coefficient that describes Brownian motion in transverse momentum space. This phenomenon, called “color decoherence,” leads to a modified joint probability of losing energy off an initially color correlated pair of particles [43]. On the level of the jet spectrum in heavy-ion collisions, accounting for early resolved vacuumlike emissions inside the medium and accounting for their energy loss leads to an enhanced suppression factor at high- p_T where the additional “Sudakov factor” scales with the “quenched” phase space [9].

The physics of color decoherence explains why single inclusive hadron R_{AA} is larger than jet R_{AA} [9], i.e., $R_{AA}^{\text{jet}}(p_T) < R_{AA}^{\text{hadron}}(p_T)$, and how this fact is intimately related to the high- z enhancement of the measured jet fragmentation function modification [44,45]. The first phenomenological implementation of the inclusion of finite resolution effects, or to which degree do two colored charges generated by vacuumlike evolution engage with the medium independently, in a Monte Carlo jet quenching shower algorithm was done in [46], and later in close analogy to the theoretical description described above in [10]. The amount of selection bias towards narrower jets has been shown to be strongly affected by the size of the resolution parameter, leaving clear imprints on the distribution of the relatively hard, measurable splitting selected by the SoftDrop [47] grooming procedure (see Refs. [48,49]) or the Dynamical Grooming [50] procedure (see Refs. [51,52] as well as [53] for a closely related substructure observable). Lately, renewed interest in the energy-energy correlator observable [54–56] has arisen in the field of jet quenching [57–61] due, in part, to its potential sensitivity to emergent medium scales such as θ_c .

The present work shows how the R dependence of jet v_2 closely relates to the existence of this medium resolution scale, providing additional tools with which to experimentally constrain these physics and build a coherent picture of jet energy loss across all observables. With a full-fledged embedding of our theoretical framework into a realistic heavy-ion environment, we show that we can describe currently available jet v_2 data with a very good agreement.

Further, by studying the R dependence of this observable, we find an interesting departure from the one found for jet R_{AA} , consisting in a clear and persistent ordering of jet v_2 as a function of the jet R across kinematic and centrality selections. With simple analytical estimates we are able to pinpoint the origin of this new phenomenon, which is rooted in the marked length dependence of the medium angular resolution scale, θ_c , combined with the fact that jet azimuthal anisotropy is a length-differential jet suppression observable. Moreover, we exploit centrality evolution to target different values of the typical decoherence angle θ_c , and in this way establish an arrangement of predictions which actually collapse into universal, analytically understandable curves, when presented in terms of the ratio of the two most relevant angular scales of the system, i.e., θ_c/R .

The rest of the paper is organized as follows. In Sec. II we present our theoretical formalism and explain how we embed it into a realistic heavy-ion environment. In Sec. III we present the results for jet R_{AA} and jet v_2 at RHIC and LHC energies as a function of R using our full semi-analytical framework. Next, in Sec. IV, we gain further insight on the phenomenological implications of the key ingredients of our computation by analytically studying a simplified scenario that allows a transparent qualitative interpretation of the full results of Sec. III. Finally, in Sec. V we conclude and look ahead.

II. THEORETICAL FORMALISM

In this section we provide a description of the theoretical formalism used to produce the full results of Sec. III. These ingredients are essentially the same as those used to produce the results of publication [12], which we reproduce, often in a more detailed form, for convenience of the reader. The key elements are illustrated in Fig. 1.

A. Vacuum-induced jet energy loss

The cross section to measure a jet with transverse momentum p_T contained within a cone of extent R in the (η, ϕ) angular plane can be computed in vacuum as [31]

$$\sigma^{pp}(p_T, R) = \sum_{k=q,g} f_{\text{jet}/k}^{(n-1)}(R|p_T, R_0) \hat{\sigma}_k(p_T, R_0), \quad (3)$$

where R_0 is an initial large angular scale ($R_0 \sim 1$), and $n \equiv n_k(p_T, R_0)$ is the power-index of the cross section of the hard parton with flavor k , $\hat{\sigma}_k(p_T, R_0)$. The latter is calculated at leading order (LO) at the factorization scale Q_{fac}^2 via the convolution of the parton distribution functions (PDFs) $f_{i/p}(x, Q_{\text{fac}}^2)$ with the $2 \rightarrow 2$ QCD scattering cross section $\hat{\sigma}_{ij \rightarrow kl}$, i.e., $\hat{\sigma}_k = f_{i/p} \otimes f_{j/p} \otimes \hat{\sigma}_{ij \rightarrow kl}$. The moment of the fragmentation function of an initial hard parton with flavor k is

$$f_{\text{jet}/k}^{(n)}(R|p_T, R_0) \equiv \int_0^1 dx x^n f_{\text{jet}/k}(R, x|p_T, R_0), \quad (4)$$

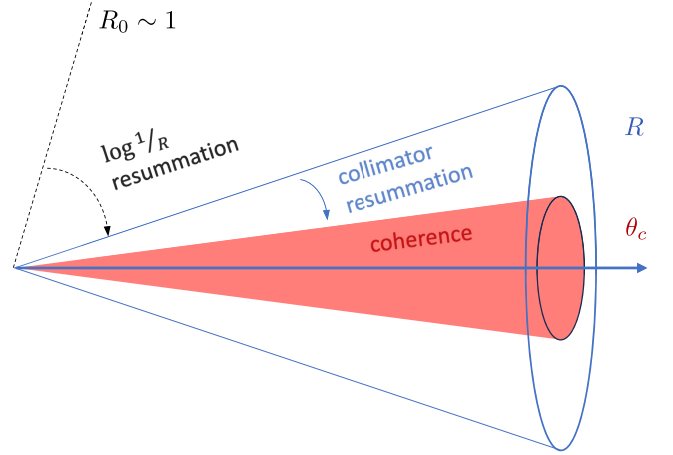


FIG. 1. Sketch of the main steps involved in our computation. Starting from a large angular scale R_0 , we evolve the jet spectrum down to a given cone R using the microjets formalism. This evolution can lead to out-of-cone radiation due to vacuum physics. Depending on the amount of in-medium resolved charges between R and the critical angle θ_c , the collimator quantifies the shift in the jet spectrum for a given R due to medium-induced effects.

receiving both quark and gluon contributions, i.e., $f_{\text{jet}/k}^{(n)} = \sum_{i=q,g} f_{i/k}^{(n)}$, via flavor conversion during the DGLAP evolution [31–33,62]. Here, $f_{\text{jet}/k}(R, x|p_T, R_0)$ is the probability of finding a subjet carrying momentum fraction x at cone angle R within a bigger jet with cone angle R_0 and transverse momentum p_T .

The full, inclusive jet cross section therefore becomes $\sigma^{pp}(p_T, R) = \sum_{i=q,g} \sigma_i^{pp}(p_T, R)$, where

$$\sigma_i^{pp}(p_T, R) = \sum_{k=q,g} f_{i/k}^{(n-1)}(R|p_T, R_0) \hat{\sigma}_k(p_T, R_0) \quad (5)$$

is the spectrum for a final quark or gluon jet measured at cone angle R .

The DGLAP evolution is what describes the microjet spectrum at different angular scales R , effectively implementing out-of-cone radiation, or in other words, vacuum jet energy loss. Its evolution equations in terms of the moments of the fragmentation functions are

$$\begin{aligned} \frac{d}{dt} f_q^{(n)}(t) &= \gamma_{qq}^{(n)}(t) f_q^{(n)}(t) + \gamma_{qg}^{(n)}(t) f_g^{(n)}(t) \\ \frac{d}{dt} f_g^{(n)}(t) &= \gamma_{gg}^{(n)}(t) f_g^{(n)}(t) + \gamma_{gq}^{(n)}(t) f_q^{(n)}(t), \end{aligned} \quad (6)$$

where the evolution variable $t \equiv \ln \theta/R$, with $\theta > R$ and $R < R_0$, so that t is evolved from $\ln R_0/R$ to 0. The anomalous dimensions at LO are [31]

$$\begin{aligned}
\gamma_{qq}^{(n)}(t) &= \int_0^1 dz \frac{\alpha_s(t)}{\pi} p_{qq}(z)(z^n - 1), \\
\gamma_{qg}^{(n)}(t) &= \int_0^1 dz \frac{\alpha_s(t)}{\pi} p_{qg}(z)z^n, \\
\gamma_{gg}^{(n)}(t) &= \int_0^1 dz \frac{\alpha_s(t)}{\pi} [p_{gg}(z)(z^n - z) - p_{gq}(z)], \\
\gamma_{gq}^{(n)}(t) &= \int_0^1 dz \frac{\alpha_s(t)}{\pi} p_{gq}(z)z^n,
\end{aligned} \tag{7}$$

with the unregularized Altarelli-Parisi (AP) splitting functions

$$\begin{aligned}
p_{qq}(z) &= C_F \frac{1+z^2}{1-z}, \\
p_{qg}(z) &= C_F \frac{1+(1-z)^2}{z}, \\
p_{gg}(z) &= C_A \frac{(1-z(1-z))^2}{z(1-z)}, \\
p_{gq}(z) &= \frac{N_f}{2} (z^2 + (1-z)^2).
\end{aligned} \tag{8}$$

Finally, the running coupling $\alpha_s(t) = \alpha(k_t)$, with $k_t = z(1-z)p_T\theta$, is evaluated at LO with five active flavors and regularized as $\alpha_s(k_t) = \min[1, 2\pi/(\beta_0 \log k_t/Q_0)]$, with $\beta_0 = 23/3$ and $Q_0 = 0.09$ GeV.

In practice, we fit the shape of the initial spectra at $R = R_0 = 1$, for a given pseudorapidity η range, using the event generator PYTHIA8 [63,64], parametrized as $d\hat{\sigma}^{(k)}/dp_T = \sigma_0^{(k)} (p_{T,0}^{(k)}/p_T)^{n^{(k)}(p_T)}$ and $n^{(k)}(p_T) = \sum_{i=0}^5 c_i^{(k)} \log^i(p_{T,0}^{(k)}/p_T)$. This effectively includes the LO $2 \rightarrow 2$ QCD scattering cross section convoluted with the PDFs [free-proton PDFs in pp collisions, nuclear PDFs in nucleus-nucleus (AA) collisions], plus some large-angle evolution down to R_0 . We then use the LO DGLAP evolution equations (6) to obtain the spectra for $R < 1$ using Eq. (3), a step which we refer to as “log 1/ R resummation” in the sketch depicted in Fig. 1. In pp collisions, and to LL accuracy, this completes the jet spectra computation.

B. Medium-induced jet energy loss

In the presence of deconfined matter, interactions of the colored jet charges with the medium constituents induce further out-of-cone energy loss contributions to the jet spectra. The size of these contributions depends on the interplay between the partonic structure contained within cone R and the QGP properties. Accounting for these medium effects requires extending angular resummation to scales smaller than R , as described in the present subsection.

The cross section to measure a jet with momentum p_T contained within a cone R in AA collisions is expressed as

$$\begin{aligned}
\sigma^{AA}(p_T, R) &= \sum_{i=q,g} \int_0^\infty d\epsilon P_i(\epsilon, R, p_T) \tilde{\sigma}_i^{pp}(p_T + \epsilon, R), \\
&= \sum_{i=q,g} Q_i(p_T, R) \tilde{\sigma}_i^{pp}(p_T),
\end{aligned} \tag{9}$$

where $\tilde{\sigma}_i^{pp}$ is the quark/gluon contribution to the total cross section at angular scale R , as computed in Eq. (3) [the tilde serves as a reminder that the proton PDFs are replaced by the nuclear PDFs EPS09 computed at LO [65], i.e., $f_{i/p} \rightarrow f_{i/A}$ in Eq. (3)], and $P_i(\epsilon)$ represents the probability distribution of losing energy ϵ out of the jet cone. We define the flavor dependent resummed quenching factor (QF) as [66]

$$Q_i(p_T, R) \equiv \int_0^\infty d\epsilon P_i(\epsilon, R, p_T) \frac{\tilde{\sigma}_i^{pp}(p_T + \epsilon)}{\tilde{\sigma}_i^{pp}(p_T)}. \tag{10}$$

In the limit of large power index n , we can use the following asymptotic expansion, valid for $\epsilon \ll p_T$ and $n\epsilon \sim p_T$,

$$\begin{aligned}
\frac{\tilde{\sigma}^{pp}(p_T + \epsilon)}{\tilde{\sigma}^{pp}(p_T)} &= \frac{p_T^{n(p_T)}}{(p_T + \epsilon)^{n(p_T + \epsilon)}} \\
&= e^{-\nu(p_T)\epsilon} [1 + \mathcal{O}(\nu_2(p_T)\epsilon^2)],
\end{aligned} \tag{11}$$

where the flavor subscript has been omitted for clarity, and where $\nu(p_T) = \frac{d \ln \tilde{\sigma}^{pp}}{d p_T}$ and $\nu_2(p_T) = \frac{d^2 \ln \tilde{\sigma}^{pp}}{d p_T^2}$. For our purposes, we will approximate $\nu(p_T) \approx n(p_T)/p_T$. Then, to first nonvanishing order in this expansion [67], one can identify the QF with the Laplace transform of the energy loss probability $P_i(\epsilon, R)$,

$$\begin{aligned}
Q(p_T, R) &\simeq \tilde{P}\left(\frac{n}{p_T}, R\right), \\
&\equiv \int_0^\infty d\epsilon P(\epsilon, R) e^{-\nu\epsilon} \Big|_{\nu=n/p_T}.
\end{aligned} \tag{12}$$

The factorization (9) reduces trivially to the jet production cross section in the absence of final-state interactions, Eq. (3), when $P(\epsilon) = \delta(\epsilon)$ (hence $Q = 1$) and by replacing the nuclear PDFs by proton PDFs.

The amount of quenching experienced by a quark/gluon jet of energy E and angular extent R , $Q_i(E, R)$ depends on the amount of sources of energy loss it contains. These sources are understood as the jet substructure fluctuations induced by the vacuumlike DGLAP evolution at scales below R . The finite spacetime extent of the medium requires knowledge of the spacetime evolution of these substructure fluctuations as well. A key ingredient in this picture is the quantum-mechanical formation time of an emission process, which essentially corresponds to the time it takes for an emission to decohere from its emitter, e.g., when $|qg\rangle \rightarrow |q\rangle + |g\rangle$ for the case of gluon emission off a quark. This formation time can be estimated using the

Heisenberg uncertainty principle to be $t_f \sim (E/Q)/Q$ [69], where Q is the virtuality of the process (not to be confused with the QF Q_i) and E/Q is a boost factor to the comoving frame of the particle. More explicitly, when considering the transverse momentum squared of a splitting event that yields two offsprings with longitudinal momenta zE and $(1-z)E$, respectively, which reads $k_\perp^2 = z(1-z)Q^2$, its formation time can be expressed as $t_f \sim z(1-z)E/k_\perp^2$.

A given fluctuation will only contribute to the QF if it is formed within the medium, characterized with an extent L , so $t_f < L$. Another requirement arises from the physics of color coherence within the medium. By computing stimulated emission off a color connected dipole traversing deconfined matter, it was found that independent emissions off each of the legs of the antenna could only happen after a time t_d [38,39,41,70], the decoherence time. This is the time it takes for the members of a dipole to lose their color correlation via color rotations induced by interactions with the medium constituents. The decoherence time t_d defines the time when the size of the dipole $r_\perp(t) = \theta t$ is of order the inverse transverse momentum scale accumulated by multiple scatterings during t , i.e., $Q_\perp(t) \equiv \sqrt{\hat{q}t}$. Or equivalently, when $r_\perp(t)Q_\perp(t) \sim 1$. Solving for t reads the decoherence time, $t_d \sim (\hat{q}\theta^2)^{-1/3}$. A given fluctuation will only contribute if $t_d < L$. Also, naturally, one requires $t_f < t_d$ for the splitting to be considered vacuumlike. When $t_f > t_d$ the splitting is medium induced unless it is produced outside the QGP.

The nonlinear angular evolution that couples the resummed QF of quarks and gluons is then given by [9]

$$\frac{\partial Q_i(p, \theta)}{\partial \ln \theta} = \int_0^1 dz \frac{\alpha_s(k_\perp)}{2\pi} p_{ji}(z) \Theta_{\text{res}}(z, \theta) \times [Q_j(zp, \theta) Q_k((1-z)p, \theta) - Q_i(p, \theta)], \quad (13)$$

where the resolved phase space constraint encapsulates the requirements stated above, i.e., $\Theta_{\text{res}}(p, R) = \Theta(t_f < t_d < L)$, where $t_f = 2/(\omega\theta^2)$ and $t_d = (\hat{q}\theta^2/12)^{1/3}$. The requirement $t_d < L$ can be translated into an angular constraint, as $\theta > \theta_c$, where the critical angle $\theta_c = (\hat{q}L^3/12)^{-1/2}$ is the minimum angle a dipole can have such that it is resolved by the medium [71]. In terms of angular scales, Eq. (13) encodes the ‘‘collimator resummation’’ evolution depicted in Fig. 1 at scales below R . This evolution has no effect for angles below θ_c , labeled as ‘‘coherence’’ in the sketch.

Even if a jet had a trivial inmedium evolution, for instance if $\theta_c > R$, it would still lose energy, but only from the source associated with the total color charge of the multipartonic system. This corresponds to the initial conditions of Eq. (13), and are called the bare quenching factors, computed as $Q_i(p, 0) = Q_{\text{rad},i}^{(0)}(p_T) Q_{\text{el},i}^{(0)}(p_T)$, for radiative and elastic energy loss mechanisms, respectively [66,73]. The radiative factor is related to the Laplace transform of the quenching weight, and given by

$$\tilde{P}(\nu) = \exp \left[- \int_0^\infty d\omega \frac{dI}{d\omega} (1 - e^{-\nu\omega}) \right], \quad (14)$$

where $dI/d\omega$ represents the spectrum of medium-induced emissions and $Q_{\text{rad}}^{(0)}(p_T) \equiv \tilde{P}(\nu = n/p_T)$. We discuss both contributions in detail below, cf. Eqs. (20) and (21) for final results.

Let us continue with a brief reminder of the framework used in the present and previous work [12] to compute medium-induced radiation. It is based on the so-called improved opacity expansion (IOE) scheme, where the relatively rare hard momentum exchanges are computed as perturbations on top of a sea of frequent multiple soft scatterings. It amounts to expanding around the harmonic oscillator. This framework naturally incorporates the BDMPS-Z (multiple soft scattering) and GLV [74,75] (single hard scattering) limits of the radiation spectrum, while it leaves out the Bethe-Heitler (BH) regime (corresponding to soft emissions with $\omega \lesssim T$) [76].

When the medium-induced emission is soft compared to the parent parton, i.e., $\omega \ll E$, the radiation spectrum can be expressed as [80–82] $dI_{\text{NHO}}/d\omega = dI^{(0)}/d\omega + dI^{(1)}/d\omega$, where ‘‘NHO’’ (or next-to-harmonic oscillator approximation) refers to the order of the expansion in the IOE, and read

$$\frac{dI^{(0)}}{d\omega} = \frac{2\alpha_s C_R}{\pi\omega} \ln |\cos \Omega L|, \quad (15)$$

$$\frac{dI^{(1)}}{d\omega} = \frac{\alpha_s C_R \hat{q}_0}{2\pi} \text{Re} \int_0^L ds \frac{-1}{k^2(s)} \ln \frac{-k^2(s)}{Q^2 e^{-\gamma_E}}. \quad (16)$$

In the last expression, $\Omega = (1-i)\sqrt{\hat{q}/(4\omega)}$, $k^2(s) = i\omega\Omega[\cot \Omega s - \tan \Omega(L-s)]/2$, and the strong coupling constant runs with the typical transverse momentum of the emission, as $\alpha_s = \alpha_s((\hat{q}\omega)^{1/4})$. These expressions are formally valid when the medium density is constant. We will nevertheless make use of them with rescaled medium parameters to account for the dynamical expansion; see Sec. II D (for a discussion of the scaling properties, see Refs. [83,84]).

In the IOE, the effective transport coefficient \hat{q} acquires logarithmic corrections associated with the matching scale Q , differing from the bare \hat{q}_0 as

$$\hat{q} = \hat{q}_0 \ln \frac{Q^2}{\mu_*^2}, \quad (17)$$

where Q naturally corresponds to the transverse momentum acquired by the emission via medium interactions during its formation time. This scale is found by solving the implicit equation $Q^4 = \hat{q}_0 \omega \ln Q^2/\mu_*^2$, where the cutoff scale is $\mu_*^2 = m_D^2 \exp[-2 + 2\gamma_E]/4$ [81,82]. The bare jet transport coefficient that follows is obtained within the Hard

Thermal Loop (HTL) effective theory, and reads $\hat{q}_0 = g_{\text{med}}^2 N_c m_D^2 T / (4\pi)$, while the expression for the Debye screening mass m_D , for three active flavors, reads $m_D^2 = 3g_{\text{med}}^2 T^2 / 2$. Finally, the coupling g_{med} between the jet partons and the medium constituents is taken to be fixed, and corresponds to the only free parameter of our framework.

While Eqs. (15) and (16) capture the ω -distribution of the emitted quanta correctly in the regimes of interest, they have been integrated over transverse momentum k_\perp and offer no information about the broadening they experienced. In order to compute jet energy loss, one needs additional knowledge about their final angular distribution with respect to the jet axis, $\theta \sim k_\perp / \omega$, since only those emissions that end up outside of the jet cone R contribute to the QF. To that end, we have adopted a multiplicative Ansatz that depends on the regime in frequency of the emission: hard or soft.

We distinguish gluons pertaining to three broad regimes. First, there are gluons with energies $T < \omega < \omega_s$, where $\omega_s \equiv (g_{\text{med}}^2 N_c / (2\pi)^2) \pi \hat{q}_0 L^2$, for which formally $dI^{(0)}/d\omega \gg dI^{(1)}/d\omega$ in the IOE expansion. These gluons are produced with $\mathcal{O}(1)$ probability during the passage of the parent parton through the medium and trigger an inverse turbulent cascade via democratic branchings that results into an efficient accumulation of modes around $\omega \sim T$ [85–88]. Emissions pertaining to this regime are then assumed to have been thermalized, becoming part of the medium and contributing as correlated background. After the subtraction of the uncorrelated background, as it is typically done in experiments, their contribution to the final energy distribution will be that produced by the wake generated via the perturbation of the relativistic hydrodynamic equations of motion that govern the evolution of a liquid system such as the QGP. While realistically modeling this nonperturbative piece of dynamics is beyond the scope of the present work, we estimate its effects on the QF by considering the possibility that some of this energy can be recaptured. The larger the jet cone, the more likely it is to recapture some of the energy.

Focusing first on the turbulent regime, corresponding to $T < \omega < \omega_s$, the amount of lost energy *at any angle* can be estimated as

$$\Delta E_{\text{turb}} = \int_T^{\omega_s} d\omega \omega \frac{dI^{(0)}}{d\omega}. \quad (18)$$

Now, assuming that a fraction ξ of this energy gets thermalized and redistributed back inside the cone, just $(1 - \xi)\Delta E_{\text{turb}}$ is actually lost. Using geometrical arguments this ratio ξ should correspond to the ratio of the area covered by the jet cone and the area where the energy is distributed, arriving at $\xi \approx R^2 / R_{\text{rec}}^2$, where the parameter R_{rec} quantifies how efficient opening up the cone is in recapturing the thermalized energy. Assuming a flat angular

distribution across the hemisphere of a given jet would correspond to $R_{\text{rec}} = \pi/2$, while a narrower angular distribution, as inspired by linearized hydrodynamics wake calculations in the absence of radial flow, would correspond to $R_{\text{rec}} < \pi/2$. Inspired by the study in [12], we have currently chosen $R_{\text{rec}} = (5/6)\pi/2$. In our previous publication we checked that the choice of this nonperturbative parameter has very little influence on jet suppression up to $R \lesssim 0.6$, even for extreme values of R_{rec} , specially when compared to other sources of uncertainties, such as the determination of the resolved phase space (which is perturbatively calculable). For this reason, and for jet observables up to $R \lesssim 0.6$, g_{med} still remains as the only relevant free parameter.

In order to realize this smearing, we note that $\Delta E = -\frac{d}{d\nu} \tilde{P}(\nu)|_{\nu=0}$; see Eq. (14). Hence, we define a “turbulent” quenching factor $Q_{\text{rad,turb}}^{(0)}(p_T) = \tilde{P}_{T < \omega < \omega_s}(\nu(1 - R^2/R_{\text{rec}}^2))$, where we have indicated that the integral over $d\omega$ in (14) is constrained by the thermal scale and ω_s . For reference, the full bare quenching factor is given in Eq. (20). In summary, emissions between $T < \omega < \omega_s$ are assumed to thermalize quickly, and their angular distribution is approximately flat in the hemisphere of the jet.

Second, semihard gluons are those emitted with $\omega_s < \omega \lesssim \omega_c$, where the critical energy $\omega_c \equiv \hat{q}_0 \ln(\hat{q}_0 L / \mu_*^2) L^2 / 2$ is the maximal energy that a medium-induced emission produced by multiple soft scatterings can have. Third, harder gluons, with $\omega > \omega_c$, can be produced via rare single large momentum transfers and we have now $dI^{(0)}/d\omega \ll dI^{(1)}/d\omega$. Semihard gluons will typically experience rescattering within the medium via further frequent transverse kicks, resulting in a diffusion process that features a Gaussian probability distribution. While harder gluons will also experience Gaussian broadening, they are already produced with a relatively large $k_\perp \gg \hat{q}L$ (since $t_f < L$), and a characteristic power-law tail $\sim \hat{q}_0 L / k_\perp^4$. The relative size of these contributions in the computation of the QF can be obtained by integrating the broadening distribution computed using the IOE framework [89] for k_\perp scales larger than wR , i.e., $B(\omega R; Q_{\text{broad}}^2) = (dI/d\omega)^{-1} \int_{(\omega R)^2}^{\infty} dk_\perp^2 dI / (d\omega dk_\perp^2)$, where Q_{broad}^2 represents the characteristic broadening scale of the different mechanisms.

In this way, the full radiation spectrum for emissions that end up beyond the jet cone, under this multiplicative Ansatz, reads [87]

$$\begin{aligned} \frac{dI_{>}}{d\omega} = & B(\omega R; Q_s^2/2) \frac{dI^{(0)}}{d\omega} \\ & + B(\omega R; \max[Q_s^2, 16\omega/(\pi^2 L)]) \frac{dI^{(1)}}{d\omega}. \quad (19) \end{aligned}$$

In the LO term, the broadening scale was set to $Q_{\text{broad}}^2 = Q_s^2/2$, where the saturation scale is $Q_s^2 = q_0 \ln(\hat{q}_0 L / \mu_*^2) L$.

This is the typical momentum acquired via multiple soft scatterings averaged over all possible production points between 0 and L [86,87]. The NLO term considers the broadening scale that results into a larger amount of broadening, either via the emission process itself, or by further multiple soft scatterings. The choice of scales in the NLO term correctly reproduce the GLV limit at high w and k_{\perp} , up to logarithmic factors. For more details on the steps here employed to incorporate broadening we refer the reader to the Appendix of our previous work [12] and to [90].

The bare quenching weight associated with medium-induced radiation is then

$$Q_{\text{rad}}^{(0)}(p_T) = \exp \left[- \int_{\omega_s}^{\infty} d\omega \frac{dI_{>}}{d\omega} (1 - e^{-\nu\omega}) - \int_T^{\omega_s} d\omega \frac{dI^{(0)}}{d\omega} (1 - e^{-\nu\omega(1 - (\frac{R}{R_{\text{rec}}})^2)}) \right], \quad (20)$$

where we have approximated $dI_{\text{NHO}}/d\omega \simeq dI^{(0)}/d\omega$ in the turbulent regime $T < \omega < \omega_s$. The bare quenching factor for elastic energy loss is

$$Q_{\text{el}}^{(0)}(p_T) = \exp \left[-\hat{e}L\nu \left(1 - \left(\frac{R}{R_{\text{rec}}} \right)^2 \right) \right]. \quad (21)$$

This expression is simply obtained as the Laplace transform of $\delta(\epsilon - \hat{e}L(1 - R^2/R_{\text{rec}}^2))$. For a weakly coupled plasma, as assumed in this work, the transport coefficient \hat{e} [91] is related to \hat{q} by using the fluctuation-dissipation relation $\hat{e} = \hat{q}/(4T)$ (where $\hat{e}_q = \frac{C_F}{N_c} \hat{e}_g$ and $\hat{e}_g = \hat{e}$ for gluons) [92].

C. Observables

The jet nuclear modification factor R_{AA} can be written as

$$R_{\text{AA}}(p_T, R) = \frac{\sum_{i=q,g} Q_i(p_T, R) \tilde{\sigma}_i^{pp}(p_T, R)}{\sum_{i=q,g} \sigma_i^{pp}(p_T, R)} \quad (22)$$

$$= Q_q(p_T, R) \tilde{f}_q(p_T, R) + Q_g(p_T, R) \tilde{f}_g(p_T, R). \quad (23)$$

As before, the tilded pp cross section implies that the free-proton PDFs have been replaced by nuclear PDFs that account for nuclear modifications and isospin effects. In particular, the modified quark and gluon fractions, defined as

$$\tilde{f}_i(p_T, R) = \frac{\tilde{\sigma}_i^{pp}(p_T, R)}{\sum_{i=q,g} \sigma_i^{pp}(p_T, R)}, \quad (24)$$

do not add up to one, i.e., $\sum_{i=q,g} \tilde{f}_i(p_T, R) \neq 1$, reflecting some of the nuclear modifications that take place even in the absence of final state energy loss.

The other main observable studied in this work are the jet azimuthal anisotropies, and more specifically the so-called elliptic flow coefficient, v_2 . It is defined as the second order Fourier coefficient in the expansion of the spectrum over the azimuthal angle ϕ [93,94], and can be computed in AA collisions as

$$v_2(p_T, R) \equiv \frac{\int_{-\pi}^{\pi} d\phi \cos(2\phi) (d\sigma^{\text{AA}}(p_T, R)/d\phi)}{\int_{-\pi}^{\pi} d\phi (d\sigma^{\text{AA}}(p_T, R)/d\phi)}, \quad (25)$$

where we used that the uncertainty in the event-plane angle resolution Ψ_R is negligible compared to other sources of uncertainties [95], and we have set it to $\Psi_R = 0$, as we do in our simulations.

D. Event-by-event medium properties

In this subsection we provide simple arguments that clarify the need to take into account fluctuations of the inmedium histories of jets in order to realistically compute jet observables in heavy-ion collisions. We also describe how we include them in our model.

Jets are produced at different locations in the transverse plane, with different orientations, and therefore explore different inmedium lengths and medium properties, such as temperature or flow, before escaping it. To clarify this point with an example, let us focus on the effect that fluctuations on the traversed length in the QGP, L , have on quenching. One can define the average value of a given quantity X by

$$\langle X \rangle = \frac{\int_0^{\infty} dL f(L) X(p_T, L)}{\int_0^{\infty} dL f(L)}, \quad (26)$$

where $f(L)$ represents the probability to traverse a length L . In the most trivial event-by-event (EbE) scenario, we assume a flat distribution between values L_{min} and L_{max} , yielding $f(L)_{\text{EbE}} = \Theta(L - L_{\text{min}})\Theta(L_{\text{max}} - L)$. We compare the results obtained with this distribution against those obtained by computing the given quantity X using only the average value (Ave) of L , where $f(L)_{\text{Ave}} = \frac{L_{\text{max}} + L_{\text{min}}}{2} \delta(L - (L_{\text{max}} + L_{\text{min}})/2)$.

Let us now choose X to be the amount of suppression of the spectrum. Again, for simplicity, let us just neglect the dependence on color factors and consider the bare quenching factor, ignoring thereby the phase-space resummation [96]. Then, $X = Q = e^{-aL}$, where $a = \sqrt{2\pi\hat{\alpha}^2\hat{q}\nu}$ with $\nu = n/p_T$. Here, we explicitly assume that \hat{q} is constant throughout the medium and independent of time. The ratio of the mean quenching factor between the ‘‘event-by-event’’ and ‘‘average’’ scenarios is

$$\frac{\langle Q \rangle_{\text{EbE}}}{\langle Q \rangle_{\text{Ave}}} = \frac{e^{-aL_{\text{min}}} - e^{-aL_{\text{max}}}}{a(L_{\text{max}} - L_{\text{min}})e^{-\frac{a}{2}(L_{\text{max}} + L_{\text{min}})}}. \quad (27)$$

This ratio tends to 1 as a goes to 0, for instance at high- p_T , but is larger than 1 at $a \gg 1$ [97]. This difference can induce a change in the slope of Q and can affect the agreement with high-precision experimental data in this regime. This very simple example illustrates that the average energy loss over many path lengths is not in general the energy loss of the average path length.

This discrepancy is not necessarily that striking for all quantities, though, as it obviously depends on the functional dependence of that quantity on L . For instance, in the case of v_2 , the EbE result coincides exactly with the Ave (in the small eccentricity approximation), but only if one uses the bare quenching weight,

$$\langle v_{2,i} \rangle_{\text{EbE}} = -\frac{2e}{L_{\text{max}} - L_{\text{min}}} \int_{L_{\text{min}}}^{L_{\text{max}}} \frac{Q'_i}{Q_i} L dL = \langle v_{2,i} \rangle_{\text{Ave}}, \quad (28)$$

simply because Q'_i/Q_i is independent of L [98]. This relation is no longer true when jet phase-space effects are included by making use of the resummed quenching factor because of its nontrivial L -dependence.

Therefore, in order to accommodate the variations in jet energy loss resulting from the diverse trajectories the jet may traverse within the expanding QGP, we must integrate our theoretical framework into a realistic heavy-ion environment. We follow the same procedure carried out in our previous publication [12], whose steps we summarize here for the reader's convenience.

We first determine the production point in the transverse plane (x, y) by using the overlap of the thickness functions of the colliding nuclei, $T_{AB}(x, y; b) = T_A(x - b/2, y) \times T_B(x + b/2, y)$, with b being the impact parameter of the nuclear collision. The thickness function is derived from the transverse density of nucleons within the Lorentz-contracted nuclei and is distributed following the Woods-Saxon density function [99]. Then, we randomly assign an orientation within the transverse plane and a random rapidity value within the range $-2 \leq y \leq 2$. While tracing the trajectory of the jet within the QGP, we calculate the integrated values of the necessary physical variables, which typically depend on the local temperature T and fluid velocity u until the jet exits the QGP phase, occurring at a pseudocritical temperature that we choose to be $T_c = 145$ MeV. These values of T and u are extracted from event-averaged hydrodynamic profiles [100] that describe the evolution of an expanding droplet of liquid QGP in PbPb collisions at $\sqrt{s} = 5.02$ A TeV, considering various centrality classes. Because parameters such as the fluid temperature are expressed in the local fluid rest frame, it is necessary to take into account the infinitesimal distance dx_F covered by the jet in this particular frame during each time increment dt in the laboratory frame.

Ignoring numerical constants, the essential physical variables required are obtained through integration along the trajectory $\gamma(t)$ of a jet as follows:

$$\begin{aligned} L &= \int_{\gamma(t)} dx_F, \\ T &= \frac{1}{L} \int_{\gamma(t)} dx_F T(x), \\ m_D^2 &\propto \frac{1}{L} \int_{\gamma(t)} dx_F T^2(x), \\ \hat{q}_0 &\propto \frac{1}{L} \int_{\gamma(t)} dx_F T^3(x) \left(\frac{p \cdot u(x)}{p^0} \right), \\ \hat{e} &\propto \frac{1}{L} \int_{\gamma(t)} dx_F T^2(x) \left(\frac{p \cdot u(x)}{p^0} \right), \\ \omega_c &\propto 2 \int_{\gamma(t)} dx_F L(t) T^3(x) \left(\frac{p \cdot u(x)}{p^0} \right), \\ \theta_c^{-\frac{1}{2}} &\propto 3 \int_{\gamma(t)} dx_F L^2(t) T^3(x) \left(\frac{p \cdot u(x)}{p^0} \right), \end{aligned} \quad (29)$$

where the length $L(t)$ is the inmedium path of the jet up to time t . The dilution factor $(p \cdot u)/p^0$, with p the four-momentum of the jet and u the fluid four velocity, arises in transport coefficients due to a flowing medium [101]. We want to make clear that even though we account for fluctuations in the jet production point and orientation, we do not encode any additional fluctuations associated with the variations of the relevant quantities along the path of the jet.

Central to the physics discussion of this work is the length dependence of the coherence angle θ_c . In Fig. 2 we show the distribution of the values of θ_c at RHIC energies, for gold-gold (AuAu) collisions, for a few centrality classes. To obtain these distributions we have followed the procedure outlined above. Given that the impact

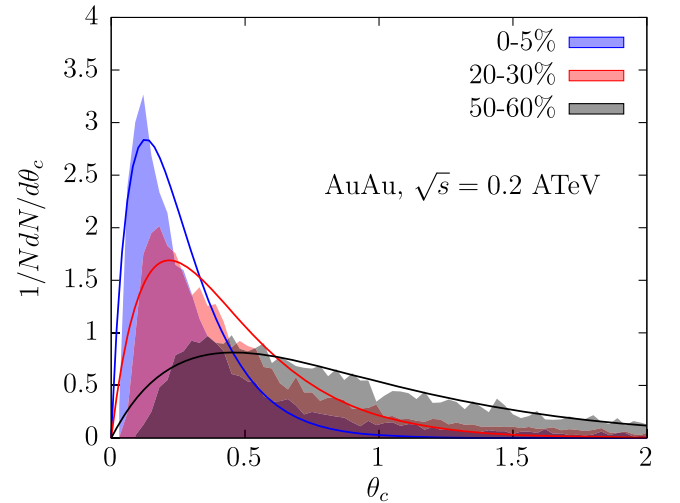


FIG. 2. Probability distribution of θ_c , as determined by the jet embedding procedure outlined in Sec. II D. We have just selected a few centrality classes for the collision kinematics at RHIC. The solid lines represent the fits obtained using Eq. (30).

parameter increases towards more peripheral collisions, the traversed lengths tend to decrease notably. This means that the typical value of θ_c increases with decreasing centrality. This is precisely what is seen in Fig. 2. Due to the reduced center-of-mass energies used at RHIC, the initial temperature of the QGP is lower, and therefore it reaches T_c earlier, reducing the typical traversed length L . This implies that, for a given centrality class, the typical values of θ_c at RHIC will be larger than those at LHC. One can easily see that the θ_c distributions from Fig. 2 are of the type

$$P(\theta_c) = \frac{\theta_c}{\theta_c^{*2}} e^{-\theta_c/\theta_c^*}, \quad (30)$$

where θ_c^* is the *mode*, i.e., the most frequent value, of the θ_c distribution for a given centrality class in a given collision system. This distribution has its origin in the distribution of path lengths explored by the different inmedium jet histories, as well as in the distribution of values of \hat{q} , for each centrality class and collision system. The value of θ_c^* , both at RHIC and LHC for a wide range of centrality classes, will be useful to interpret the results, as discussed in Sec. IV. The qualitative agreement of this set of fits can be seen by comparing the solid lines with the data histograms in Fig. 2, and the corresponding values of θ_c^* are tabulated in Table I. Further fluctuations associated with different initial state configurations will be studied in future work.

III. RESULTS

In this section we provide results for jet suppression, R_{AA} , and jet azimuthal anisotropy, v_2 , as a function of the jet cone angle R and jet transverse momentum p_T for a number of centrality classes that range from 0–5% to 60–70%, both at RHIC and LHC. They are computed using the framework described in Sec. II, which is the same as in our previous publication [12]. We compare our predictions to experimental data that was in many cases released after the publication of our previous work.

TABLE I. Fitted values for θ_c^* , using Eq. (30), both at RHIC (AuAu collisions at $\sqrt{s} = 0.2$ A TeV) and LHC (PbPb collisions at $\sqrt{s} = 5.02$ A TeV), for different centrality classes.

Centrality	θ_c^*	
	RHIC	LHC
0–5%	0.13	0.09
5–10%	0.15	0.10
10–20%	0.17	0.12
20–30%	0.22	0.15
30–40%	0.27	0.19
40–50%	0.35	0.24
50–60%	0.45	0.32
60–70%	0.58	0.41

We emphasize that our calculation relies on two free parameters, g_{med} and R_{rec} , namely the strength of the coupling to the medium, g_{med} , and the energy recovery parameter, R_{rec} , which were determined in our prior analysis of jet R_{AA} at LHC energies [12]. Accordingly, in all of our results, the theory error bands correspond to taking the minimum and maximum values among the four combinations for the parameters $g_{\text{med}} = \{2.2, 2.3\}$ and $R_{\text{rec}} = \{1, 5/6\} \frac{\pi}{2}$. We do not tune their values any further to obtain the results presented in the current work.

A. Jet suppression at RHIC and LHC

We start by showing in Fig. 3 our results for jet R_{AA} for three different radii, $R = 0.2$, $R = 0.3$, and $R = 0.4$, in the different panels, confronted against experimental data measured by STAR [102]. Our results are computed for jets with $|\eta| < 0.5$, while those from STAR have $|\eta| < 1$. Error bars on data denote statistical uncertainties, while bands refer to systematic uncertainties. The dashed line indicates the region below which the experimental results are biased by the requirement of the presence of a high- p_T track within the jet, which is something that the STAR collaboration needed to apply in order to reduce the contamination from fake background jets, and that cannot be replicated by our present semianalytical framework. Another difference is the fact that STAR only used charged tracks to reconstruct the jets, while our calculation is done for full jets. A rather crude way to account for this difference, which we did not attempt to do, would be to shift the jet p_T by a factor 2/3, although this would have little effect on the unbiased region due to the apparent relative flatness of R_{AA} with jet p_T . We observe good agreement between our predictions and experimental data in the unbiased region, above $p_T \gtrsim 15$ GeV, correctly reproducing the small dependence of jet suppression on jet-size R .

In order to have a closer look into this mild R -dependence of jet suppression, we present ratios of different $R_{AA}(R)$ over the result for $R = 0.1$ for RHIC energy in Fig. 4. At low p_T , increasing the jet radius translates into less suppression (larger R_{AA}), since the emissions off the very few resolved color charges can still be captured within the larger cone. As one increases the jet p_T , the phase space within the jet increases, more so for a larger R such as $R = 0.6$, and the number of resolved charges increases. This leads to a larger suppression [still $\mathcal{O}(10\%)$ effects at most] of larger R at higher p_T , overcoming the capturing effect that explains the behavior at low p_T . The R dependence of jet suppression is thus sensitive to very relevant aspects of the jet-medium interaction, such as the distribution of the radiated energy and the size of the resolved phase space. Future jet measurements at the sPHENIX detector [103] are expected to be precise enough to corroborate this picture.

After the publication of our previous work [12], the ALICE collaboration presented their results for the R

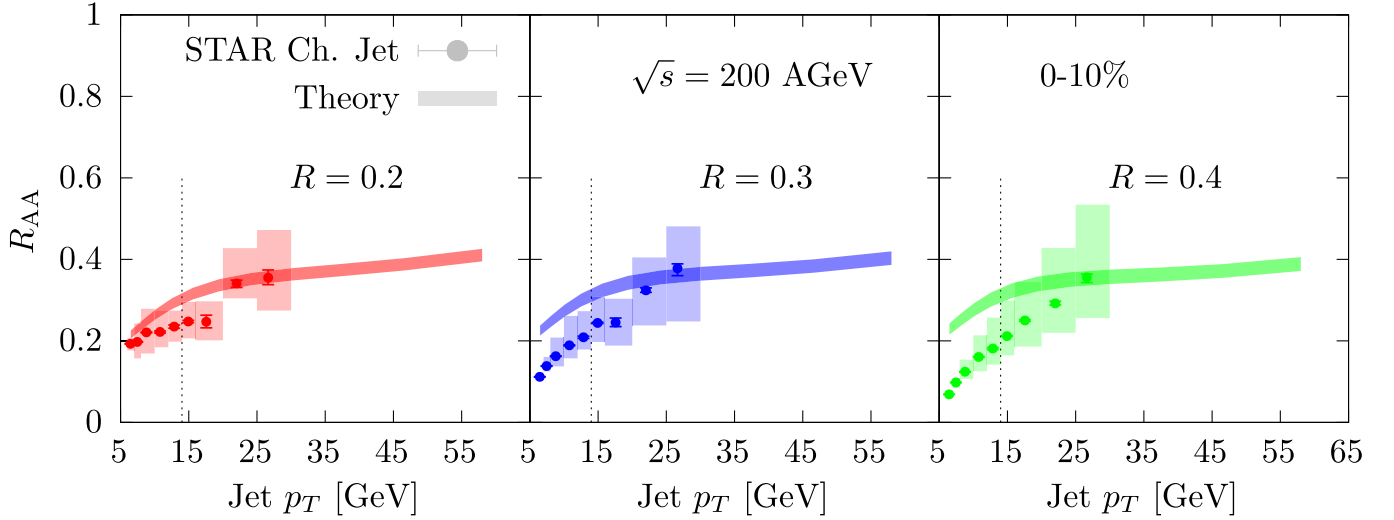


FIG. 3. Comparison of jet R_{AA} against STAR data [102] measured at RHIC at $\sqrt{s} = 0.2$ A TeV, where different panels show different R . The dashed vertical line indicates the jet p_T below which the existence of an experimental bias prevents a direct comparison with our semianalytical results.

dependence of jet suppression at fairly low jet p_T , ranging from $R = 0.2$ to $R = 0.6$ [105] (the R dependence of this low jet p_T range was first measured years ago by ATLAS [104], although in terms of a double central-to-peripheral, R_{CP} ratio that we do not attempt to reproduce in the present work). In order to account for the differences in acceptance with respect to ATLAS' (whose results we compared against in our previous publication [12]), i.e., $|y| < 2.8$, we redid our calculations by modifying the rapidity ranges of the initial spectrum at $R_0 = 1$ from which we compute the vacuum and medium jet evolution, such that $|\eta| < 0.9 - R$. Any expected causes for differences in R_{AA} due to the different rapidity ranges, such as the change of the spectral index n (larger at larger η) and the quark-initiated jet fraction (larger at larger η) turned out to be very small. We show our predictions, confronted against ALICE data, in Fig. 5, obtaining an excellent agreement.

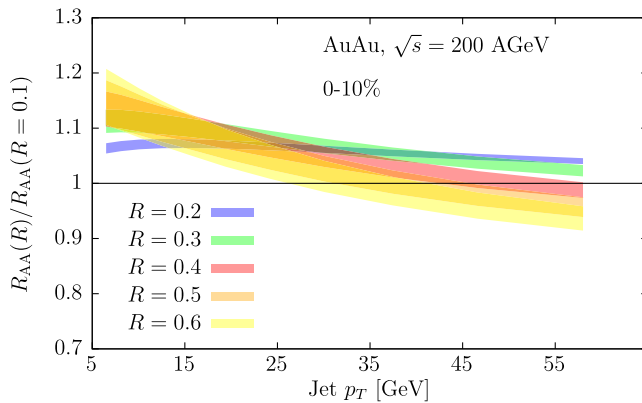


FIG. 4. Ratio of R_{AA} for different R over that for $R = 0.1$, at $\sqrt{s} = 0.2$ A TeV for AuAu collisions in the 0–10% centrality class.

In Fig. 6 we show jet R_{AA} , also at LHC, but for a much higher jet p_T . Our results are compared to recent ATLAS data for $R = 0.2$ [106], ATLAS data for $R = 0.4$ [108], and CMS data [107] for $R = 0.2, 0.4, 0.6$. The first point, at $p_T \approx 100$ GeV of ATLAS results for $R = 0.4$ jets [108], was used to fit our value of g_{med} in [12]. The small larger suppression experienced by $R = 0.4$ jets as compared to $R = 0.2$ jets, which tends to disappear with decreasing jet p_T , as measured by ATLAS, is correctly captured by our prediction. As predicted in our previous work [12], jet $R_{AA}(R = 0.2)$ should be somewhat less suppressed than those with larger R , which hold even more similar values of suppression among them. To a large extent, this is so due to the fact that the typical critical angle for this centrality is around $\theta_c \sim 0.12$, representing an extensive part of the phase space of an $R = 0.2$ jet. However, this mild trend was not observed by CMS, where $R = 0.2$ and $R = 0.4$ lie on top of each other for the limited available range in jet p_T . Moreover, CMS has measured that also $R = 0.6$ jets are as little suppressed as $R = 0.2$ jets, together with $R = 0.4$ jets, within current uncertainties. The manifest, although mild tension between the ATLAS and CMS results for $R = 0.4$ jets motivates extending the p_T range and precision of measurements. Further improvements from the theory side involve increasing the precision of the resolved phase-space calculation, encapsulated in the value of θ_c in our framework, as it can induce $\mathcal{O}(20\%)$ changes in jet suppression [12], specially at very high p_T .

In summary, the excellent agreement as a function of jet $p_T \gtrsim 100$ GeV up to $p_T \lesssim 1$ TeV, for $R = 0.4$ jets belonging to a wide range of centrality classes shown in our previous work [12], combined with the no less remarkable results here shown, which extend the p_T reach down to $p_T \gtrsim 20$ GeV, both for smaller ($R = 0.2$) and larger

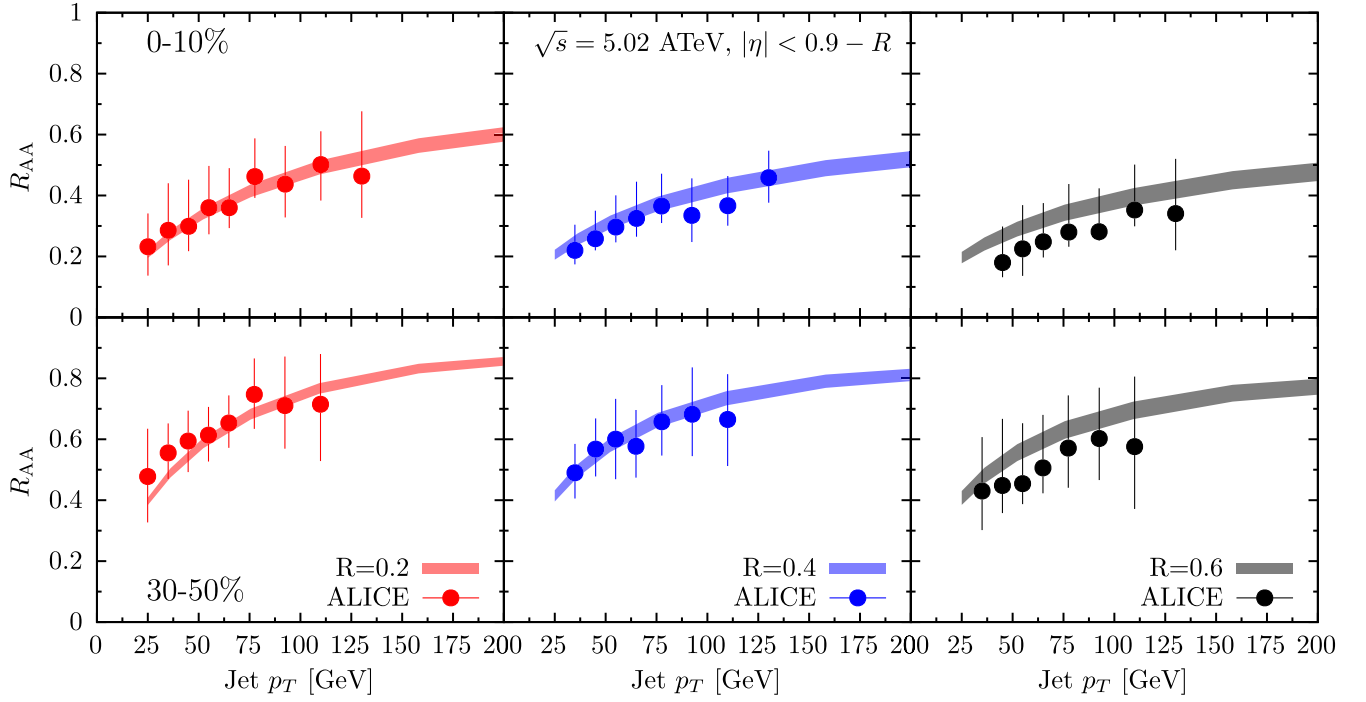


FIG. 5. Comparison of theoretical predictions of R_{AA} against ALICE data [105] for PbPb collisions at $\sqrt{s} = 5.02$ TeV, for the 0–10% (30–50%) centrality class in the upper (lower) row. Different columns display different values of jet cone R .

($R = 0.6$) jet radii, and for several centrality classes as well, present overall an encouraging picture of our first-principles understanding of jet suppression in deconfined QCD matter. Very good agreement is achieved across centralities, spanning across almost two orders of magnitude in jet p_T , and several jet radii, up to $R = 0.6$, which we deemed to be the radius above which nonperturbative effects due to correlated background (medium response) start to be important [12].

A complete set of predictions for jet R_{AA} at RHIC and LHC, for different centralities and jet radii up to $R = 0.6$, can be found in Appendix.

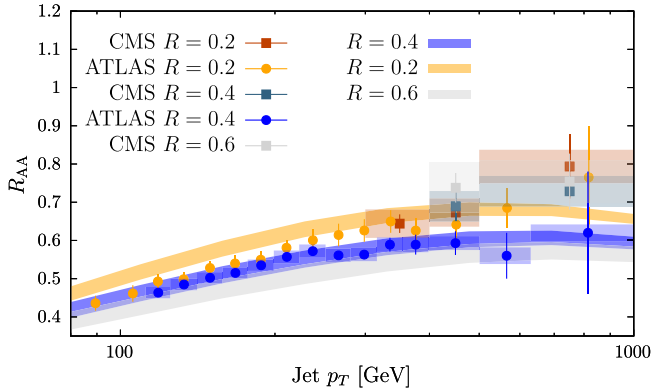


FIG. 6. Our predictions for R_{AA} with $R = 0.2$ confronted against ATLAS data [106] and CMS data [107], for PbPb collisions at $\sqrt{s} = 5.02$ TeV. In our previous publication [12], we used the first point of ATLAS data for $R = 0.4$ [108] for 0–10% centrality class to fit our value of g_{med} .

B. Jet azimuthal anisotropy at RHIC and LHC

We now turn to the discussion of the results for jet v_2 , defined in Eq. (25). In our main set of results we do not consider fluctuations in the initial collision geometry, as our hydrodynamic backgrounds are event averaged within a given centrality class, and so we have set the event plane to $\Psi_R = 0$. Just as it is the case for the flow coefficients of the soft particles from the bulk of the system, these fluctuations are expected to affect the values of higher-order flow harmonics, such as v_3 , for high- p_T objects as well [19].

In order to compute the jet v_2 as in Eq. (25), one simply needs to label each inmedium jet history with the azimuthal angle along which it propagated, ϕ . One then readily obtains the results for $R = 0.2$ jets at LHC, shown in Fig. 7, which are compared against ATLAS data [109]. While ATLAS jet R_{AA} results typically feature an extended rapidity coverage, $|y| \leq 2.8$, these jet v_2 results are for jets within $|y| < 1.2$. We have recomputed our double differential jet spectra, both in p_T and ϕ , for this narrower rapidity range. In contrast to the case of jet R_{AA} , where we found negligible effects from narrowing the rapidity range when comparing to ALICE data, jet v_2 is more sensitive to the decrease of the quark-fraction that is induced by the reduction of the rapidity range about midrapidity, especially so towards lower jet p_T [52]. The reason is that, in the soft limit, jet $v_2 \sim C_R$, as will be illustrated in Sec. IV. Our predictions (solid bands) yield in general very good agreement with experimental data, both in total magnitude and in the jet p_T -dependence.

The comparisons shown in Fig. 7 point to the presence of tensions in the 0–5% centrality bin, where we overpredict the

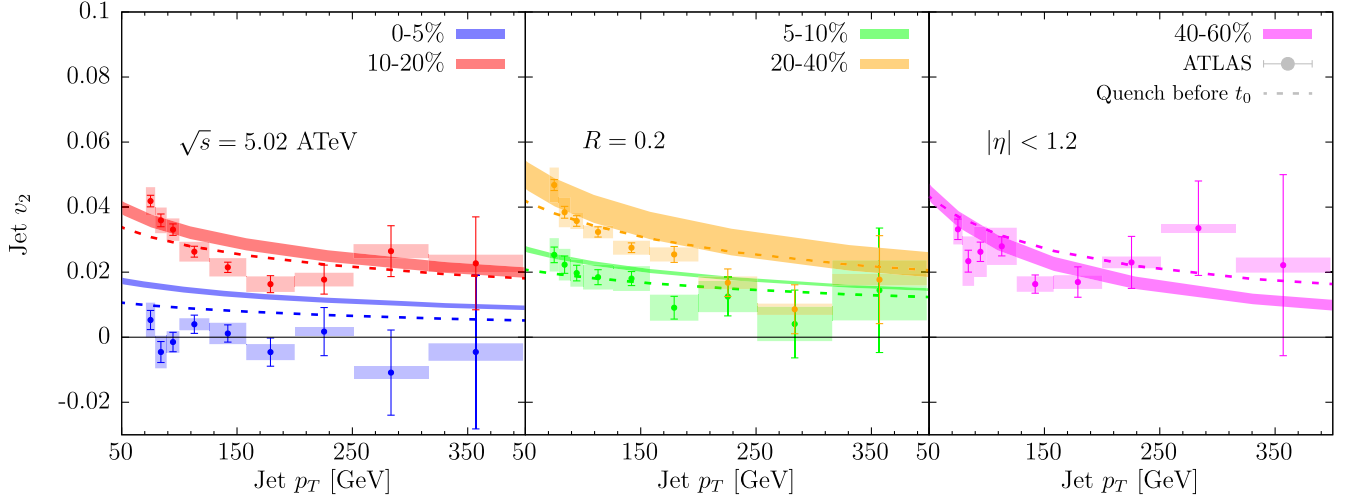


FIG. 7. Results for jet $v_2(p_T)$ for $R = 0.2$ jets at LHC for PbPb collisions at $\sqrt{s} = 5.02$ A TeV, confronted against ATLAS data [109]. See text for details.

size of jet v_2 at lower jet p_T . One possible reason could be our oversimplified treatment of quenching physics before hydrodynamization time, $t_0 = 0.6$ fm/c in our simulations, where we simply ignore energy loss effects. It has been shown that neglecting quenching in the earliest stages overestimates the size of v_2 [110,111]. Including the effects of quenching in the initial stages, for which theoretical calculations are now becoming available [112,113], will be studied in future work. In the present work, we estimate the effect by allowing the possibility of elastic and radiative energy loss before t_0 , as if our present formulas were valid out of equilibrium as well. In practice, we set $T(t < t_0, \vec{x}) = T(t_0, \vec{x})$ and $v_x(t < t_0) = v_y(t < t_0) = 0$, while still keeping $v_z = z/t$. Quenching during more time naturally reduces the value of jet R_{AA} , so in order to test the effect on v_2 for the same amount of jet suppression, we need to reduce the value of the coupling with the medium. While the bands represent our default results with quenching $t \geq t_0$, obtained with $g_{\text{med}} \in \{2.2, 2.3\}$, the dotted lines are the results with a reduced $g_{\text{med}} = 2$, so that jet R_{AA} for $R = 0.4$ jets at 0–10% centrality lies within the band. We observe that, with the exception of the most peripheral class, the dotted lines tend to be below the bands, yielding a smaller jet v_2 , in consistency with what has been found in the literature [22,110,111]. We note that the largest relative effect is, for the most central class, 0–5%, where the reduction of jet v_2 is almost by a factor of two. While the reduction of jet v_2 does not significantly modify the level of agreement with experimental data for more peripheral classes, in the case of the 0–5% class the sizeable relative reduction brings our theoretical predictions much closer to ATLAS data.

We further explore the R dependence of jet v_2 in Appendix, both at RHIC and LHC. In general, jet v_2 increases as a function of R , as expected, since a larger R tends to imply a larger resolved phase space. This is the main reason why R_{AA} tends to decrease with increasing

$R \lesssim 0.6$ at high jet p_T . This trend tends to be reversed towards lower jet p_T , as illustrated for example in Fig. 4, where the effect of increasing the amount of recaptured energy by opening the cone overcomes the associated enlargement of the resolved phase space. However, something distinct happens in the case of jet v_2 , where always $v_2(R, p_T) > v_2(R', p_T)$ if $R > R'$, regardless of the relative ordering in R_{AA} . This is best appreciated in Fig. 13, for RHIC results, where the typically lower jet p_T compared to LHC kinematics yield $R_{AA}(R, p_T) > R_{AA}(R', p_T)$, over a wide range in the accessible jet p_T , while $v_2(R, p_T) > v_2(R', p_T)$, for $R > R'$.

In particular, we observe an interesting grouping among different R , evolving with centrality, which is again best perceived for RHIC kinematics. For central collisions, say 5–10% centrality class, $v_2(R = 0.1, p_T) \approx v_2(R = 0.2, p_T) \lesssim v_2(R = 0.3, p_T) < v_2(R = 0.4, p_T) \lesssim v_2(R = 0.6, p_T)$, while by the 50–60% centrality class $v_2(R = 0.1, p_T) \approx v_2(R = 0.2, p_T) \approx v_2(R = 0.3, p_T) \approx v_2(R = 0.4, p_T) < v_2(R = 0.6, p_T)$. This type of grouping is absent for R_{AA} , and reflects the connection between jet v_2 and the path-length dependence of the critical angle θ_c , whose most frequent value θ_c^* strongly evolves with centrality, as shown in Fig. 2. These aspects are explained in Sec. IV based on simple arguments. At this stage, we merely illustrate the behavior of this grouping by computing the difference in jet v_2 between jets with a given $R > 0.1$ and those with $R = 0.1$, for a fixed jet $p_T = 20$ GeV, as a function of centrality, as shown in Fig. 8. We observe the sequential collapse of the value of jet v_2 for jets with a given R towards the value of jet v_2 for the smallest $R = 0.1$ at different centrality classes. According to the projected uncertainties associated to jet v_2 measurements with the sPHENIX detector [114], we believe that these subtle features can be confronted against upcoming RHIC data [115].

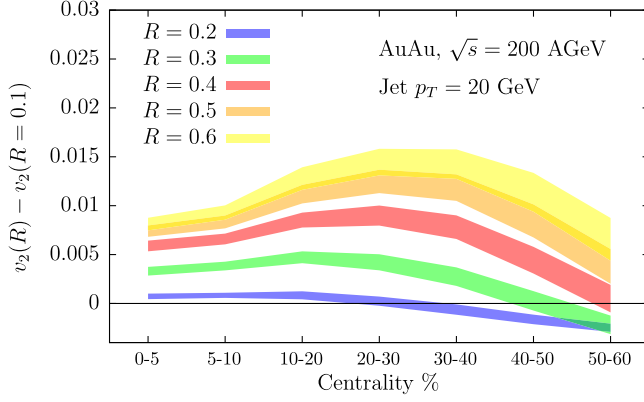


FIG. 8. Difference between the jet $v_2(R)$ for $R > 0.1$ with respect to $v_2(R = 0.1)$, for jet $p_T = 20$ GeV and as a function of centrality at RHIC for AuAu collisions at $\sqrt{s} = 0.2$ A TeV.

IV. DISCUSSION

To gain insight into the results presented in the previous section, particularly to better comprehend the seemingly distinct behavior of jet v_2 concerning centrality compared to that of jet R_{AA} , we will try to summarize the main features with an analytical model that captures the qualitative features of our fully-fledged event-by-event calculations above.

Our main goal is to reveal the importance of jet coherence, an aspect that has been largely neglected in existing literature. The geometry-sensitive nature of v_2 can be utilized to probe the length dependence of both quenching *and* resolution effects. To illustrate our point, consider measuring a quenched jet “in plane” versus “out of plane.” The jet traveling “in plane” will not only experience a shorter path length, leading to a smaller bare quenching, but also be less resolved leading also to a larger θ_c and a smaller phase space for additional substructure quenching. In contrast, the jet traveling “out of plane” is both more quenched *and* more resolved. This additional effect emphasizes the path-length differences embodied in an observable such as v_2 and will lead to: (i) an enhanced v_2 of jets compared to hadrons, and (ii) a characteristic increase of v_2 with R for a given centrality, as seen, for example, in Fig. 8.

A. R dependence of R_{AA}

Taking the strong quenching limit, i.e., $Q_i^{(0)} \ll 1$, in Eq. (13), see Ref. [12] for details, and using the fact that the phase space is dominated by the double log which allows us to approximate $z \sim 1$ (we further assume a fixed coupling constant), only the virtual term contributes and the solution can be written directly as

$$Q_i(p_T, R) \simeq Q_i^{(0)}(p_T) e^{-\Omega_{\text{res}}(p_T, R)}. \quad (31)$$

Here, we effectively account the quenching of the leading parton (total charge) via $Q_i^{(0)}$. For the moment, to simplify

our discussion in this section, we neglect any weak R dependence of this factor. This phase space is sensitive to resolved subjects in the QGP and is given by

$$\begin{aligned} \Omega_{\text{res}} &= \int_0^R \frac{d\theta}{\theta} \int_0^1 dz \frac{\alpha_s}{\pi} P(z) \Theta(t_f < t_d < L), \\ &\approx 2\bar{\alpha} \int_{\max(\theta_c, \theta_{\min})}^R \frac{d\theta}{\theta} \int_{(\frac{2}{3}\hat{q}/\theta^4)^{1/3}}^{p_T} \frac{d\omega}{\omega}, \end{aligned} \quad (32)$$

where we used that the splitting function $P(z) \approx 2C_R/z$ in the soft limit of gluon emission ($\omega = zp_T$). In this limit, we get that the formation time is $t_f = 2/(\omega\theta^2)$ and the decoherence time is $t_d = (\hat{q}\theta^2/12)^{1/3}$. The limit on the angular integral is then $\theta_c = (\hat{q}L^3/12)^{-1/2}$ or $\theta_{\min} \equiv (\frac{2}{3}\hat{q}/p_T^3)^{1/4}$ [9]. In the high- p_T limit, precisely for $p_T > \omega_c/3$, and for $R > \theta_c$, we get

$$\Omega_{\text{res}} = 2\bar{\alpha} \ln \frac{R}{\theta_c} \left(\ln \frac{3p_T}{\omega_c} + \frac{2}{3} \ln \frac{R}{\theta_c} \right), \quad (33)$$

where the characteristic energy scale is $\omega_c = \hat{q}L^2/2$ and the characteristic angular scale is $\theta_c = (\hat{q}L^3/12)^{-1/2}$. Clearly, if $R \leq \theta_c$ the phase space vanishes $\Omega_{\text{res}} = 0$ and the jet is quenched *coherently*, only according to its total color charge, i.e.,

$$Q_i(p_T, R)|_{R \leq \theta_c} = Q^{(0)}(p_T), \quad (34)$$

where $Q^{(0)} = Q_{\text{rad}}^{(0)} Q_{\text{el}}^{(0)}$, corresponding to radiative and elastic energy losses, respectively. We will not consider elastic energy loss in this section, but note that most of the conclusions regarding the scaling properties of R_{AA} and v_2 still hold under the general condition that $Q^{(0)} = e^{-f(p_T, R)L}$, where $f(p_T, R)$ is an arbitrary function of the jet transverse momentum and cone angle.

For example, in what we refer to as the “soft” approximation, i.e., when we assume that the medium-induced spectrum scales as $dI/d\omega \propto \omega^{-3/2}$, the bare quenching weight becomes $Q_i^{(0)}(p_T) = \exp(-2\sqrt{\pi\bar{\alpha}^2\omega_c n/p_T})$ [66], and only $\bar{\alpha} = \alpha_s C_i/\pi$ distinguishes between the quark and gluon energy loss. This suppression factor describes how a single parton suffers energy loss through soft gluon emissions *at any angle*. It becomes directly relatable to the quenching of single-inclusive hadron spectra if we replace the parton p_T by the transverse momentum of the leading hadron, or $p_T \rightarrow p_T' \equiv p_T/\langle z \rangle$; see also [117]. For example, for massless hadrons $\langle z \rangle \approx 0.3-0.5$ over a wide jet p_T range [118].

Assuming that the spectrum is dominated by a single parton species and neglecting nuclear PDF (nPDF) effects, we find that

$$R_{AA} \simeq Q(p_T, R). \quad (35)$$

Including the full flavor dependence will influence jet R_{AA} (and hence also jet v_2) on the quantitative level only. Furthermore, for this discussion we omit effects coming from the recovery of emitted energy at large angles.

The regime of strong quenching, given by $Q_i^{(0)} \ll 1$, is relevant for the mid- to high- p_T range at RHIC and LHC. This is applicable in the mid- to high- p_T regime where $p_T \approx 40$ GeV. Hence, assuming the linearized form of the full quenching factor, cf. Eq. (31), the ratio is

$$\frac{Q(p_T, R)}{Q(p_T, 0.1)} \approx \exp[\Omega_{\text{res}}(p_T, 0.1) - \Omega_{\text{res}}(p_T, R)] \approx 1 - 2\bar{\alpha}\mathcal{L}_c \ln\left(\frac{\max(R, \theta_c)}{\max(0.1, \theta_c)}\right), \quad (36)$$

where $\mathcal{L}_c \equiv \ln(3p_T/\omega_c)$, up to terms that are not enhanced by a logarithm of p_T . We conclude that the ratio of jet suppression factors is mainly logarithmically dependent on the coherence angle θ_c .

B. R dependence of v_2

Simplifying the precise details of the geometrical aspects of the suppression one can estimate the jet anisotropy as [119]

$$v_2 \simeq \frac{1 R_{AA}(L_{\text{in}}) - R_{AA}(L_{\text{out}})}{2 R_{AA}(L_{\text{in}}) + R_{AA}(L_{\text{out}})}, \quad (37)$$

see also [84], where L_{in} is the shorter path length experienced by the jets that are emitted in plane, while L_{out} is the path length of jets emitted perpendicular to it, i.e., out of plane. For slightly eccentric systems such as those corresponding to central and semicentral collisions we can assume that the typical differences in path lengths between in plane and out of plane are small, i.e., $L_{\text{in}} = L$ and $L_{\text{out}} = L + \Delta L$, where $\Delta L/L \ll 1$. Assuming that the nuclear overlap resembles a perfect ellipse, the path length difference ΔL is directly related to the eccentricity of the ellipse e via $\Delta L/L = 2e$. It follows that Eq. (37) takes the simple form,

$$v_2 \simeq -\frac{\Delta L}{4} \frac{d \ln R_{AA}}{dL}, \quad (38)$$

upon expansion to first order in ΔL .

To gain analytical insight into the behavior of v_2 as a function of jet p_T and the cone angle R , we will at this point introduce further simplifications. Again, we assume that the spectrum is dominated by a single parton species and we will neglect nPDF effects. Furthermore, we will assume that the full quenching factor can be approximated by its linearized version with “soft” bare quenching factors, as done in the previous subsection, Sec. IV A. These sets of simplifications allow us to better isolate the effects of the

relation between the resolved phase space of a jet and energy loss.

In line with previous approximations, we expand the quenching weight up to the first order in ΔL , i.e., $Q(L + \Delta L) \approx Q(L)(1 + \Delta L \delta Q)$ to obtain

$$v_2 \approx -\frac{\Delta L}{4} \delta Q. \quad (39)$$

We define the perturbed resummed quenching factor as $\delta Q(p_T, R) = \partial \ln Q(p_T, R)/\partial L$, where the path-length dependence is implicit through the dependence on the medium scales ω_c and θ_c . The remaining part of this section will be dedicated to studying the path-length dependence of the quenching factor through δQ based on Eq. (31). As in the previous subsection, we will now focus on the regime of strong quenching, where $Q^{(0)}(p_T) \ll 1$. This immediately implies $\delta Q \ll 1$.

Let us start considering the path-length dependence of the bare quenching factor $Q^{(0)}(p_T)$. We find that

$$\delta Q_0 = \frac{\partial \ln Q^{(0)}(p_T)}{\partial L} \approx -C_R \sqrt{\frac{2\alpha_s^2 \hat{q} n}{\pi p_T}}, \quad (40)$$

where the last step is valid for the “soft” approximation. Due to its specific dependence on the variable ω_c/p_T , we can also trade the derivative on the path length for a derivative on p_T , i.e., $\delta Q_0 = -\frac{2p_T}{L} \frac{\partial \ln Q^{(0)}(p_T)}{\partial p_T}$ [24].

Another contribution to the perturbation of the quenching factor resulting from the Sudakov suppression factor needs to be accounted for, i.e., $\delta Q = \delta Q_0 + \delta Q_1$, where $\delta Q_1 = \delta \Omega_{\text{res}}(p_T, R)$; see Eq. (33). This additional term encapsulates the path length variation of additional jet splittings that occur in the medium. It turns out that the correction to the phase space is

$$\delta \Omega_{\text{res}}(p_T, R) = \frac{3\bar{\alpha}}{L} \mathcal{L}_c, \quad (41)$$

again with $\mathcal{L}_c \equiv \ln(3p_T/\omega_c)$, which is independent of the jet cone angle up to the condition $R > \theta_c$. Although derived within simplified assumptions, the terms that we have identified allow us to extract several useful qualitative characteristics of jet v_2 that we can compare to the results from our full numerical simulations. Inspecting the two terms in Eqs. (40) and (41), it becomes immediately clear that $v_{2,i} \propto \delta Q_i \propto C_R$ for $i = q, g$. This is simply a consequence of the Casimir scaling of the quenching factors, where both $\ln Q_i^{(0)} \propto C_R$ and $\ln Q_i \propto C_R$ [120].

The flow coefficient is then found to be

$$\frac{v_2}{e} \simeq \bar{\alpha} \left(\sqrt{\frac{\pi \omega_c n}{p_T}} + \frac{3}{2} \mathcal{L}_c \Theta(R - \theta_c) \right), \quad (42)$$

where we highlighted the leading dependence of the second term on the cone angle. In what follows, we will neglect the additional weak logarithmic dependence on L in \mathcal{L}_c . This illustrates a distinctive jump in jet v_2 at fixed p_T and as a function of R arising from the coherence physics encoded in the full quenching factor. Other subleading contributions to the R dependence are not discussed further here (see the comment in Sec. IV A).

Note also that in the ‘‘coherent regime’’, i.e., for $R \leq \theta_c$, we uncover a simple relation between jet quenching and the azimuthal asymmetry,

$$\left. \frac{v_2}{e} \right|_{R \leq \theta_c} \approx -\frac{1}{2} \ln R_{AA}|_{R \leq \theta_c}, \quad (43)$$

expected to hold for modest path-length differences equivalent to small eccentricity. Assuming a relatively central collision, 20–30%, for which $e \sim 0.2$ [121], with a charged hadron $R_{AA} \sim 0.6$ around $p_T \approx 40$ GeV [122] the above pocket formula predicts the estimated value $v_2 \sim 0.05$, which is in the ballpark of the experimental data [123]. Note that this correspondence is exact in our model study, where we neglected expansion effects and employed the ‘‘soft’’ approximation. It should nevertheless hold under general assumptions as long as the ‘‘bare’’ quenching factor scales with L , i.e., $\ln Q^{(0)} \propto L$. Finally, it is worth pointing out that the azimuthal asymmetry is a small $v_2/e \sim \mathcal{O}(\alpha_s)$ effect, in contrast to the overall jet suppression factor $R_{AA} \sim \mathcal{O}(1)$.

The sensitivity to the critical angle θ_c is a direct consequence of coherence physics and therefore, a concrete prediction of our calculation. For example, instead of the coherent phase space, if we assume that all emissions created in the medium, i.e., with $t_f < L$, fully contribute to jet quenching, we arrive at a collimator function with the phase-space $\Omega_{\text{incoh}}(p_T, R) = \frac{\tilde{q}}{2} \ln^2(\frac{1}{2} p_T R^2 L)$ [9], again for large enough p_T so that the argument of the log is positive.

In this case, there is a similar logarithmic enhancement of jet v_2 since $\delta\Omega_{\text{incoh}} = \frac{\tilde{q}}{L} \ln(\frac{1}{2} p_T R^2 L)$. However, this contribution is expected to appear for all R , leading to a much weaker relative R dependence.

Our above discussion also serves to draw a direct link between single-inclusive hadron and fully reconstructed jet v_2 at high- p_T . The hadron v_2 originates from a parton v_2 at a p_T which is larger by a factor $1/\langle z \rangle$ and, since the quenching of the total charge contribution diminishes with p_T , it will always be smaller. On the other hand, jet v_2 is enhanced compared to the hadron v_2 due to the same reasons that R_{AA} is lower.

C. Identifying an angular scale from jet suppression

The different qualitative behavior depending on the relative size of R and θ_c found in previous sections suggests a representation of observables based on the ratio θ_c/R . This approximate scaling behavior highlights the interplay between quenching and resolution effects affecting jets in the quark-gluon plasma.

In Sec. II D we have not only learned that the typical value of θ_c evolves strongly with centrality, but also that within a given centrality the corresponding θ_c distribution is relatively wide. We have extracted the most frequent value of θ_c for each centrality and collision system, $\theta_c^*(\sqrt{s}, \text{cent.})$, shown in Table I. We will now make use of these set of values to present results in terms of the ratio θ_c^*/R .

We plot the jet R_{AA} and jet v_2 results obtained with the full semianalytical calculations employed in Sec. III in Fig. 9. In order to factor out absolute effects related to the total amount of suppression, we have presented the results for jets with $R \geq 0.2$ divided by the results for jets with $R = 0.1$. Small- R jets, such as $R = 0.1$, are meant to be proxies for relatively unresolved color structures. The x axis of these plots is a convenient measure of centrality, with 0–5% being the most central and 60–70% the most

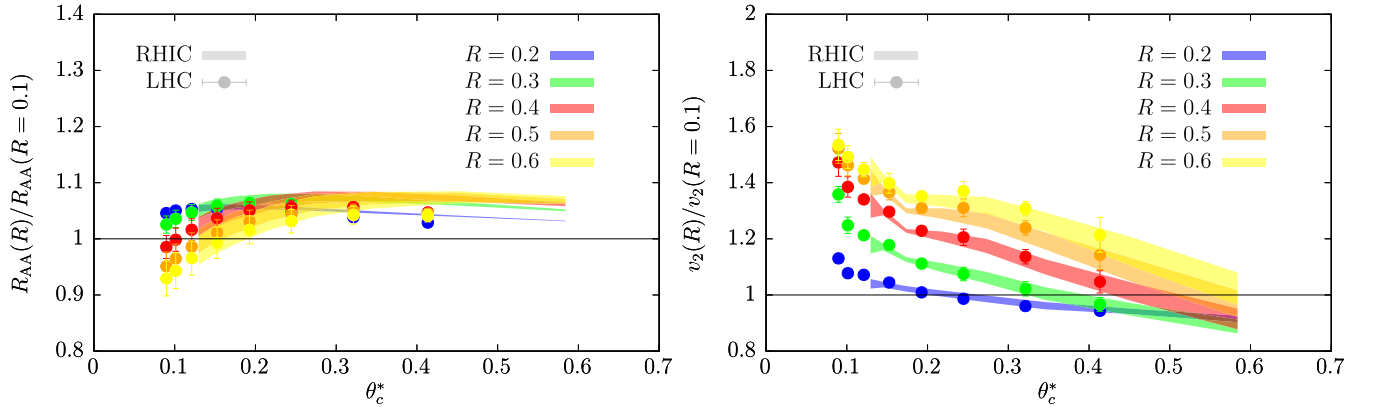


FIG. 9. Left: Double ratio of R_{AA} between jets with different radii R and jets with $R = 0.1$. Right: Double ratio of jet v_2 between jets with different radii R and jets with $R = 0.1$. Results are plotted against the variable θ_c^* , which is the most frequent value of the θ_c distribution for a given collision system and centrality. Jet momentum has been set to $p_T \approx 38$ GeV.

peripheral (as illustrated in Table I). Given that θ_c decreases with path length, we have chosen θ_c^* as such a measure (a smaller θ_c corresponds in general to more central collisions). In this plot, the data points represent calculations for LHC kinematics with model parameters identical to the analysis of the experimental data above. The shaded lines are equivalent calculations for RHIC kinematics.

While the R_{AA} ratios for different cone sizes R are fairly similar to each other, with deviations $\mathcal{O}(10\%)$, we see a strong ordering of the jet v_2 ratios with R , as expected from the qualitative discussion above. The jet p_T has been set to be $p_T \approx 38$ GeV, since this is a regime accessible in both collision systems. Next, we note that the calculations for RHIC and LHC kinematics fall on top of each other. Focusing on the v_2 plot in Fig. 9 (right), we can see that the 20–30% centrality class for LHC kinematics (corresponding to the third data point in the figure) overlaps with the prediction for RHIC kinematics in the 0–5% centrality class (the leftmost part of the band). This reveals how the underlying distributions of θ_c are similar (see Table I for the most frequent value θ_c^*) and that this is driving the behavior of v_2 .

In Fig. 10 we present results for jet R_{AA} (left panel) and jet v_2 (right panel) as a function of the scaling variable θ_c^*/R . We can observe a striking collapse of all data points onto clearly defined paths along θ_c^*/R , both for the jet R_{AA} and jet v_2 ratios shown in each panel. As above, the jet p_T has been set to be $p_T \approx 38$ GeV, since this is a regime accessible in both collision systems. Computations for RHIC kinematics reach larger values of θ_c^*/R than LHC kinematics, and LHC results reach smaller values of θ_c^*/R than RHIC results.

Starting with the right panel of Fig. 10, where we plot the scaling of jet v_2 ratios, one of the most visible features is the congruence at $\theta_c^*/R \approx 1$. This is a clear manifestation of the collapse already hinted at in Fig. 8, and is expected based on the estimates from Eq. (42) for unresolved jets.

Most importantly, this reveals a radical change of behavior of jet quenching at $\theta_c^*/R < 1$, where jet substructure fluctuations are contributing to the overall quenching, and $\theta_c^*/R > 1$, where the jets are quenched coherently. Succinctly, this plot illustrates the main features of our predictions for length-differential jet quenching at different center-of-mass energies, centralities, and cone sizes R .

The observed general trends can be easily understood using the simple analytical results of the previous sections. As already pointed out in Sec. II D, attempting to compare this analytical toy model to our complete framework requires at least the inclusion of some measure of the fluctuations of the dominant variables. Therefore, for a *qualitative* comparison, we allow the fluctuation of θ_c while keeping all other parameters fixed (e.g., ω_c). In this way, the steplike behavior of Eq. (42) now becomes similar to an exponential decay, i.e.,

$$\Theta(R - \theta_c) \rightarrow \int_0^R d\theta_c P(\theta_c) = 1 - \frac{1+t}{te^t}, \quad (44)$$

where we used the parametrization in Eq. (30) and $t = \theta_c^*/R$. A similar procedure is performed to “improve” the toy model of the ratio of jet suppression factors.

Replacing the step function with this expression gives a very good approximation of the trend observed in Fig. 10. We plot the results using Eqs. (31) and (42), with the procedure to account for fluctuations in θ_c described above, in Fig. 11. For this qualitative comparison we have chosen $p_T/\omega_c \approx 1.2$ (roughly corresponding to $\hat{q} = 1.5$ GeV²/fm, $L = 3$ fm, and $p_T = 40$ GeV) and $n = 5$. The dotted lines show the result had we kept θ_c at a constant value, while the full lines include fluctuations of θ_c according to Eq. (30). As expected, fluctuations cause the scaling behavior of v_2 to drastically change shape, but the analytical trends are strikingly similar to the behavior obtained from the full numerical calculation.

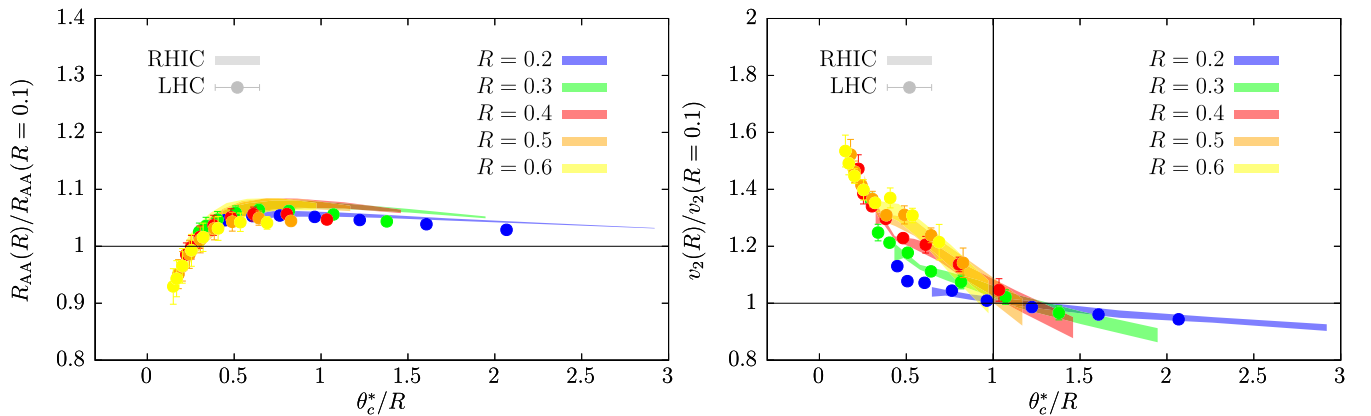


FIG. 10. Left: Double ratio of R_{AA} between jets with different radii R and jets with $R = 0.1$. Right: Double ratio of jet v_2 between jets with different radii R and jets with $R = 0.1$. Results are plotted against the variable θ_c^*/R , where θ_c^* is the most frequent value of the θ_c distribution for a given collision system and centrality. Jet momentum has been set to be $p_T \approx 38$ GeV.

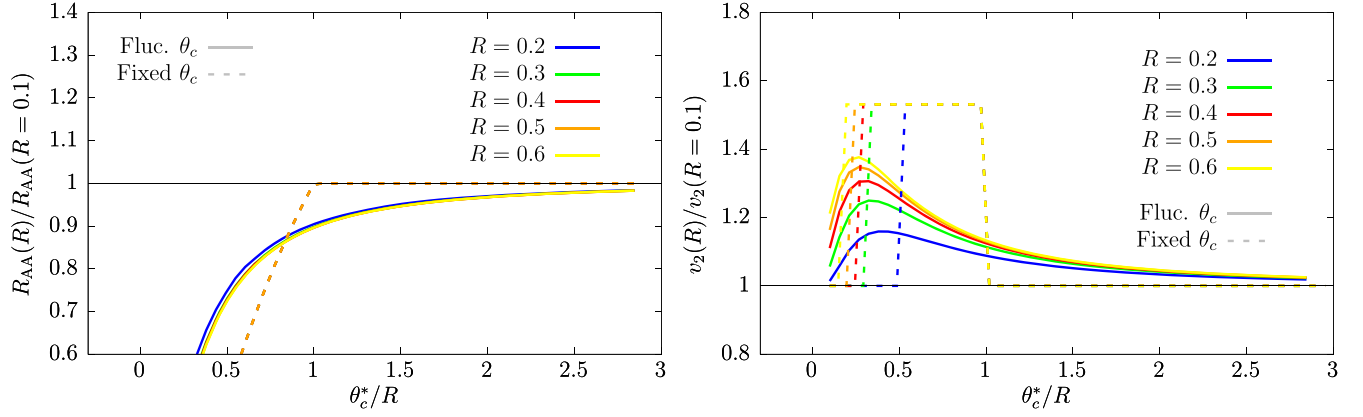


FIG. 11. Simple analytical estimates for the jet R_{AA} (left panel) and jet v_2 (right panel) ratios between jets with different $R > 0.1$ and those with $R = 0.1$, at a fixed $p_T/w_c \approx 1.2$ and $n = 5$. Results are presented as a function of θ_c^*/R . The dashed curves labeled with “Fixed θ_c ” ignore the effect of θ_c fluctuations within a given centrality class, while the solid curves, labeled as “Fluc. θ_c ,” do include them.

Some remarks are in order. Note that in the left panel of Fig. 11 the double ratio never goes above one in contrast to the left panel of Fig. 10, showing the full numerical results. Similarly, the toy-model predictions for the v_2 ratios in the right panel of Fig. 11 do not go below one at $\theta_c^*/R > 1$, in contrast to the right panel of Fig. 10. This is because our simplified formulas do not account for the possibility that radiation can stay within the cone, implying that larger R always translates into a larger energy loss. In this coherent regime, the physics of jet quenching is dominated by the energy loss of a single color charge and we know that opening up the cone angle leads straightforwardly to the slow recapturing of lost energy at large angles.

The fact that the clear jet v_2 scalings presented in the right panel of Fig. 10 can be obtained simply by using the most frequent value of θ_c in a given centrality class, c.f. Table I, is highly nontrivial. In comparison, the scaling observed for jet R_{AA} is not as striking given that jet suppression factor has a mild R dependence to begin with, for a given centrality.

Despite all the differences associated with the quenching of jets with different R , produced in different collision systems, with mediums of different initial energy densities and lifetimes, to mention a few of them, we predict that these jet R_{AA} and jet v_2 ratios should follow general trends if expressed in terms of θ_c^*/R . We reemphasize the fact that these trends are obtained from the results generated with the full semianalytical framework, within a realistic heavy-ion environment, which successfully describe jet R_{AA} and jet v_2 for essentially all jet p_T and R , centrality classes, and collision systems—after fitting a single parameter, $g_{\text{med}} \approx 2.25$.

V. CONCLUSIONS

In this paper we have provided the first analytical calculation of the elliptic flow coefficient v_2 for jets in heavy-ion collisions. Given that jets are extended

multipartonic objects, their degree of suppression due to medium interactions is in general larger than for single-inclusive hadrons. This fact leads, on average, to lower values of R_{AA} , the jet yields suppression observable, and to larger values of v_2 , the azimuthally asymmetric jet yields observable. A key ingredient of our framework is the determination of the quenched phase space of a jet, determined by the physics of color decoherence in the medium. At LL accuracy, dipoles with an angle smaller than the critical angle θ_c will not be resolved by the medium, and will lose energy as a single color charge, i.e., will lose less energy. The fact that the critical angle possesses a marked length dependence, $\theta_c \sim L^{-3/2}$, together with the fact that jets belonging to different centrality classes typically traverse very different amounts of QGP, naturally motivates the study of the physics of color decoherence for jet quenching observables as a function of centrality, targeting wide ranges of critical angles. The identification of such a program is one of the main conclusions of the present paper.

As a relevant phenomenological example, we have studied the jet v_2 evolution with centrality as a function of the jet opening angle R . We have compared to ATLAS data for narrow jets with $R = 0.2$, obtaining very good agreement, using the same framework and same single free parameter used in our previous work where we described the centrality evolution of inclusive jet suppression at LHC, for different jet radii [12]. In contrast to the R dependence of jet R_{AA} , we have found that there is a clear ordering for the jet v_2 results for jets with different R , where larger R jets always possess larger values of jet v_2 as compared to smaller R jets. The reason is tightly connected to the typical value of θ_c in a given centrality class. Given that v_2 quantifies the different degree of suppression of jets traversing different amounts of length (i.e., it is a length-differential observable), those jets whose paths lead to a larger value of θ_c (the short paths) will be less suppressed

than those whose paths lead to smaller values of θ_c (the longer paths), thereby raising the value of jet v_2 . This is so unless the jet as a whole is to be unresolved regardless of the path taken, i.e., when $R < \theta_c$. In this case, small differences in the traversed path length will not lead to different amounts of resolved phase space, reducing the size of jet v_2 when compared to those jets with $R > \theta_c$. We indeed observe how, as a function of centrality, jets whose R goes from being above to below the typical value of θ_c in that centrality change their behavior in relation to jets with other R , such as in relation to the smallest R jets studied in this work, $R = 0.1$, as shown in Fig. 8.

Usage of a semianalytical framework has allowed us to qualitatively explain, with simple expressions, this characteristic behavior of the R dependence of v_2 . These results suggest a representation of jet v_2 ratios in terms of the ratio between the two most relevant angular scales in this problem, i.e., as a function of θ_c/R . In Fig. 10 we have shown that this procedure leads to the approximate collapse of all such jet v_2 ratios, for any centrality and collision system, onto a single curve, for a fixed jet p_T . Our simple expressions, Eqs. (36) and (42), provide an analytical explanation for this behavior, provided that the substantial role of the θ_c fluctuations within a given centrality class are taken into account. The success of our framework to describe all presently available jet suppression data for any jet p_T , size R , centrality class, and collision system grants robustness to these predictions. They represent a transparent test of the physics of color decoherence which can be confronted with future precise experimental data to be measured both at RHIC and LHC. This dynamical picture can be tested by acquiring precise data on jet v_2 for as many cone sizes R as possible, in as many centrality classes as possible and in as many collision systems as possible (such as PbPb, AuAu, and OO collisions [124,125]).

In this work we have also revisited the predictions for jet R_{AA} , comparing our predictions with newly released experimental data and obtaining remarkably good agreement. Further directions include incorporating a running coupling between the jet parton and medium constituents, which is now fixed and is in practice the only free parameter of our framework (g_{med}). Another particularly relevant theoretical improvement will consist in the inclusion of theoretically well-motivated dynamics for quenching in the nonequilibrium regime. Quenching in the initial stages can have sizeable effects on jet v_2 , as shown in Fig. 7 using a simplified prescription, and in consistency with what was found for single-inclusive high- p_T hadron v_2 [110,111]. The strongest relative effect appears for the most central class, and is necessary to describe experimental data. Properly accounting for quenching in the initial stages is also necessary to understanding the possible existence of energy loss in small systems, such as proton-lead (pPb) or pp collisions.

The current implementation of medium response effects, now simply accounted for by a single parameter, R_{rec} , and

which have been shown to have important effects for large- R jets, such as $R > 0.6$, leaves room for improvement. A more precise treatment of the kinematics of the particles involved in the scattering processes, such as the recoiling constituents of the QGP, can be achieved by numerically computing bare quenching factors using QCD effective kinetic theory simulations [126,127]. Moreover, including the effects of thermodynamic gradients and flow on the stimulated radiation and broadening kernels recently computed in [128–132], so far neglected in most studies in the literature, will provide a more realistic description of the coupling between the jet and the dynamically evolving medium. In particular, the existence of a flowing medium implies that the thermalized energy and momentum deposited by the jet will flow along with this medium, leading to different wake patterns depending on the jet inmedium history [133,134]. Large R jets are also sensitive to the amount of quenching experienced by the recoiling jet in the same event due to the long-range correlations induced by the wake [135], provided that they are close enough in rapidity [136]. One way to account for these physics will be to use computationally efficient semianalytical calculations of the wake profiles [137,138] in order to determine the R_{rec} parameter jet by jet.

Furthermore, it will also be very interesting to extend this framework in order to go beyond inclusive jet observables, allowing it to becoming differential in jet substructure properties, or including the possibility of multijet events.

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APPENDIX: FULL SCAN ON THE R DEPENDENCE OF R_{AA} AND JET v_2 WITH CENTRALITY AT RHIC AND THE LHC

In this section we show a comprehensive set of results for jet R_{AA} and jet v_2 for LHC (Fig. 12) and RHIC (Fig. 13) collision systems. In the following we comment on the main aspects of the R and p_T dependence of these observables, which offer complementary information to that presented in the main text.

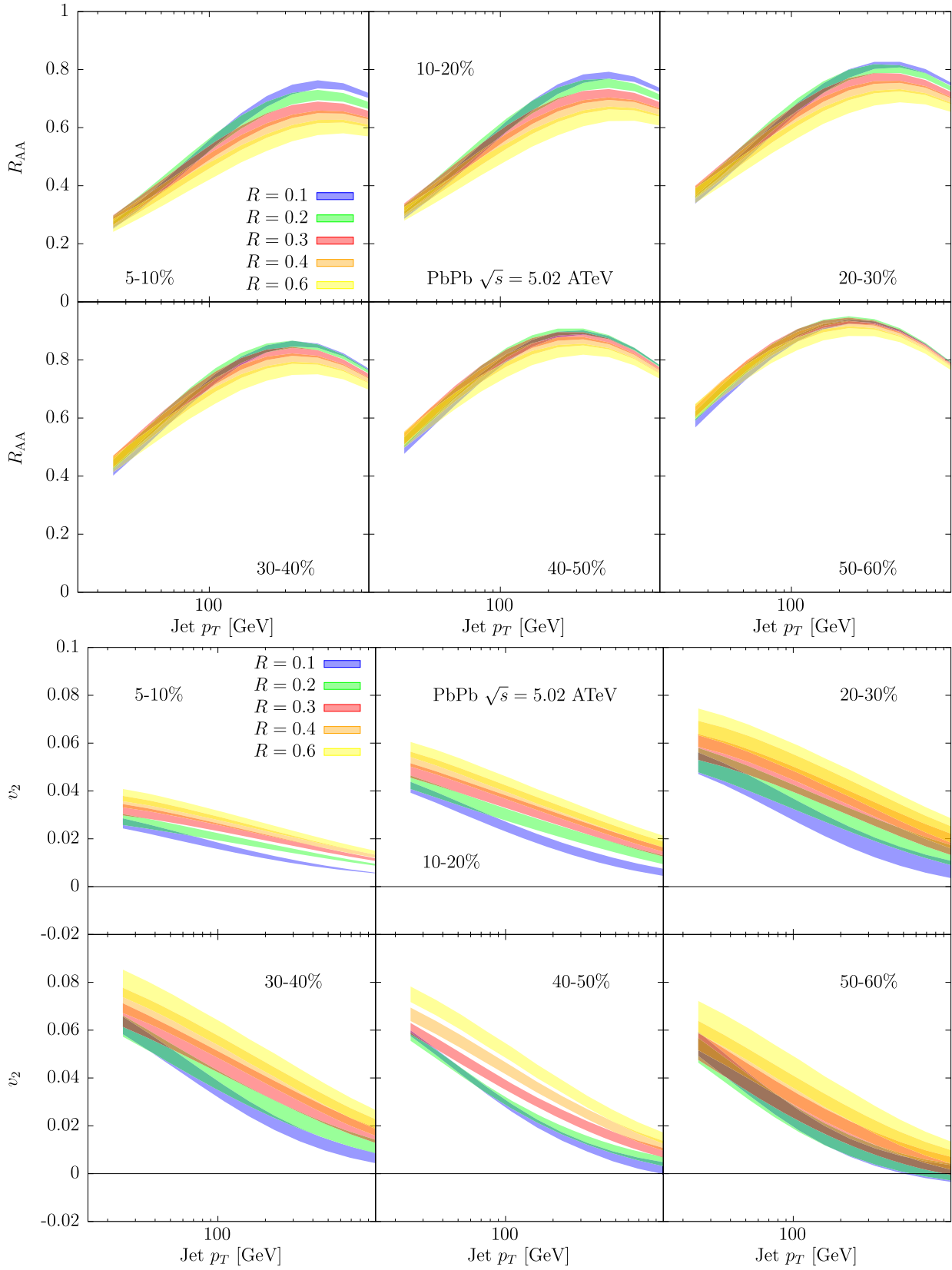


FIG. 12. Jet $R_{AA}(R, p_T)$ (top) and jet $v_2(R, p_T)$ (bottom) at LHC for PbPb collisions at $\sqrt{s} = 5.02$ A TeV.

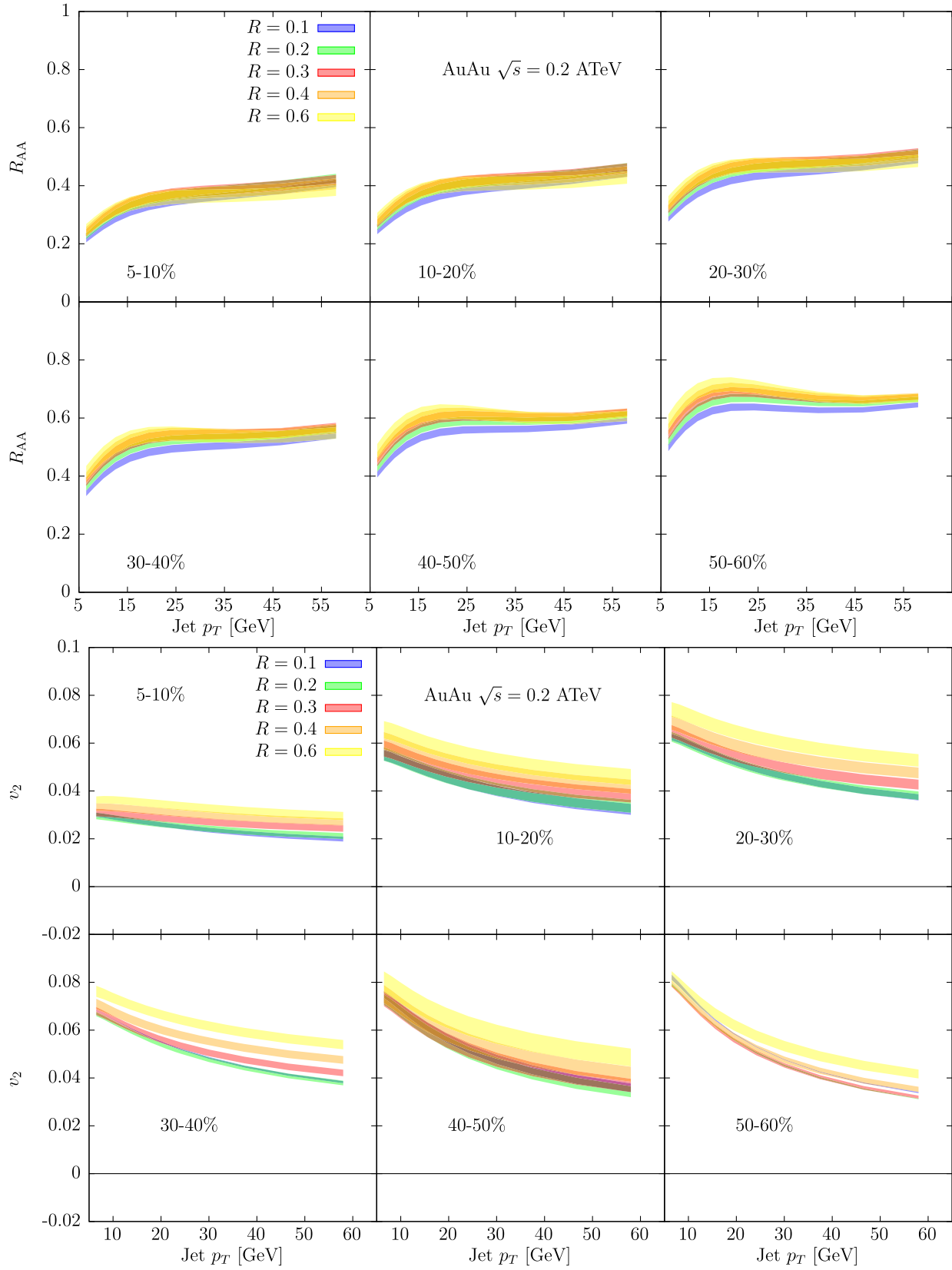


FIG. 13. Jet $R_{AA}(R, p_T)$ (top) and jet $v_2(R, p_T)$ (bottom) at RHIC for AuAu collisions at $\sqrt{s} = 0.2$ ATeV.

We first discuss the LHC results shown in Fig. 12. The results in the top panel indicate that jet R_{AA} for jets with different R are much more similar at low p_T than at high p_T , where larger R jets are more suppressed, specially for the most central classes. This is because of the competition between phase space and recapture of energy effects. The size of the resolved phase space grows with jet p_T and R . Even though a large R jet is more capable of retaining energy emitted from its core, the larger amount of energy loss sources present at high p_T trump this effect and the net result is an increased suppression. In contrast, at low p_T the recapture of the emitted energy plays a more important role, yielding overall a very mild dependence of suppression on R . As we move towards more peripheral classes, the size of the resolved phase space diminishes as the typical value of θ_c increases, yielding an increasingly flatter dependence with R across a wider jet p_T range. We even observe how for the most peripheral class, in the bottom right panel, large R jets are somewhat less suppressed than small R ones.

Interestingly, for the case of jet v_2 (lower panel of Fig. 12) we in general find that $v_2(p_T, R') \geq v_2(p_T, R)$ if $R' > R$, for any centrality class. This is explained in detail with simple analytical estimates in Sec. IV, and is in essence related to the interplay between the marked length dependence of θ_c and the fact that jet v_2 is a length-differential jet suppression observable. As the typical value of θ_c increases with decreasing centrality, we find that jet v_2 values tend to be similar among jets with $R \leq \theta_c$. This is specially so towards lower jet $p_T \approx 50$ GeV, where phase-space effects are not too large. The important remark to make is that, in contrast to the jet R_{AA} results, relative

differences in jet v_2 between small R and large R jets do in general remain large over a wide range of jet p_T and centrality classes.

By moving to the RHIC results of Fig. 13 we observe qualitatively similar trends. Jet R_{AA} , shown in the top panel, is naturally covering a much narrower jet p_T range due to smaller center of mass energies. The relatively small size of the resolved phase space implies that large R jets are slightly less suppressed than small R jets at fairly low values of jet p_T (where the applicability of our framework should be put to question), while yielding a very mild dependence towards higher jet $p_T \approx 40$ GeV. Again, however, we find that jet v_2 , shown in the lower panel of Fig. 13, does present the same ordering discussed above for the LHC results. To appreciate the interesting decorrelation between jet v_2 and jet R_{AA} behavior with varying R , it is instructive to note that even though jet R_{AA} can increase (less quenching) with increasing R , jet v_2 will always increase with increasing R . The fact that the size of the resolved phase space remains relatively small over the whole p_T range makes the grouping of jets whose R is smaller than the typical value of θ_c for a given centrality class clearer than it was in the case for the LHC results of Fig. 12.

We conclude by highlighting that the results of this section imply that the θ_c^*/R scalings predicted in Sec. IV for jet v_2 ratios are a very transparent manifestation of the physics of color decoherence in jet suppression, dissimilar to the case of the jet R_{AA} ratios, where relative differences among jets with different R are already small due to competing effects.

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