Radiative Energy Loss in Small and Large Systems

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Abstract

We use perturbative quantum chromodynamics to compute the corrections to the energy loss of a hard particle due to short separation distances between the creation of the particle and the in-medium scattering center that stimulates bremsstrahlung radiation. The result has several surprising features. The correction 1) does not go to zero for large path lengths; 2) breaks color triviality; 3) is formally zero in the large formation time approximation, but numerically dominates at large $\sim 100$ GeV parent parton energies out to long $\sim 3$ fm paths. Although motivated by the recent shocking evidence of collectivity in small systems, 1) and 3) mean our result has critically important implications for all energy loss model comparisons to data. In particular, 3) shows that a large fraction of the gluons radiated do not have a formation time that is large compared to the Debye screening length of the medium, violating a core approximation used in the derivations of all pQCD-based radiative energy loss formulae.

Keywords: pQCD, QGP, Energy Loss, Small Systems

1. Introduction

Recent startling results from the Relativistic Heavy Ion Collider (RHIC) and Large Hadron Collider (LHC) show that key signatures of quark-gluon plasma (QGP) formation are found in high-multiplicity p+p and p/d+A collision systems. In particular, collective behavior \cite{1, 2}, strangeness enhancement \cite{3, 4}, and quarkonium suppression \cite{5} appear to be sensitive only to the measured multiplicity of the collision, and not to the size of the nuclear fireball as naively implied by the type of the colliding particles.

Jet quenching is another key observable of QGP formation \cite{6, 7}, providing a unique femtoscope for probing the precise dynamics of the relevant degrees of freedom in this novel phase of nuclear matter. Energy loss models \cite{8–11} based on perturbative quantum chromodynamics (pQCD) have had enormous qualitative success in describing the momentum dependence and angular distribution of the suppression of high-momentum, $\sim 5$–150 GeV single particle pions \cite{12, 13} and charged hadrons \cite{14–16} from primordial hard light flavors and gluons and electrons \cite{17–19} as well as $D$ \cite{20} and non-prompt $J/\psi$ mesons \cite{21} from open heavy flavor decays at mid rapidity in A+A systems from $\sqrt{s} = 0.2$ ATeV to 2.76 ATeV.

Early experimental analysis shows a tantalizing correlation between centrality and suppression of jets in p/d+A collisions at RHIC and LHC \cite{22, 23}. These early results and sure-to-come future measurements call for quantitative theoretical predictions for jet tomography in small colliding systems.

There are two major complications to comparing theoretical predictions to experimental measurements in small colliding systems. First, phenomenologically, there is an inherent bias between rare high-multiplicity events and the rare collisions initially populated with one or more high transverse momentum (high-$p_T$) particles \cite{24, 25}. Second, theoretically, derivations of energy loss based on pQCD use simplifying assumptions \cite{26} that make them inapplicable to small systems.

The first complication makes it difficult to properly normalize the usual observable used in tomographic studies, the nuclear modification factor $R_{AB}$. $R_{AB}$ is the ratio of a spectrum in A+B collisions to the same particle spectrum in p+p collisions suitably normalized such that $R_{AB} = 1$ for particles unaffected by the presence of a QGP. Because of the aforementioned bias, properly normalizing $R_{AB}$ in high-multiplicity p+p and p/d+A events is problematic. One solution may be to divide the spectrum of interest by a known unaffected electroweak
spectrum with the same event selection criteria, forming a $\gamma_{AB}, W_{AB}$, or $Z_{AB}$.

The work of this letter was motivated by the second complication. In the usual DGLV opacity expansion [27, 28], the energy loss derivation assumes a large separation distance $\Delta z \equiv z_1 - z_0 \gg \lambda_{mfp} \gg 1/\mu$ between the initial production position $z_0$ of the hard parent parton and the position $z_1$ where it scatters off a QGP medium quasiparticle. The mean free path of the high-$p_T$ particle is $\lambda_{mfp} = 1/\rho \sigma \sim 1 - 2$ fm while the Debye mass in an infinite, static thermal QGP of temperature $T \sim 350$ MeV is $\mu = gT \sim 0.5$ GeV, as derived from thermal field theory [29]. In the collision of p+p or p/d+A, one expects a system of radius $\lesssim 2$ fm; for these collision systems then, most high-$p_T$ particles cannot have a large separation distance between production and scattering.

In this letter we derive the $N = 1$ in opacity generalization of DGLV for all separation distances $\Delta z$ in order to apply this radiative energy loss formalism to small systems.

We find two major unexpected results from our derivation: first, under the usual assumptions of eikonality, collinearity, and large gluon formation time common to all pQCD-based energy loss formalisms [26], the all-distance DGLV generalization is formally identical to the large separation distance DGLV expression; second, a numerical evaluation of the correction term, which is formally zero under the aforementioned assumptions, is actually dominant compared to the large separation distance DGLV expression at high (~ 100 GeV) parton energies.

We will show that the first curiosity is due to a highly non-trivial interference in the squared matrix element between non-zero small separation distance contributions to the first and second order Dyson terms in the $N = 1$ in opacity amplitude.

We will also show that the second surprise is due to the breakdown of the large formation time approximation in the DGLV formalism. We emphasize that while previous work has demonstrated the extreme sensitivity of all energy loss calculations to the collinear approximation [26, 30], and therefore the need to move beyond its use in all energy loss models, the sensitivity we find from the large formation time approximation is both new and different from the sensitivity to the collinear approximation.

As such all current jet quenching models that include radiative energy loss based on pQCD must individually assess their sensitivity to the large formation time approximation when making quantitative comparisons with data.

2. Setup

In this paper we use precisely the setup of the DGLV calculation [28], but the details of the current calculation can be found in [31]. For clarity, we treat the high-$p_T$ eikonal parton produced at an initial point $(t_0, z_0, x_0)$ inside a finite QGP, where we have used $p$ to mean transverse 2D vectors, $\mathbf{p} = (p_x, p_y)$ for 3D vectors and $p = (p^0, \mathbf{p}) = (p^0 + \mathbf{p}, p^0 - \mathbf{p}, \mathbf{p})$ for four vectors in Minkowski and light cone coordinates respectively. As in the DGLV calculation, we consider the target to be a Gyulassy-Wang Debye screened potential [32] with Fourier and color structure given by

$$V_n = V(q_n) e^{-i\mathbf{q}_n \cdot \mathbf{x}},$$

$$= 2\pi \delta(q^0) \nu(q_n, q_n^\perp) e^{-i\mathbf{q}_n^\perp \cdot \mathbf{x}} T_{n}(R) \otimes T_{\omega}(n).$$

The color exchanges are handled using the applicable $SU(N_c)$ generator $T_{\omega}(n)$ in the $d_{h}$ dimensional representation of the target or $T_{n}(R)$ in the $d_{h}$ dimensional representation of the high-$p_T$ parent parton.

In light cone coordinates the momenta of the emitted gluon, the final high-$p_T$ parton, and the exchanged medium Debye quasiparticle are

$$k = \left[xP_x, m^2 + k^2, k\right],$$

$$p = \left[(1-x)P_x, M^2 + k^2, (1-x)P^x, -k\right],$$

$$q = [q^0, q^\perp, \mathbf{q}],$$

where the initially produced high-$p_T$ particle of mass $M$ has large momentum $E^+ = p^0 = 2E$ and negligible other momentum components. Notice that we include the Ter-Mikayelian plasmon effect with an effective emitted gluon mass $m_g$ [28, 29]. See Fig. 1 for a visualization of these momenta.

A shorthand for energy ratios will prove useful notationally. Following [28] we define $\omega = x E^+/2 = x P^+/2$, $\omega_0 \equiv \mathbf{k}^2/2\omega$, $\omega_i \equiv (\mathbf{k} - \mathbf{q})^2/2\omega_i$, $\omega_{ii} \equiv (\mathbf{k} - \mathbf{q})^2/2\omega$, and $\tilde{\omega}_{ii} \equiv (m_{\mu}^2 + M^2 x^2)/2\omega$.

We will also make the following crucial assumptions: 1) the eikonal, or high energy, approximation, for which $E^+$ is the largest energy scale of the problem; 2) the soft (radiation) approximation $x \ll 1$; 3) collinearity, $k^+ \gg k^-$; 4) that the impact parameter varies over a large transverse area; and, most crucially for this letter, 5) the large formation time assumption $\omega_i \ll \mu_i$, where $\mu_i^2 \equiv \mathbf{q}_i \cdot \mathbf{q}_i$.

Note that the above approximations, in addition to allowing us to systematically drop terms that are small, permit us to 1) (eikonal) ignore the spin of the high-$p_T$
parton; 2) (soft) assume the source current for the parent parton varies slowly with momentum $J(p - q + k) \approx J(p + k) \approx J(p)$; 3) (collinearity) complete a separation of energy scales

$$E^+ \gg k^+ \gg k^- \equiv \omega_{ij} \sim \omega_{i(j, \ldots)} \gg \frac{(p + k)^2}{P^+}; \quad (3)$$

and 4) take the ensemble average over the phase factors, which become $\langle e^{-i\mathbf{q} \cdot \mathbf{q}'^b} \rangle = \frac{2\pi n}{A} \delta^2(\mathbf{q} - \mathbf{q}')$.

In the original DGLV calculations [28], the large formation time played only a minor role. However, when considering short separation distances between the scattering centers, the large formation time assumption naturally increases in importance.

With the above approximations, we reevaluated the 10 diagrams contributing to the $N = 1$ in opacity energy loss amplitude [28] without the additional simplification of the large separation distance $\Delta z \gg 1/\mu$ assumption.

3. Calculation and Results

In the original evaluation of the 10 diagrams contributing to the $N = 1$ in opacity energy loss derivation, the large separation distance approximation $\Delta z \gg 1/\mu$ allowed the neglect of terms proportional to $\exp(-\mu \Delta z)$. In our reevaluation of these 10 diagrams we retained all terms proportional to $\exp(-\mu \Delta z)$. However, we found an enormous simplification due to the large radiated gluon formation time approximation $\omega_{ij} \ll \mu$: all but 2 of the 10 diagrams’ 18 new small distance correction pole contributions are suppressed under the large formation time assumption. We show the two diagrams with non-zero contributions at the amplitude level $\mathcal{M}_{1,1,0}$ and $\mathcal{M}_{2,2,0}$ in the large formation time approximation in Fig. 1.

The full result for these two amplitudes under our approximation scheme is then

$$\mathcal{M}_{1,1,0} \approx -J(p)e^{i(p + k)\mathbf{z}_0} T_{\mathcal{A}_1} a_1 \int d^2 \mathbf{q}_1 (2\pi)^2 \langle 0, \mathbf{q}_1 \rangle e^{-i\mathbf{q}_1 \cdot \mathbf{b}_1} \times \frac{\mathbf{k} \cdot \mathbf{\epsilon}}{\mathbf{k}^2 + m_{g}^2 + x^2 M^2} \left[ e^{i(\mathbf{a}_1 \cdot \mathbf{z}_0)} (1 - \frac{1}{2} \mu e^{-i(\mathbf{z}_1 - \mathbf{z}_0)}) \right]$$

(4)

$$\mathcal{M}_{2,2,0} \approx J(p)e^{i(p + k)\mathbf{z}_0} T_{\mathcal{A}_2} a_2 a_1 \langle 0, \mathbf{q}_1 \rangle \langle 0, \mathbf{q}_2 \rangle \times igT_{\mathcal{A}_1} T_{\mathcal{A}_2} a_1 a_2 \langle 0, \mathbf{q}_1 \rangle \langle 0, \mathbf{q}_2 \rangle \times \frac{\mathbf{k} \cdot \mathbf{\epsilon}}{\mathbf{k}^2 + m_{g}^2 + x^2 M^2} \left[ e^{i(\mathbf{a}_1 \cdot \mathbf{z}_0)} + e^{-i(\mathbf{z}_2 - \mathbf{z}_0)} \frac{1 - \mu e^{-i(\mathbf{z}_1 - \mathbf{z}_0)}}{2(\mu_1 + \mu_2)} \right]$$

(5)

The double differential single inclusive gluon emission distribution is given by [28]

$$d^3 N^g_{1,0} d^3 N_J \approx d^3 \mathbf{P} d^3 \mathbf{K} \left( \frac{1}{dT} \text{Tr} \langle \mathcal{M}_1 \rangle^2 + \frac{2}{d_T} \mathcal{R} \text{Tr} \langle \mathcal{M}_1 \mathcal{M}_2 \rangle \right)$$

(6)

from which the energy loss, given by the energy-weighted integral over the gluon emission distribution $\Delta E = E \int dx dxdN_g/dx$, can be calculated from the amplitudes.

The main analytic result of our letter is then the $N = 1$ first order in opacity small distance generalization of the DGLV induced energy loss of a high-$p_T$ parton in a QGP:

Figure 1: $\mathcal{M}_{1,1,0}$ (top panel) and $\mathcal{M}_{2,2,0}$ (lower panel) are the only two diagrams that have non-zero short separation distance corrections in the large formation time limit. $\mathcal{M}_{2,2,0}$ is the double Born contact diagram, corresponding to the second term in the Dyson series in which two gluons are exchanged with the single scattering center.
\[ \Delta E^{(s)}_{\text{ind}} = \frac{C_R \alpha_s \Lambda E}{\pi} \int dx \int d^2 q_1 \frac{\mu^2}{(\mu^2 + q_1^2)^2} \int \frac{d^2 k}{\pi} \]

\[ \times \int d\Delta z \frac{\rho(\Delta z)}{E} \left[ -2 \frac{1 - \cos((\omega_0 + \omega_m)\Delta z))}{(k - q_1)^2 + m_\pi^2 + x^2 M^2} \right. \]

\[ \times \left. \left( \frac{k^2 + m_\pi^2 + x^2 M^2}{(k - q_1)^2 + m_\pi^2 + x^2 M^2} \right) \right] + \frac{1}{2} \frac{e^{-\mu_i \Delta z}}{k^2 + m_\pi^2 + x^2 M^2} \left( 1 - \frac{2C_R}{C_A} \right) \left( 1 - \cos((\omega_0 - \omega_m)\Delta z) \right) \]

\[ \times \left( \cos((\omega_0 - \omega_m)\Delta z) - \cos((\omega_0 - \omega_1)\Delta z) \right). \]

Figure 2: Fractional energy loss of charm and bottom quarks in a QGP with \( \mu = 0.5 \text{ GeV} \) and \( \lambda_{mfp} = 1 \text{ fm} \) for (a) fixed path length \( L = 4 \text{ fm} \), (b) fixed energy \( E = 10 \text{ GeV} \), and (c) fixed energy \( E = 100 \text{ GeV} \). In the figures, “DGLV” dashed curves are computed from the original \( N = 1 \) in opacity large separation distance DGLV formula while “DGLV + corr” solid lines are from our all separation distance generalization of the \( N = 1 \) DGLV result, Eq. 7.

The small separation distance correction shown in the last four lines of Eq. 7 has the properties we expect: 1) the correction goes to zero as the separation distance becomes large, \( \Delta z \to \infty \) (or, equivalently, as the Debye screening length goes to 0, \( \mu \to \infty \)) and 2) the correction term vanishes as the separation distance vanishes, \( \Delta z \to 0 \), due to the destructive interference of the LPM effect.

An immediate puzzle is the breaking of color triviality seen to all orders in opacity in the large separation distance approximation [27]; in the small separation distance correction, the color triviality breaking comes from the term proportional to \( 2C_R/C_A \). Unexpectedly, Euler’s formula for the exponential representation of the trigonometric functions reveals that the arguments of the cosines in Eq 7 are approximately zero in the large formation time limit (compared to the \( \mu_i \Delta z \)). Therefore, even though the amplitudes have a non-zero contribution in the large formation time limit, remarkably, the squared summed small separation distance correction term actually vanishes identically in the large formation time limit, yielding precisely Eq. 83 of [28].

Despite formally being 0 in the large formation time limit, we nevertheless numerically investigated the importance of the short separation distance correction term in Eq. 7 to produce Figs. 2a, 2b and 2c. The numerical results use the same values as [28]: \( \mu = 0.5 \text{ GeV} \), \( \lambda_{mfp} = 1 \text{ fm} \), \( C_R = 4/3 \), \( C_A = 3 \), \( \alpha_s = 0.3 \), \( m_{\text{charm}} \equiv m_c = 1.3 \text{ GeV} \) and \( m_{\text{bottom}} \equiv m_b = 4.75 \text{ GeV} \), and the QCD analogue of the Ter-Mikayelian plasmon effect was taken into account by setting \( m_{\text{gluon}} \equiv m_g = \mu/\sqrt{2} \).

As in [29], kinematic upper limits were used for the momentum integrals such that \( 0 \leq k \leq 2x(1-s)E \) and
0 ≤ q ≤ 5E/µ, due to finite kinematics. This choice of \(k_{\text{max}}\) guarantees that the final momentum of the parent parton is collinear to the initial momentum of the parent parton and that the momentum of the emitted gluon is collinear to the momentum of the parent parton. The fraction of momentum carried away by the radiated gluon, \(x\) was integrated over from 0 to 1. The distribution of scattering centers was assumed to be exponential in order to account for the rapidly expanding medium, \(\rho(z) = 2 \exp(-2\Delta z/L)/L\).

In Fig. 2a we show the fractional energy loss of charm and bottom quarks of varying energy propagating through a 4 fm long static QGP brick. Notice first that the small distance correction term is generally an energy gain due to the sign of the color triviality breaking term and, second, that the size of the correction relative to the long distance DGLV result grows with energy.

In Fig. 2b we plot the fractional energy loss of charm and bottom quarks of energy \(E = 10\) GeV for path lengths up to 5 fm. One sees that the small separation distance correction has a non-negligible effect even for large path lengths. Although initially unanticipated, the non-zero effect is due to the integration over all separation distances between the production point and the scattering position; even for large path lengths, some of the interaction distances between the parent parton and the target occur at separation distances that are small compared to the Debye screening scale. Additionally, the relative size of the small distance correction term and the leading DGLV result diminishes at fixed energy as the path length grows.

The main revelation of our numerical analysis is shown in Fig. 2c, which presents the fractional energy loss of 100 GeV charm and bottom quarks propagating up to 5 fm through a QGP. The small distance “correction” term dominates over the leading DGLV result for the first ~3 fm of the path.

The significant energy gain out to large ~3 fm paths seems difficult to reconcile with the measured experimental suppression of charged particles in central A+A collisions at LHC [14–16].

4. Asymptotic Analysis and Conclusions

In order to better understand the dominance of our computed short separation distance correction over the leading DGLV result, one may derive an asymptotic formula for the all separation distance expression.

Define \(\Delta E^{(1)}_{\text{ind}} \equiv \Delta E_{\text{DGLV}}^{(1)} + \Delta E_{\text{corr}}^{(1)}\), where \(\Delta E_{\text{ind}}^{(1)}\) is given by Eq. 7. Starting with the correction term \(\Delta E_{\text{corr}}^{(1)}\), we take all thermal and quark masses to zero, and analytically evaluate the integral over the scattering separation distance \(\Delta z\). Then, we remove the kinematic bound on the momentum kick from the medium \(q_{\text{max}} \to \infty\), shift the momentum integral, analytically evaluate the angular integrals in momentum space, and perform the integrals over \(k\) and \(q\). The result is

\[
\Delta E_{\text{corr}}^{(1)} = \frac{C_R \alpha_s}{2\pi} \frac{L}{\lambda_g} \left( -\frac{2C_R}{C_A} \frac{1}{2 + \mu L} \right) \times 2 \int_0^1 dx \log \left( \frac{L k_{\text{max}}}{2 + \mu L} \right).
\]

Taking for simplicity \(k_{\text{max}} = 2x E\) we find

\[
\Delta E_{\text{corr}}^{(1)} = \frac{C_R \alpha_s}{2\pi} \frac{L}{\lambda_g} \left( -\frac{2C_R}{C_A} \frac{1}{2 + \mu L} \right) \times \log \left( \frac{2EL}{2 + \mu L} \right).
\]

The equivalent asymptotic expression for the large separation distance leading massless DGLV result was derived in [27]. The result, with \(k_{\text{max}} \to \infty\), is

\[
\Delta E_{\text{DGLV}}^{(1)} \approx \frac{C_R \alpha_s}{4} \frac{L^2 \mu^2}{\lambda_g} \log \frac{E}{\mu}.
\]

There are several important features of Eqs. 9 and 10 to note. First, the terms not proportional to the color triviality breaking \(2C_R/C_A\) factor in Eq. 9 cancel at this level of approximation and the correction is purely an energy gain. Second, the correction term is log divergent in the upper bound of the perpendicular momentum of the emitted gluon \(k_{\text{max}}\) whereas the large separation distance DGLV term is finite for infinite \(k_{\text{max}}\). Third, the correction term is linear in \(L\), while the leading term is proportional to the usual \(L^2\). Fourth, the asymptotic correction term breaks color triviality as its magnitude is proportional to \(L/\lambda_g\), where \(\lambda_g\) is the mean free path of the parent parton (whether quark or gluon), instead of proportional to \(L/\lambda_g\), where \(\lambda_g\) is the mean free path for gluons.

Most important, destructive interference between the contributions to the large separation distance DGLV result lead to an energy loss that grows only logarithmically with energy \(E\). The short separation distance correction piece does not suffer from a similar interference and grows linearly with \(E\).

It is precisely this linear in \(E\) behavior compared to the \(\log(E)\) of the large separation distance DGLV term that leads to the correction term dominating over the leading term at higher energies.

The fact that the short separation distance “correction” term can dominate over the leading large separation distance DGLV result even out to path lengths
$L \sim 6\mu$ when relaxing the large formation time assumption (effects that should tend to zero under the large formation time assumption), suggests that the large formation time assumption is invalid in the DGLV approach.

Before further investigating the large formation time assumption, recall the issue of the collinear approximation in pQCD-based energy loss formulae [26]. It was shown in [30] that a significant fraction of the gluon radiation from $N = 1$ large separation distance DGLV is not emitted collinearly, despite the use of the collinear approximation is violated for much of the phase space of the emitted gluon. One may understand this breakdown of the collinear approximation in the leading DGLV formula by considering the required ordering $k^+ \gg k^-$. From Eq. 2, $k^+ = 2x\mu$ and $k^- = k^2/2xE$ we require

$$2x\mu \gg \frac{k^2}{2xE} \Rightarrow 2 \gg O(1),$$

where the $O(1)$ term ranges from $\sim 1/2$ up to 2. The lower limit of $1/2$ comes from considering the typical momentum fraction taken from the parent parton by the emitted gluon, $x_{np} \sim \mu/E$ [30]; the upper limit of 2 results from using $k_{max} = 2x\mu$. Thus the collinear approximation is violated for much of the phase space of the emitted gluon.

Similarly, the large formation time approximation requires that

$$\mu \gg \omega_i \sim \frac{1}{\tau} = \frac{k^2}{2xE} \Rightarrow \mu \times O(1),$$

where, again, the $O(1)$ term ranges from $\sim 1/2$ up to 2. It is important to note that the large formation time is a separate approximation from the collinear approximation; it is only when $k \sim \mu$ that the two approximations are equivalent.

Since all energy loss formalisms, GLV, BDMS-ZAW, AMY, and HT (see [7] and references therein) exploit the large formation time approximation, we are faced with a need to reassess the applicability of the large formation time assumption in any description of energy loss. While the influence of the assumption of collinearity was relatively easy to quantify across formalisms by simply varying the maximum allowable perpendicular momentum of the emitted gluon, estimating the importance of the large formation time approximation in the other formalisms may be a challenge. Similarly, deriving expressions in the other formalisms that do not rely on either the collinear or large formation time approximations might also be formidable. The lack of theoretical control over these assumptions calls into doubt the quantitative extraction of medium parameters through the use of jet quenching [10]. We leave addressing these issues for future work.

5. Acknowledgments

The authors wish to thank the SA-CERN Collaboration, the South African National Research Foundation (NRF), and the National Institute for Theoretical Physics (NITheP) for their generous support. The authors also wish to thank CERN for the hospitality extended to them during the completion of part of this work. Finally, the authors wish to thank Miklos Gyulassy, Ulrich Heinz, and Carlos Salgado for valuable discussions.

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