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Collimation of energy in medium-modified QCD jets

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Abstract: The collimation of energy inside medium-modified jets is investigated in the leading logarithmic approximation of QCD. The Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) evolution equations are slightly modified by introducing splitting functions enhanced in the infrared sector. As compared to elementary collisions in the vacuum, the angular distribution of the jet energy is found to broaden in QCD media.

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1 Introduction

The jet quenching has been established as the phenomenon of strong high-p_T suppression of partons produced in heavy ion collisions [1–3]. Experimentally, it has been confirmed that highly virtual partons produced in such reactions suffer an energy degradation prior to their hadronization in the vacuum. The high-p_T partons energy is radiated away as gluon bremsstrahlung, thus increasing the amount of such gluons after several interactions of the leading parton with the scattering centers in the quark gluon plasma (QGP) occur [4]. Jet production in hadronic collisions is a paradigmatic hard QCD process. An elastic ($2 \to 2$) or inelastic ($2 \to 2+X$) scattering of two partons from each one of the colliding hadrons (or nuclei) results in the production of two or more partons in the final-state. At high p_T , the outgoing partons have a large virtuality Q which they reduce by subsequently radiating gluons and/or splitting into quark-antiquark pairs. Such a parton branching evolution is governed by the QCD branching probabilities given by the DGLAP equations [5] down to virtualities $\mathcal{O}(1\,\mathrm{GeV}^2)$.

In this paper we study the energy collimation of jets produced in heavy ions collisions from a toy QCD-inspired model introduced by Borghini and Wiedemann [6]. In this model, the DGLAP splitting functions are enhanced in the infrared region in order to mimic medium-induced soft gluon radiation. More precisely, the $1/x$ dependence of the branching probabilities $q \to qg$, $g \to gg$ is enhanced by introducing the factor $N_s = 1 + f_{med}$ in $P_{ab}(x) = (1 + f_{med})/x + \mathcal{O}(1)$. In order to quantify the collimation of energy, we consider a jet of half opening angle Θ_0 initiated by a parton A, followed by the production of a subjet of half opening angle Θ (with $\Theta < \Theta_0$) initiated by a parton B with large energy fraction $x \sim 1$, which contains the bulk of the jet energy registered by the calorimeter. Experimentally, this corresponds to the calorimetric measurement of the energy flux deposited within a certain solid angle Θ. The first version of this quantification was given in [7, 8] in the vacuum and it was found to scale like,

$$\frac{\Theta}{\Theta_0} = \left(\frac{E\Theta_0}{\Lambda}\right)^{-\gamma_A(x)},$$

where $\gamma_A(x)$ ($A = q, \bar{q}, g$) is a certain function derived from the DGLAP evolution equations at large x and $\Lambda \equiv \Lambda_{QCD}$ is the mass scale of QCD. In this context, we will introduce the N_s dependence in $\gamma(x) \to \gamma(x, N_s)$ in order to compare our results in the medium ($\gamma(x, N_s)$) with previous results in the vacuum ($\gamma(x) = \gamma(x, 1)$). The same logic was applied to the study of the collimation of average multiplicities in [9], where jets were found to broaden as compared to those in the vacuum after inserting the N_s-dependence in the evolution.

The same model has been largely used, i.e. for the estimation of other observables like medium-modified inclusive p_T-distributions [10] and fragmentation functions [11], in all cases results look promising and appealing for future comparison with experimental data. The large suppression of these distributions [10, 11] at large x (high-p_T) and the enhancement at small x (small-p_T) are respectively good examples of the jet quenching in dense QGP.

2 Theoretical framework

Let us start from the production of one gluon or quark ($A = g, q, \bar{q}$) initiated jet with energy E, followed by its fragmentation into another parton B at angle Θ such that, $Q_0/E \leq \Theta \leq \Theta_0$, where Θ_0 is the half opening angle of the jet and Q_0, the minimal transverse momentum of the emitted partons before the hadronization into hadrons occurs. The collimation of the energy is characterized by the energy fraction $x \sim 1$ where the bulk of the jet energy inside the given cone $\Theta < \Theta_0 \ll 1$ is deposited. The probability for the energy fraction x to be deposited in a cone of aperture Θ is related to the inclusive spectrum of partons through the formula [7],

$$D_A(x, E\Theta_0, E\Theta) = \sum_{B=g,q} D_A^B(x, E\Theta_0, E\Theta), \tag{1}$$

where the nature of partons B is not identified. In the above relation (1) two scales $E\Theta_0$ and $E\Theta$ have been written in the argument of the inclusive spectrum so as to account for the given process in the energy range $E\Theta \leq Q \leq E\Theta_0$; this notation will be shortened below. In order to quantify the collimation of energy, it should be considered that the deposited fraction of the jet energy is large, that is for $x \to 1$. The DGLAP evolution equations in the vacuum takes the simple form [5],

$$\frac{d}{d\ln Q^2} D(x, Q^2) = \frac{\alpha_s(Q^2)}{4\pi} \int_x^1 \frac{dz}{z} P(z) D\left(\frac{x}{z}, Q^2\right), \tag{2}$$

where P is the *Hamiltonian* matrix of splitting functions characterizing the parton splitting probabilities. The splitting functions are accurately obtained from the FO approach in perturbation theory for small coupling constant $\alpha_s \ll 1$. The coupling constant is written in the form,

$$\alpha_s(Q^2) = \frac{\pi}{N_c \beta_0 \ln(Q^2/\Lambda_{QCD}^2)},$$

where $\beta_0 = \frac{1}{4N_c}(\frac{11}{3}N_c - \frac{4}{3}T_R)$ is the first coefficient of the beta function and $\Lambda_{QCD} \approx 250$ MeV. In Mellin space $\mathcal{D}(j, Q^2)$ is obtained from the transformation

$$\mathcal{D}(j, Q^2) = \int_0^1 dx\, x^{j-1} D(x, Q^2),$$

such that the convolution (2) simply becomes

$$\frac{d}{d\ln Q^2} \mathcal{D}(j, Q^2) = \mathcal{P}(j) \mathcal{D}(j, Q^2),$$

or more explicitly rewritten in the matrix form at LO

$$\frac{d}{d\xi} \begin{pmatrix} \mathcal{D}_{q_{\rm NS}}(j,\xi) \\ \mathcal{D}_{q_{\rm S}}(j,\xi) \\ \mathcal{D}_g(j,\xi) \end{pmatrix} = \begin{pmatrix} \mathcal{P}_{qq}(j) & 0 & 0 \\ 0 & \mathcal{P}_{qq}(j) & \mathcal{P}_{qg}(j) \\ 0 & \mathcal{P}_{gq}(j) & \mathcal{P}_{gg}(j) \end{pmatrix} \begin{pmatrix} \mathcal{D}_{q_{\rm NS}}(j,\xi) \\ \mathcal{D}_{q_{\rm S}}(j,\xi) \\ \mathcal{D}_g(j,\xi) \end{pmatrix}, \tag{3}$$

where $\mathcal{D}_{q_{\rm NS}}$ and $\mathcal{D}_{q_{\rm S}}$ stand respectively for the flavor non-singlet (or valence) and flavor-singlet quark distributions, and $\mathcal{P}_{ik}(j)$ is the Mellin transform of the LO splitting functions. The variable [5]

$$\xi(Q^2) = \frac{1}{4N_c\beta_0} \ln\left(\ln \frac{Q^2}{\Lambda_{QCD}^2}\right), \quad \frac{d}{d\ln Q^2} = \frac{\alpha_s}{4\pi} \frac{d}{d\xi},$$

is introduced for the sake of simplicity. In order to account for the medium-induced gluon radiation in heavy-ion collisions, we make use of the QCD-inspired model proposed in [6], which allows for a simple computation of the equations at large x. In this model the infrared parts of the splitting functions are enhanced by the factor $N_s = 1 + f_{med}$, where $f_{med} > 0$ accounts for medium-induced gluon radiation. Realistic values of f_{med} $(= 0.6, 0.8)$ are extracted from fits of the nuclear modification factor R_{AA} [3] to the RHIC data [2]; for the LHC we will take $f_{med} = 1$. The medium-modified splitting functions read,

$$P_{gg}(z) = 4N_c \left[\frac{N_s}{z} + \left[\frac{N_s}{1-z} \right]_+ + z(1-z) - 2 \right], \quad P_{gq}(z) = 2\, T_R [z^2 + (1-z)^2], \tag{4a}$$

$$P_{qg}(z) = 2\, C_F \left(\frac{2N_s}{z} + z - 2 \right), \quad P_{qq}(z) = 2\, C_F \left(\left[\frac{2N_s}{1-z} \right]_+ - 1 - z \right), \tag{4b}$$

with the $[\dots]_+$ prescription defined as $\int_0^1 dx [F(x)]_+ g(x) \equiv \int_0^1 dx F(x)[g(x) - g(1)]$. Performing the Mellin transform of Eq. (4a,4b) gives [11]

$$\begin{aligned}
\mathcal{P}_{gg}(j) &= -4N_c \left[N_s\, \psi(j+1) + N_s \gamma_E - \frac{N_s - 1}{j} - \frac{N_s - 1}{j-1} \right] \\
&\quad + \frac{11 N_c}{3} - \frac{2 n_f}{3} + \frac{8 N_c (j^2 + j + 1)}{j(j^2-1)(j+2)}, && \text{(5a)} \\
\mathcal{P}_{gq}(j) &= T_R\, \frac{j^2 + j + 2}{j(j+1)(j+2)}, && \text{(5b)} \\
\mathcal{P}_{qg}(j) &= 2\, C_F\, \frac{(2N_s - 1)(j^2 + j) + 2}{j(j^2-1)}, && \text{(5c)} \\
\mathcal{P}_{qq}(j) &= -C_F \left[4N_s\, \psi(j+1) + 4N_s \gamma_E - 4\frac{N_s - 1}{j} - 3 - \frac{2}{j(j+1)} \right], && \text{(5d)}
\end{aligned}$$

where $\psi(j) = \frac{d}{dj} \ln \Gamma(j)$ is the digamma function with $\gamma_E \approx 0.5772$ the Euler constant. It can be easily checked that Eq. (5) reduces to the ordinary splitting functions given in [5, 7] after setting $N_s = 1$. In the large $z \sim 1$ limit, or equivalent large $j \gg 1$ we are interested in, the expressions of the anomalous dimensions (5a) and (5d) can be re-expressed in the form

$$\mathcal{P}_{qq}(j) \approx 4 C_F N_s \left(-\ln j + \frac{3}{4N_s} - \gamma_E \right), \quad \mathcal{P}_{gg}(j) \approx 4 N_c N_s \left(-\ln j + \frac{\beta_0}{N_s} - \gamma_E \right), \tag{6}$$

where the asymptotic behavior $\psi(j+1) \approx \ln j$ has been replaced for $j \gg 1$. The Mellin transform can be inverted, with the inversion given by

$$D(x, \Delta\xi) = \frac{1}{2\pi i} \int_C dj\, x^{-j} \mathcal{D}(j, \Delta\xi), \tag{7}$$

where the contour C in the complex plane is parallel to the imaginary axis and lies to the right of all singularities of $\mathcal{D}(j, \Delta\xi)$. In (7)

$$\Delta\xi = \frac{1}{4 N_c \beta_0} \ln \left[\frac{\ln\left(\frac{E\Theta_0}{\Lambda_{QCD}}\right)}{\ln\left(\frac{E\Theta}{\Lambda_{QCD}}\right)} \right].$$

[3] R_{AA} corresponds to the ratio of medium-modified and unmodified single inclusive hadron spectra.

Integrating (3) after inserting (6) leads to the valence distribution at large $x \sim 1$ [5],

$$D_A^A(x, \Delta\xi) \simeq (1-x)^{-1+4C_A N_s \Delta\xi} \frac{\exp[4C_A N_s(\frac{3}{4N_s} - \gamma_E)\Delta\xi]}{\Gamma(4C_A N_s \Delta\xi)}. \tag{8}$$

where $\beta_0 = 3/4$ for $n_f = 3$ was replaced from (6) and the N_s dependence is new. The behavior of the valence distribution (8) when Θ changes can be studied by taking its derivative with respect to ξ. In [7], the valence distribution (8) was shown to present a maximum at some angle Θ for fixed values of x where the bulk of the jet energy is concentrated. Taking $\partial_\xi D_A^A = 0$ results in,

$$\psi(4C_A N_s \Delta\xi) = \ln(1-x) + \frac{3}{4N_s} - \gamma_E, \tag{9}$$

where here again $\psi(x)$ is the digamma function. For $\Delta\xi \to 0$, or $\Theta \to \Theta_0$, almost the whole energy will be deposited inside the cone Θ_0 while for Θ decreasing down to $\Theta \geq E/\Lambda$ the emission outside the cone increases. The equation (9) with the N_s dependence is the main result of this paper. It should be inverted numerically in order to provide the behavior of the ratio Θ/Θ_0 as a function of the energy of the leading parton and the nuclear factor N_s for fixed values of $x \sim 1$. Symbolically, the inversion can be written in the simple form,

$$\frac{\Theta}{\Theta_0} = \left(\frac{E\Theta_0}{\Lambda}\right)^{-\gamma_A(x,N_s)}, \quad \gamma_A(x, N_s) = 1 - \exp\left[-\frac{N_c\beta_0}{C_A N_s}\psi^{-1}\left(\ln(1-x) + \frac{3}{4N_s} - \gamma_E\right)\right], \tag{10}$$

such that, the collimation is stronger for higher energies in both vacuum and medium jets. As observed in (10), the N_s-dependence contained in the denominator of the exponent should soften the collimation of jets produced in heavy ion collisions. Furthermore, the function $\gamma_A(x, N_s)$ appearing in the exponent provides the slope of the collimation as a function of x and N_s. Of course, though (10) is model dependent, it may capture the main features of a more realistic QCD energy loss picture in heavy ion collisions. For instance, it would be interested to extend the vacuum result ((10) with $N_s = 1$) to the energy loss model introduced in [12], which modifies the DGLAP splitting functions by accounting for the medium length L and the transport coefficient $\hat{q}$; however, it stays out of the scope of the present paper.

3 Phenomenology

We display the collimation of energy through the ratio $\frac{\Theta}{\Theta_0}$ from the numerical inversion of (9). Let us first consider a jet produced in a heavy ion collision at energy scales greater than 100 GeV at the LHC and let us set $\Theta_0 \sim 1$ for simplicity. For the gluon and quark jets, one sets $C_A = N_c = 3$ and $C_A = C_F = 4/3$ respectively in (9). In figure 1 and 2, we display the collimation of energy for two fixed values of energy fractions $x = 0.5$ and $x = 0.9$ which are in agreement with the large x ($x \sim 1$) approximation applied in this frame. As expected from (9) and observed in figures 1 and 2, the curves of the collimation in the medium ($f_{med} = 0.6, 0.8, 1$) are shifted towards higher values of the subjet opening angle Θ as compared with the curve for Θ in the vacuum. Moreover, the slope of the collimation decreases as f_{med} increases. For fixed Θ_0, E and x, the same amount of energy contained inside a subjet of opening angle Θ_1 in the vacuum is distributed over a broader angular aperture Θ_2, such that $\Theta_2 > \Theta_1$. It follows

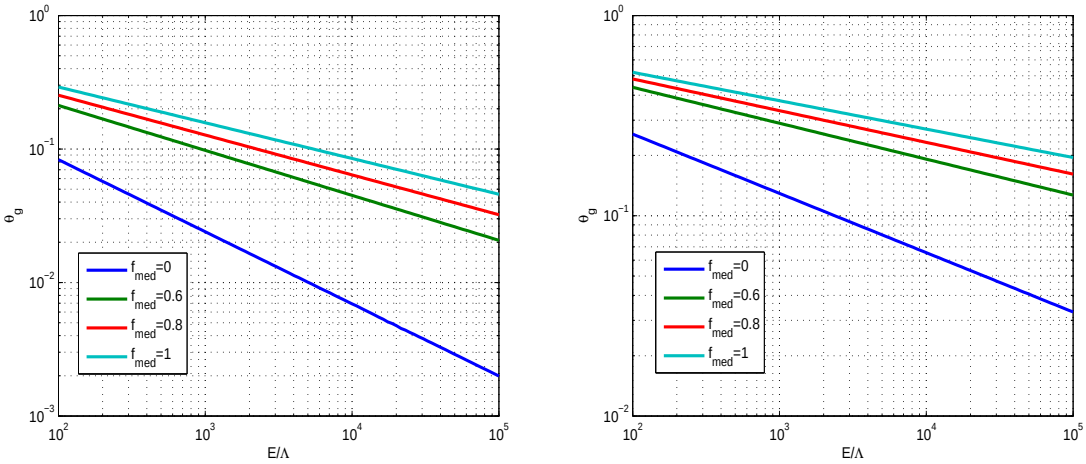

Figure 1: Collimation of energy inside a gluon jet for $x = 0.5$ (left) and $x = 0.9$ (right).

that gluon jets are less collimated than quark jets in the vacuum [7] and not surprisingly, also in the medium. For $x = 0.5$ jets are more collimated than for higher values of the energy fraction. Finally, it is straightforward to check that the values of $\gamma_A(x, N_s = 1)$ for $x = 0.5, 0.9$ are in agreement with the expected values $\gamma_g(0.5, 1) \approx 0.54$, $\gamma_g(0.9, 1) \approx 0.30$, $\gamma_q(0.5, 1) \approx 0.83$ and $\gamma_q(0.9, 1) \approx 0.55$ in the vacuum [7, 8]. This analysis could be performed in future measurements by the LHC experiments as a probe for the QGP produced in PbPb-collisions.

4 Summary

In this paper we have studied the collimation of energy inside jets produced in heavy ion collisions. From the production of one jet of opening angle Θ_0 initiated by a parton A, we considered the production of a subjet Θ ($\Theta < \Theta_0$) initiated by a parton B, where the definite fraction $x \sim 1$ of the jet energy is deposited. We made use of the DGLAP evolution equations at large x with medium-modified splitting functions, which accounts for medium-induced soft gluon radiation. Our study was extended to the distributions $D_A^B(x, Q^2)$ at $x \sim 1$ so as to provide their dependence on the nuclear parameter N_s and therefore, it allowed to determine the behavior of the collimation of the jet energy as a function of the same parameter.

Since these results are model-dependent, more efforts are required towards the construction of a more realistic QCD energy loss picture where the medium parameters and medium-induced soft gluon dynamics [13] could be both taken into account. However, as physically expected from the model, jets in the vacuum are more collimated than medium-modified jets and therefore, there is an evidence for the broadening of jets as gluon radiation is enhanced. Quark jets are more collimated than gluon jets in both the vacuum and the medium. Therefore, the collimation of energy treated in this paper and the collimation of average multiplicity treated in [9] are both good candidates to investigate the jet quenching.

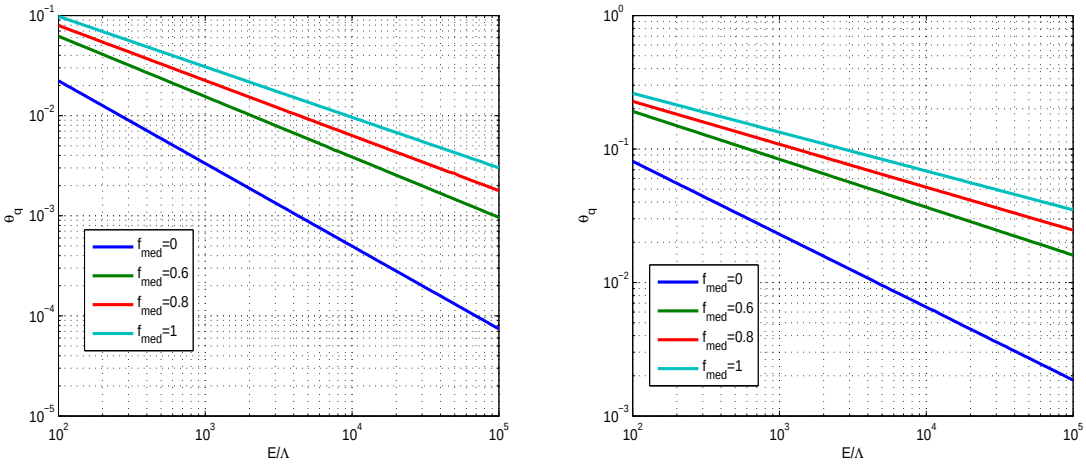

Figure 2: Collimation of energy inside a quark jet for $x = 0.5$ (left) and $x = 0.9$ (right).

Finally, in the forthcoming part of this work, we will compare the collimation with the Monte Carlo in-medium shower YaJEM event generator [14, 15] and with preliminary CMS data [16].

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