Jet correlation measurement in heavy-ion collisions: from RHIC to LHC

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Abstract. We attempt to deduce simple options of ‘jet quenching’ phenomena in heavy-ion collisions at \( \sqrt{s_{\text{NN}}} = 5.5 \text{ TeV} \) at the LHC from the present knowledge of leading-hadron suppression at RHIC energies. In light of the nuclear modification factor for leading particles we introduce the nuclear modification factor for jets, \( R_{\text{jet}}^{AA} \), and for the longitudinal momenta of particles along the jet axis, \( R_{pL}^{AA} \).

1. Introduction
At RHIC energies, hard probes, mainly light quarks and gluons, are experimentally accessible in heavy-ion collisions with sufficiently high rates. Their production is quantified using inclusive single-particle spectra at high momentum (\( p_T > 2 - 3 \text{ GeV} \)). The suppression of these high-momentum (leading) particles in central nucleus–nucleus relative to peripheral or pp collisions is one of the major discoveries at RHIC. In Au–Au collisions at various centre-of-mass energies (\( \sqrt{s_{\text{NN}}} = 62.4, 130 \), and \( 200 \text{ GeV} \)) the two experiments with high transverse-momentum capabilities, PHENIX and STAR, but also PHOBOS and BRAHMS, have measured:

- the suppression of single-particle yields at high \( p_T \) (\( \gtrsim 2 - 4 \text{ GeV} \)) [1–5];
- the disappearance, in central collisions, of jet-like correlations in the azimuthally-opposite side (away-side) of a high-\( p_T \) leading particle [6, 7] and, quite recently, the reappearance of the particles on the away-side manifested in low-momentum hadrons [8, 9];

The absence of these effects in d–Au collisions at the same centre-of-mass energy (\( \sqrt{s_{\text{NN}}} = 200 \text{ GeV} \)) [1, 10–12] confirms that in (central) nucleus–nucleus collisions final-state (as opposed to initial-state) effects modify the measured particle spectra.

The experimental observations have been explained in terms of various quenching models, where the energetic partons produced in the initial hard scattering ‘lose’ energy as a consequence of the interaction with the dense, partonic matter. Most models implement parton energy loss according to medium-induced gluon radiation of a hard parton traversing dense partonic matter of finite size [13–21]. Also hadronic interactions [22, 23] have been investigated and partially found to contribute to the observed depletion of the hadron spectra.

At the LHC at 30 times higher centre-of-mass energy, hard probes will be abundantly produced, even at energies of more than one order of magnitude higher than at RHIC. These energetic partons might be identified by their fragmentation into hadronic jets of high energy. In contrast with RHIC, the initial energy of \( E_T > 50 \text{ GeV} \) is high enough to allow the full reconstruction of the hadronic jet, even in the underlying heavy-ion environment [24]. Therefore, measurements of changes in the properties of identified jets in nucleus–nucleus with respect to pp collisions become possible.
In the present work we attempt to deduce perspectives for ‘jet quenching’ measurements at LHC given present data and models of hadron suppression at RHIC. In section 2 we introduce PQM, a model that describes high-$p_T$ suppression at RHIC and its prediction of the nuclear modification factor for LHC. In section 3 we outline the surface-emission effect which limits the sensitivity of leading particles to the density of the medium. In section 4 we present a PYTHIA simulation of jet quenching at LHC conditions and introduce the nuclear modification factor for jets, $R_{AA}^{jet}$, and for the longitudinal momenta of particles along the jet axis, $R_{AA}^{p_L}$.

2. Leading-particle spectroscopy at RHIC and LHC conditions

The effect of the medium on the production of a hard probe is typically quantified via the nuclear modification factor,

$$R_{AB}(p_T, \eta; b) = \frac{1}{\langle N_{coll}(b) \rangle} \times \frac{d^2 N^{\text{hard}}_{AB}/dp_T d\eta}{d^2 N^{\text{hard}}_{pp}/dp_T d\eta},$$

which measures the deviation of the nucleus–nucleus from the superposition of independent nucleon–nucleon collisions. In absence of strong nuclear initial-state effects it should be unity, if binary collision scaling holds. Figure 1 (left) shows $R_{AA}(p_T)$ for central Au–Au collisions at $\sqrt{s_{NN}} = 200$ GeV as published by PHENIX [3, 4] and STAR [2] together with a pQCD calculation of parton energy loss obtained with our Monte Carlo program PQM (Parton Quenching Model) [20].

PQM combines a recent calculation of the BDMPS-Z-SW parton energy loss [25] with a realistic description of the collision geometry at mid-rapidity. Our approach allows to calculate the transverse momentum and centrality dependence of single-hadron and di-hadron correlation suppression, as well as the ‘energy-loss induced’ azimuthal anisotropy of particle production in non-central collisions. The model has one single parameter, $k$, that sets the scale of the density of the medium. Once the parameter is fixed, e.g. on the basis of the nuclear modification data at $\sqrt{s_{NN}} = 200$ GeV, we scale it to different energies and collision systems assuming its proportionality to the expected volume-density of gluons. The application of our model is restricted to the high-$p_T$ region, above $4−5$ GeV at RHIC energies and above 10 GeV at LHC energy, since we do not include initial-state effects.
The leading-particle suppression in nucleus–nucleus collisions is obtained evaluating

\[
\frac{d^2\sigma_{\text{quenched}}}{dp_T d\eta} |_{\eta=0} = \sum_{a,b,j=q,\bar{q}} \int dx_a \, dx_b \, d\Delta E_j \, dz_j \, f_a(x_a) \, f_b(x_b) \frac{d^2\hat{\sigma}^{ab\to j X}}{dp_{T,j}^2 d\eta_j} |_{\eta_j=0} \times \\
\times \delta \left(p_{T,j}^{\text{init}} - (p_{t,j} + \Delta E_j)\right) P(\Delta E_j; R_j, \omega_{c,j}) \frac{D_{h/j}(z_j)}{z_j^2},
\]

in a Monte Carlo approach. Equation (2) describes the production of high-\(p_T\) hadrons within the perturbative QCD collinear factorization framework, at mid-rapidity, \(\eta = 0\), including medium-induced parton energy loss in the BDMPS-Z-SW formalism. \(f_a(b)\) is the parton distribution function for a parton of type \(a(b)\) carrying the momentum fraction \(x_a(b)\). \(\hat{\sigma}^{ab\to j X}\) are the partonic hard-scattering cross sections and \(D_{h/j}(z_j)\) is the fragmentation function, i.e. the probability distribution for the parton \(j\) to fragment into a hadron \(h\) with transverse momentum \(p_T = z_j p_{t,j}\). The modification with respect to the pp (vacuum) case is given by the energy-loss probability, \(P(\Delta E_j; R_j, \omega_{c,j})\), for the parton \(j\). Its two input parameters, the kinematical constraint and the characteristic scale of the radiated gluons, depend on the in-medium path length \(L\) of the parton and on the BDMPS transport coefficient of the medium, \(\tilde{q}\). The latter is defined as the average medium-induced transverse momentum squared transferred to the parton per unit path length. (See Ref. [20] on how we calculate the two parameters \(R\) and \(\omega_{c}\) for a given parton and realistic geometry.) In the original calculation [25], the energy-loss probability (quenching weight) is independent of the energy of the original parton allowing a parton with finite energy to radiate more than its energy. To account for finite parton energies we have introduced two different ways of constraining the weights, the non-reweighted and reweighted case. In the first case we constrain the loss to the energy of the parton, whenever the radiated energy is determined to be larger than that. In the second we require (by truncation of the distribution) the energy loss to be less than the energy of the parton. The resulting energy loss is larger in the non-reweighted case, because partons ‘thermalize’ with a probability \(\int E_\text{\footnotesize{\textsc{hadron}}} \, \rho(\epsilon) \, d\epsilon\) it is argued [21, 25] that the difference in the values of the observables for the two approaches illustrates the theoretical uncertainties of the model.

Coming back to fig. 1 (left) we note that for the chosen value of \(k\) the calculation reproduces \(R_{AA}(p_T)\) for central Au–Au collisions at \(\sqrt{s_{NN}} = 200\) GeV. In Ref. [20] we consistently compare model predictions to various high-\(p_T\) observables at RHIC energies. We also show that we need to scale \(k\) by a factor of seven to obtain the RHIC prediction of \(R_{AA}\) at \(\sqrt{s_{NN}} = 5.5\) TeV shown in fig. 1 (right). Our prediction for the LHC is in agreement, both in the numerical value and in the \(p_T\)-dependence, with that obtained in Ref. [21], while it is quite different from that calculated in Ref. [17].

3. Properties of leading-hadron suppression

At LHC energies, the expected \(R_{AA}\) for central Pb–Pb collisions at LHC as a function of \(p_T\) is rather flat, i.e. almost \(p_T\)-independent, similar to what is observed at RHIC at \(p_T > 4\) GeV. To study this \(p_T\)-dependence we show in fig. 2 \(R_{AA}(p_T)\) in central Pb–Pb at LHC for \((\tilde{q}) = 12\) GeV\(^2\)/fm (corresponding to the value found at RHIC), \((\tilde{q}) = 24\) GeV\(^2\)/fm and \((\tilde{q}) = 98\) GeV\(^2\)/fm (the value obtained by the scaling from RHIC to LHC), as well as the result of a calculation with fixed \(\tilde{q} = 10\) GeV\(^2\)/fm and fixed length of 4.4 fm. Clearly, the fixed-length calculation shows a stronger \(p_T\)-dependence than the PQM calculation.

The flatness of \(R_{AA}\) is a consequence of

- the steeply falling production cross-section, \(\propto \left(\frac{1}{p_T}\right)^n(p_T)\), where \(n(p_T)\) is rising from about 7 to 12 (RHIC) and from 6 to 7 (LHC) in the relevant \(p_T\) regime;
- the emission from the surface, which for large medium densities dominates [26];
as we explain in the following. \( R_{AA} \), eq. (1), at mid-rapidity, can be approximated by

\[
R_{AA}(p_T) = \int d\Delta E P(\Delta E, p_T + \Delta E) \frac{dN^{pp}(p_T + \Delta E)}{dp_T} / \frac{dN^{pp}(p_T)}{dp_T},
\]

where, \( dN^{pp}/dp_T \) is the spectrum of hadrons (or partons) in the case of no medium (i.e. pp neglecting initial state effects). The suppression computed with eq. (3) is found to give a rather good approximation to the one computed with PQM or with the full formula, eq. (1). In the case the production spectrum is (approximately) exponential the \( p_T \)-dependence cancels in the ratio and we find \( R_{AA} \) to be (approximately) independent of \( p_T \). At RHIC this is the case at about \( p_T \geq 30 \text{GeV} \). Below that value and at the measured values of 5–12 GeV, as well as at the LHC (\( n(p_T) \leq 7 \)), the spectrum is given by a power-law and it was expected [27] that \( R_{AA} \propto (1 + \Delta E/p_T)^{-n(p_T)} \), i.e. reaching unity for large \( p_T \).

However, this neglects the fact that for dense media surface emission or, more generally, the probability to have no energy-loss, \( P(\Delta E = 0, E) \), plays a significant role, an effect which is even more pronounced at low \( p_T \) (compared to \( \omega_c \)). To simplify our argumentation we allow either no loss (\( \Delta E = 0 \)) or complete loss (\( \Delta E = E \)) in the non-reweighted case,

\[
P(\Delta E, E) = p_0 \delta(\Delta E) + (1 - p_0) \delta(\Delta E - E).
\]

1 Inserting the constrained weight into eq. (3) we obtain

\[
R_{AA}(p_T) = p^* + (1 - p^*) \frac{dN^{pp}(2p_T)}{dp_T} / \frac{dN^{pp}(p_T)}{dp_T}.
\]

It is obvious that eq. (4) is just a crude approximation, but it demonstrates that the value of \( R_{AA} \) is dominated by the fraction of hadrons (or partons), which escape without losing much of their energy. For the simple power-law production spectrum the contribution from higher \( p_T \) is suppressed by about \((1 + \Delta E/p_T)^{-n(p_T)}\). Taking into account only fixed values of \( \hat{q} \) and \( L \) the probability \( p^* \) is given by the discrete weight, \( p_0 \), at \( R = 0.5 \hat{q} L^3 \), amounting to about 0.05(0.002) for quarks(gluons). Note that for a proper calculation one must take into account the right production ratio of quarks-to-gluons. For realistic

1 For a dense medium the constrained weights at low parton energy indeed do have a sharp peak at zero and at maximum possible energy loss, whereas the values in between are negligible.
The picture of surface-emission at LHC, even for very large-momentum hadrons, is supported by fig. 3 (top). It visualizes the distribution of production points \((x_0, y_0)\) in the transverse plane for partons, which escape from the dense overlap region and yield hadrons with \(p_T^{\text{hadron}} > 50\) GeV. The chosen values of the medium density correspond to the value found at RHIC and the expectation for LHC. The corresponding in-medium path-length distributions, fig. 3 (bottom), reveal that the average ‘thickness’ of the emission surface is restricted to less than 2 fm: even hadrons with \(p_T > 50\) GeV are emitted dominantly from the surface.

The dominance of the surface effect limits the sensitivity of leading particles to the density of the medium, mainly for experimentally accessible low-\(p_T\) range at RHIC. This is demonstrated in fig. 4 where we show the dependence of \(R_{AA}\) on the average transport coefficient evaluated with PQM for 0–10% most central collisions at \(\sqrt{s_{NN}} = 5.5\) TeV. For 10 GeV hadrons the nuclear modification factor is sensitive to average medium densities up to about 15 GeV^2/fm, and for 100 GeV hadrons the sensitive regime might widen to average values of about 30.

4. Jet spectroscopy and modification of jet properties

At LHC, leading-hadron spectroscopy will naturally extend to jet spectroscopy. Simulations show that jets with transverse energies of more than 50 GeV are identifiable and reconstructible on an event-by-event
basis, even in most central collisions; however with severe limitations in the resolution of the jet energy. The interaction of the hard partons with the medium created in these collisions is expected to manifest in the modification of jet properties deviating from known fragmentation processes in vacuum. Calculations predict that the additional gluons radiated by the original parton remain (fragment) inside the jet cone, although redistributed in transverse phase space [25, 28]. The corresponding jet-production cross section is expected to follow binary scaling. However, the jet shape is claimed to broaden and the jet multiplicity to soften and increase [29]. Ideally, if one reconstructs the hadronic energy for outstanding high-energy jets the jet energy may be associated with the original energy of the parton. In combination with results at lower parton energies from correlation methods jet measurements at higher energy might complete the picture of medium-induced parton-energy-loss phenomena. In the following it is our aim to introduce a simulation and analysis of jet quenching effects. Details can be found in Ref. [24]. Unquenched jets representing the pp reference are generated using the PYTHIA generator [30]. Quenched signal jets are simulated using a modified version of PYTHIA.\footnote{The quenching is implemented as an afterburner to standard PYTHIA developed by A. Morsch.} It introduces parton energy loss via final-state gluon radiation, in a rather ad-hoc way: Before the partons originating from the hard 2-to-2 process (and the gluons originating from ISR) are subject to fragment, they are replaced according

\[
\text{parton}_i(E) \rightarrow \text{parton}_i(E - \Delta E) + n(\Delta E) \text{gluon}(\Delta E/n(\Delta E))
\]

conserving energy and momentum. $\Delta E$ is calculated by PQM for the non-reweighted case and depends on the medium density, parton type and parton production point of the jet system in the collision overlap region. The number of radiated gluons, $1 \leq n(\Delta E) \leq 6$, is determined by the condition that each gluon must have less energy than the quenched parton from which it was radiated away.

For the analysis we have prepared events containing jets simulated with the modified PYTHIA version for $\langle q \rangle = 1.2, 12$ and $24 \text{ GeV}^2/\text{fm}$. These signal events are then embedded into 0–10% central HIJING [31] events. The choice of first value implies only a very little modification of the embedded quenched jets compared to embedding oj jets with standard PYTHIA. The second corresponds to the case found to describe the $R_{AA}$ at RHIC, whereas the third is a conservative choice below the extrapolation from RHIC to LHC.

The jet finding is based on final particles (no detectors effects) at mid-rapidity ($-1 < \eta < 1$). We use the standard ILCA cone finder [32]. However, to cope with the soft heavy-ion background we restrict the

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure4.png}
\caption{$R_{AA}$ as a function of the average transport coefficient, $\langle \hat{q} \rangle$, for 10 and 100 GeV hadrons in 0–10% central Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.5$ TeV. The calculation is for the non-reweighted case.}
\end{figure}
cone radius to $R = 0.3$ and apply a particle $p_T$-cut of $p_T > 2$ GeV. The same settings are used for the reference jet measurement in pp. The jet energy resolution at 50 GeV is about 80%, and more than 90% at 100 GeV.

4.1. Nuclear modification factor for jets

It has been shown [29] that the medium-induced broadening of the jet reduces the energy inside $R = 0.3$ by $\sim 15\%$ and by $\sim 7\%$ for jets with $E_T = 50$ and $E_T = 100$ GeV, respectively, and already at $R = 0.7$ the effect reduces to about $2\%$. However, one must be cautious about these findings, since this prediction has been calculated assuming a rather low value of the gluon density at LHC conditions. We have verified (without embedding in HIJING) that for $R = 0.7$ more than $90\%$ of the energy remains in the jet cone. Up to $45\%$ of the rise at jet energies lower than 50 GeV can be attributed to the contribution of the underlying heavy-ion background to the jet signal. It may well be that jet suppression (at fixed value of $R$) may be observed in the suppression of inclusive single-jet spectra in central nucleus–nucleus relative to peripheral or pp collisions. Analogous to eq. (1), we define the nuclear modification factor for jets, $R_{AA}^{jet}(E_T, \eta; b)$, according to

$$R_{AA}^{jet}(E_T, \eta; b) = \frac{1}{\langle N_{coll}(b) \rangle} \times \frac{d^2N_{jet}^{AA}}{dE_T d\eta} \times \frac{d^2N_{jet}^{pp}}{dE_T d\eta}.$$ (5)

Equation (5) implicitly depends on the (fixed) value of $R$ used to find the jets. Figure 5 shows $R_{AA}^{jet}(E_T)$ for different quenching scenarios in 0–10% central Pb–Pb.

![Figure 5](image)

**Figure 5.** The nuclear modification factor for jets, $R_{AA}^{jet}(E_T)$, as a function of the reconstructed jet energy, $E_{T \text{rec}}$, for different quenching scenarios in 0–10% central Pb–Pb. The jets in pp and Pb–Pb are identified with the cone finder using $R = 0.3$ and $p_T > 2$ GeV.

4.2. Longitudinal momentum modification factor

The signature of medium-induced gluon radiation should be visible in the modification of the jet fragmentation function as measured through the longitudinal and transverse momentum distributions of associated hadrons within the jet. The momenta parallel to the jet axis, $p_L = p_{hadron} \cos(\theta_{jet}, \theta_{hadron})$, are expected to be reduced (jet quenching), while the momenta in the transverse direction, $j_T = p_{hadron} \sin(\theta_{jet}, \theta_{hadron})$, to be increased (transverse heating). In our simulation for technical reasons we
can not give an explicit transverse-momentum kick to the radiated gluons. Therefore, we concentrate on the changes of the longitudinal fragmentation.

We define the longitudinal momentum modification factor according to

$$R_{pL}^{AA}(p_L) = \frac{N_{jets}^{AA}}{N_{jets}^{pp}} \frac{dN_{AA}/dp_L}{dN_{pp}/dp_L}$$

(6)

where for proper normalization $N_{jets}^{AA}$ and $N_{jets}^{pp}$ account the total number of jets that were found in A–A and pp, respectively. Equation (6) implicitly depends on the transverse jet-energy range of jets taking into account in the distributions. Figure 6 shows $R_{AA}(p_L)$ for jets with reconstructed energies of $E_T^{rec} > 80$ GeV in 0–10% central Pb–Pb compared for the different quenching scenarios.

**Figure 6.** Longitudinal momentum modification factor, $R_{pL}^{AA}$, for jets with $E_T^{rec} > 80$ GeV and different quenching scenarios in 0–10% central Pb–Pb. The jets in pp and Pb–Pb are identified with the cone finder using $R = 0.3$ and $p_T > 2$ GeV.

The expected behavior is clearly visible: higher medium density leads to stronger suppression for large longitudinal momenta and enhancing of smaller momenta. At low $p_L$ the effect becomes most apparent. At these low momenta there is an moderate additional contribution from particles which belong to the underlying heavy-ion background, which amounts to about 20% for $p_L < 5$ GeV. It might be interesting to study the slope of $R_{AA}(p_L)$ as a function of $E_T$.

5. Conclusions
We have presented the PQM model [20] which is able to describe high-$p_T$ suppression effects at RHIC. Its application to LHC conditions leads to the expectation that $R_{AA}$ is essentially constant with $p_T$, also for very high momenta up to 100 GeV. The nuclear modification factor for leading hadrons is largely dominated by surface effects. This limits its sensitivity to the density of the created medium. The measurement of reconstructed jets above the underlying heavy-ion background may allow to probe the medium to deeper extents. In a simulation of jet quenching at LHC we have studied two observables: the nuclear modification factor for jets, $R_{AA}^{jet}$, and the nuclear modification factor for the longitudinal momenta of particles along the jet axis, $R_{pL}^{AA}$. The definitions implicitly depend on the parameters used to find the jets. For jets with $E_T > 50$ GeV both are suitable measures to quantify deviations from the pp case: $R_{AA}^{jet}(E_T)$ decreases with increasing $E_T$ and $R_{pL}^{AA}$ is enhanced at low $p_L$ and suppressed at high $p_L$. However, for jets with $E_T < 50$ GeV suffer from the contamination of the jet cone.
by uncorrelated particles from the underlying background which at these energies severely influences measured jet properties.

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