Collisional parton energy loss in a finite size QCD medium revisited: Off mass-shell effects

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We study the collisional energy loss mechanism for particles produced off-mass-shell in a finite size QCD medium. The off-mass-shell effects introduced are to consider particles produced in wave packets instead of plane waves and the length scale associated to an in-medium particles’ life-time. We show that these effects reduce the energy loss as compared to the case when the particles are described as freely propagating from the source. The reduction of the energy loss is stronger as this scale becomes of the order or smaller than the medium size. We discuss possible consequences of the result on the description of the energy loss process in the parton recombination scenario.

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I. INTRODUCTION

The problem of collisional parton energy loss in a QCD medium has been revived due to recent RHIC data on non-photonic single electrons [1] that are not well described within radiative energy loss calculations. Collisional energy loss has been a subject of research from long ago [2, 3, 4]. During the pioneering years, an important question was to understand how to handle the infrared singularities in perturbative calculations at finite temperature. The advent of resummation techniques clarified this point and allowed to reliably compute the energy loss of a heavy parton traversing an infinite medium to lowest order in perturbation theory [5, 6]. Short after, it was estimated that radiative energy-loss in a finite size medium was a more important mechanism to account for energy losses of energetic partons [7, 8, 9]. Nonetheless, even more recent studies [10, 11] suggested that for a range of parameters relevant to RHIC energies, radiative and collisional energy losses for heavy quarks are of the same order of magnitude. These last calculations where done for infinite QCD media. The outstanding question was whether collisional energy loss for finite size media was also significant.

In this context there where two results seemingly in contradiction [12, 13]. In Ref. [12] a semiclassical approach based on linear response theory computes the collisional energy loss by means of the work done by the response chromoelectric field on the color-charged heavy parton traversing the medium. The infinite medium limit of this description agrees with the collisional energy loss result at high temperature—up to color factors—obtained from a perturbative approach using HTL effective propagators [2]. The original claim that the finite size medium induced energy loss is strongly suppressed compared to the infinite medium case was later revised by properly subtracting the kinetic energy associated to producing the particle within the medium [10].

On the other hand, in Ref. [13] a lowest order perturbative calculation using HTL propagators finds that finite size effects on the collisional energy loss are not significantly suppressed as compared to the infinite medium case. The formulation of the problem is based on the assumption that the scattered particle originates within the medium but otherwise is produced on mass-shell.

However, when particles are emitted by sources lasting a finite amount of time they are not necessarily produced on their mass-shell since the source emits over a (wide) range of energies. A physical consequence is the possibility that the particle losses its identity within the medium. In the midsts of a high energy heavy-ion collision, such possibility can be realized in the recombination of a jet parton with the partons from the surrounding medium. Recall that during the propagation inside a deconfined QCD plasma, a fast parton can have not only induced gluon radiation but also induced absorption from thermal gluons. This process can fairly well be considered as parton recombination, which is one of the accepted mechanisms that are used to describe the distinct features of meson and baryon spectra that include a baryon to meson ratio larger than one for $p_t \gtrsim 2$ GeV in central Au + Au collisions at RHIC [14] and their different azimuthal anisotropies [15]. This point is addressed in the context of the modification of parton fragmentation functions induced by medium effects in Ref. [16]. When a parton recombines it certainly loses its original identity and the energy loss should stop being described in terms of parton degrees of freedom. Parton recombination from a jet with thermal partons to form intermediate $p_t$ hadrons is a viable scenario in the case of light flavors, given the features of the proton to pion ratio and even for s quarks, given the features of the $\Lambda$ to kaon ratio [17]. Other off mass-shell effects can be of relevance as well when studying if and how the virtuality of the propagating parton affects the in-medium splitting functions [18].

In this work we study one of such off mass-shell effects, namely a possible finite life-time of the scattered
partons–originating within and traversing the finite size medium– in the description of the collisional energy loss mechanism. We introduce the possibility that the scattered parton be described in terms of a propagator containing a parameter associated to the parton’s life time. We argue that for recombining partons, this picture could be used to consider collisional energy losses only up to times when these recombine with thermal partons from the medium from where the energy loss process should start being described in terms of hadronic degrees of freedom.

The work is organized as follows: In Sec. II we rewrite the expression for the collisional energy loss in a finite size QCD medium based on the formalism used in Ref. [13], allowing for particles being emitted off mass-shell and having a finite life-time. In Sec. III we give the numerical estimates for the collisional energy losses of heavy and light quarks using parameters relevant for RHIC energies. We use the cases studied in Ref. [13] as a base to compare our results to. We finally conclude and give an outlook in Sec. IV.
In order to explore a simple scenario, let us recall that

$$\lim_{\eta \to 0^+} \frac{\eta}{(p_0 - E_p)^2 + \eta^2} = (2E_p \pi) \delta(P^2 - M^2) \theta(p_0)$$  \hspace{1cm} (5)$$

where $E_p = +\sqrt{p^2 + M^2}$. To consider a finite width, we take $\eta$ finite and write

$$S(P) = \left( \frac{2M}{E_p} \right)^\eta \frac{\eta}{(p_0 - E_p)^2 + \eta^2} \Lambda(P).$$  \hspace{1cm} (6)$$

Upon the change of variable $y = x_1 - x$ and after integration over $d^4x_2$, $d^4y$, $d^4p$, $d^4q$ and $d^4x$ in Eq. (6) we get

$$iM = 2i \left( \frac{\eta}{E_p} \right) \int_{-\infty}^{\infty} \frac{dp_0}{(2\pi)i} e^{-i(p_0 - E' - \omega)L/2\nu}$$

$$\times \sin[(p_0 - E' - \omega)L/2\nu]$$

$$\times iM_0(p_0) \tilde{j}(P),$$  \hspace{1cm} (7)$$

where $\tilde{j}$ is the Fourier transform of $j(t, x)$ and

$$iM_0(p_0) = ig^2 D_{\alpha\beta}(K' - K)$$

$$\times \bar{u}(p', s') \gamma^\alpha u(p, s) \bar{u}(k', \lambda') \gamma^\beta u(k, \lambda),$$  \hspace{1cm} (8)$$

and we used that $\Lambda(P) u(p, s) = u(p, s)$. $M_0(p_0)$ is the matrix element describing the scattering process in infinite volume and for the following discussion, we have emphasized its dependence on $p_0$.

In order to perform the integral in Eq. (7), we write

$$\frac{1}{(p_0 - E_p)^2 + \eta^2} = \frac{1}{[(p_0 - E_P) + i\eta][(p_0 - E_p) - i\eta]},$$  \hspace{1cm} (9)$$

and for convergence, close the contour of integration on the lower $p_0$ complex half-plane, which picks the pole at $p_0 = E_p - i\eta$. This results in the expression for the matrix element

$$iM = i \left( \frac{1}{E_p} \right) e^{-\eta L/\nu} e^{-i(E_p - E' - \omega)L/2\nu}$$

$$\times \sin[(E_p - E' - \omega - i\eta)L/2\nu]$$

$$\times iM_0(E_p - i\eta) \tilde{j}(P).$$  \hspace{1cm} (10)$$

Notice that for a consistent description, we require the condition $\eta \ll E_p$ meaning that the central energy of the wave packet is much larger than it’s width. Therfore in Eq. (10) we can approximate

$$M_0(E_p - i\eta) \approx M_0(E_p).$$  \hspace{1cm} (11)$$

The square of the matrix element given in Eq. (10), averaged over the initial spin $s$ and summed over all other spins is therefore

$$\frac{1}{2} \sum_{s, s', \lambda, \lambda'} |M|^2 = e^{-\eta L/\nu} \left( \frac{1}{E_p' \pi} \right) |\tilde{j}(P)|^2$$

$$\times \sin[(\omega - v \cdot q + i\eta)L/2\nu]^2$$

$$\times \frac{1}{2} \sum_{s, s', \lambda, \lambda'} |M_0(E_p)|^2,$$  \hspace{1cm} (12)$$

where we use that for an energetic incoming fermion $E_p - E' \approx v \cdot q$. It is worth mentioning that if instead of using the propagator in Eq. (6), one uses the free Feynman propagator, the result for the square of the matrix element, averaged over the initial and summed over all other spins, yields the result found in Ref. [13].

Hereafter, we specialize to the description of the scattering of the fermion (quark) in the QCD plasma. The differential collisional energy loss $dE$ is related to the differential collision interaction rate $d\Gamma$ by $dE = \nu d\Gamma$ [9]. On the other hand, $d\Gamma$ is given by

$$d^3N d\Gamma = \frac{1}{2} \sum_{s, s', \lambda, \lambda'} |M|^2 \frac{d^3p'}{(2\pi)^32E'} \frac{d^3k'}{(2\pi)^32k'} \frac{d^3k}{(2\pi)^32k}$$

$$\times \sum_{\xi=g, q} n_{eq}(k)\{1 \pm n_{eq}(k')\}$$

$$\approx \frac{1}{2} \sum_{s, s', \lambda, \lambda'} |M|^2 \frac{d^3p'}{(2\pi)^32E'} \frac{d^3k}{(2\pi)^32k} \frac{d^3k'}{(2\pi)^32k'}$$

$$\times n_{eq}(k),$$  \hspace{1cm} (13)$$

where

$$d^3N = dR |\tilde{j}(P')|^2 |u(p', s')\bar{u}(p', s')|^2 \frac{d^3p'}{(2\pi)^32E'}$$

$$= dR |\tilde{j}(P')|^2 \frac{d^3p'}{(2\pi)^32E'},$$  \hspace{1cm} (14)$$

represents the number of (non-scattered) particles into the phase space volume in the interval $p'$ and $p' + d^3p'$ with $dR = 3$ for the SU(3) fundamental representation. Also, in Eq. (13) we have used that when describing the collisional energy loss, $n_{eq}(k'\{1 \pm n_{eq}(k')\}) = n_{eq}(k)$ since the term proportional to $n_{eq}(k)n_{eq}(k')$ is odd under the exchange of $k$ and $k'$ and integrates to zero [9], and have defined

$$n_{eq}(k) = \sum_{\xi=q, g} n_{eq}(k).$$  \hspace{1cm} (15)$$

For a source producing energetic particles with a large spread in momentum, we can take the approximation $|\tilde{j}(P')|^2 \approx |\tilde{j}(P)|^2$. Therefore, the collisional energy loss, considering the finite width of the scattered wave packet...
Notice that Eq. (16) is modified with respect to the corresponding expression in Ref. [13] by the $\eta$-dependent exponential factor and the $\eta$ dependence in the arguments of the sine function and in the energy denominator. When $\eta \to 0$, the corresponding expression for the energy loss in Ref. [13] is recovered. In order to find the explicit expression for Eq. (16), recall that the effective gluon propagator can be written as

$$D_{\mu\nu} = -P_{\mu\nu} \Delta_T - Q_{\mu\nu} \Delta_L$$

(17)

where in the HTL approximation, the effective transverse and longitudinal gluon propagators are given by

$$\Delta_T^{-1} = \omega^2 - q^2 - \frac{m_D^2}{2} - \frac{(\omega^2 - q^2)m_D^2}{2q^2}$$

$$\times \left( 1 + \frac{\omega}{2q} \ln \frac{\omega - q}{\omega + q} \right)$$

$$\Delta_L^{-1} = q^2 + m_D^2 \left( 1 + \frac{\omega}{2q} \ln \frac{\omega - q}{\omega + q} \right),$$

(18)

where $m_D^2 = g^2 T^2 (1 + N_f/6)$ is the square of the Debye mass and, working in Coulomb gauge, the only non-vanishing components of the transverse and longitudinal projectors are

$$P^{ij} = \delta^{ij} - \frac{q^i q^j}{q^2}$$

$$Q^{00} = 1.$$ (19)

Using Eqs. (17) - (19), the averaged over initial, summed over all other spins, matrix element squared describing the underlying scattering process in vacuum is given by

$$\frac{1}{2} \sum_{s, s', \lambda, \lambda'} |M_0|^2 = 16 g^4 E_p^2 \left\{ |\Delta_L(q)|^2 \left( kk' + k \cdot k' \right) + 2 \Re [\Delta_L(q) \Delta_L(q)^*] \times \left[ k \left( v \cdot k' - \frac{v \cdot q k \cdot k'}{q^2} \right) + k' \left( v \cdot k - \frac{v \cdot q k \cdot k}{q^2} \right) + |\Delta_T(q)|^2 \left[ 2 \left( v \cdot k - \frac{v \cdot q k \cdot q}{q^2} \right) \times \left( v \cdot k' - \frac{v \cdot q k' \cdot q}{q^2} \right) + (kk' - k \cdot k') \left( v^2 - \frac{v \cdot q v \cdot q}{q^2} \right) \right] \right\}.$$ (20)

Since for a non-expanding medium, the energy loss does not depend on the direction of $v$, we can simplify Eq. (16) by averaging over the direction of $v$ [5]. This is most con-
FIG. 3: (Color online) Fractional energy loss for (a) light, (b) charm and (c) bottom quarks as a function of the medium’s length for a fixed quark momentum $p = 10\text{ GeV}$. The uppermost curve in each case corresponds to the description without off mass-shell effects. This is compared to the case with off mass-shell effects for two values of $\eta = 40, 80\text{ MeV}$. In each case, the fractional energy loss decreases as the value of $\eta$ increases.

Conveniently performed by introducing the auxiliary functions

$$J_i = \int \frac{d\Omega}{4\pi} \frac{\sin[(\omega - \mathbf{v} \cdot \mathbf{q})L/2\nu]}{\omega - \mathbf{v} \cdot \mathbf{q} + i\eta}^2 (\omega - \mathbf{v} \cdot \mathbf{q})^i,$$  \hspace{1cm} (21)

$i = 1, 2, 3$, in terms of which, the average over the different powers of $\mathbf{v}$ appearing in Eq. (20) can be expressed. The functions in Eq. (21) are explicitly given in the appendix. After averaging over the directions of $\mathbf{v}$, the expression for the energy loss can be written as

$$\Delta E = \frac{C_R g^4}{2\pi^4} e^{-\eta L/\nu} \int_0^{\infty} n_{eq}(k) dk \times \left( \int_0^k q dq \int_{-q}^q \omega d\omega + \int_{q_{\text{max}}}^k q dq \int_{q-2k}^q \omega d\omega \right) \times \left[ (\Delta L(q))^2 \frac{(2k + \omega^2 - q^2)^2}{2} J_0 ight. + \left. |\Delta T|^2 \frac{q^2 - \omega^2}{4q^4} [(2k + \omega^2 + q^2) J_0 + 2\omega J_1 - J_2] \right],$$  \hspace{1cm} (22)

where the limits of integration over $\omega$ and $q$ take into account that for the considered scattering amplitude, the transferred four-momentum is space-like and

$$q_{\text{max}} = \frac{2k(1 + k/E_p)}{1 - v + 2k/E_p},$$  \hspace{1cm} (23)

is obtained from the approximation that the maximum energy transferred occurs for backward scattering \[21\].

III. NUMERICAL RESULTS

In order to present the quantitative behavior for the energy loss, we take standard values for the parameters involved. We give examples of the effect for light as well as for heavy flavors both, to study the mass effect and to directly compare to the findings of Ref. [13]. The plasma temperature is taken as $T = 0.225\text{ GeV}$, the effective number of flavors $N_f = 2.5$, the strength of the coupling constant $\alpha = g^2/4\pi = 0.3$ and the Debye mass $m_D = 0.5\text{ GeV}$. The bottom quark mass is taken as 4.5 GeV whereas the charm quark mass is taken as 1.2 GeV. We take the mass of the light quarks as 0.2 GeV.

Figure 2 shows the fractional energy loss $\Delta E/E$ for light, charm and bottom quarks in a finite size medium with $L = 5\text{ fm}$, for the cases with and without off mass-shell effects. For the curves describing the off mass-shell effects, we consider two values, $\eta = 40\text{ MeV}$ and $\eta = 80\text{ MeV}$. Notice that in all cases, a finite value of $\eta$ produces the energy loss to decrease as compared to the case $\eta = 0$. The decrease is more important for larger values of $\eta$.

Figure 3 shows the fractional energy loss for light, charm and bottom quarks as a function of the medium’s length $L$ for a fixed quark momentum $p = 10\text{ GeV}$, comparing also the cases with and without off mass-shell effects. For the curves describing the off mass-shell effects, we consider once more the two values, $\eta = 40\text{ MeV}$ and $\eta = 80\text{ MeV}$. Notice that a finite value of $\eta$ causes the fractional energy loss to asymptotically reach a maximum value as the medium’s size increases. This is in
sharp contrast with the case where no off mass-shell effects are considered where for large $L$ the energy loss increases linearly.

Figure 4 shows the behavior of the fractional energy loss as a function of $\eta$ also for light, charm and bottom quarks for a fixed value of the quark momentum $p = 10$ GeV and a fixed value of the medium’s size $L = 5$ fm. The fractional energy loss decreases with increasing $\eta$ and the decrease is similar in shape, regardless of the quark mass.

IV. DISCUSSION AND CONCLUSIONS

In this work we have studied the off mass-shell effects on the collisional energy loss of particles, produced and scattered within a finite size QCD medium, associated to the introduction of a finite-width wave packet and therefore a finite particle’s life-time. We have shown that this effect decreases the energy loss as compared to the case when these particles are produced on mass-shell and therefore live longer than the medium, fragmenting outside it. We have argued that this picture should be applied in particular to energetic partons that recombine with thermal partons and thus hadronize within the medium.

Recall that the length scales playing a role for the energy loss mechanisms in a finite size, thermal, non-expanding medium are the medium’s size $L$, the average distance between collisions $d \sim 1/T$, the Debye radius $\delta_D$, the mean free path $\delta \sim 1/g^2 T$ and the particle’s formation time $t_f \sim 1/E_p$. When considering the in-medium particle’s lifetime, the length scale $\eta^{-1}$ is introduced. For the description of the scattering process in terms of a perturbative picture, it is required that the hierarchy of scales

$$1/T \ll r_D \ll \delta \quad (24)$$

be satisfied. The requirement that the scattering particle can be described in terms of an oscillating mode means that this mode is not damped too fast, which in turn translates into the condition

$$1/E_p \ll 1/\eta, \quad (25)$$

and can be thought of as the condition that the time the medium takes to produce the particle is much shorter than the in-medium particle’s lifetime. For very energetic partons, it is safe to assume that $t_f \sim 1/E_p$ is the smallest of all length scales. Moreover, for media sizes of the order expected to be produced in relativistic heavy-ion collisions, it is also safe to assume that $L$ is larger than the mean free path and therefore the hierarchy of scales

$$1/E_p \ll 1/T \ll r_D \ll \delta \lesssim L, \quad (26)$$

follows. In fact, the results in Refs. [13, 14] can be viewed as meaning that as long as the medium size is larger than the Debye radius, the collisional energy loss for on mass-shell particles in a finite size medium is not suppressed as compared to the infinite medium case.

The situation changes when introducing the in-medium particles’ life-time. The results of this work show that when $\eta^{-1} \lesssim L$ the effects are strong. They cease to matter for $\eta \to 0$.

It should be emphasized that in the context of this work, the term loss of identity does not refer to a parton change of color or phase, which are accounted for already in the description of the underlying QCD scattering process, neither to a change of flavor which would be mediated by the weak interaction and thus irrelevant for the time scales involved during the QCD plasma phase. This term is rather related to the onset of a hadronization mechanism that happens during the interaction of a fast (hard) parton with soft ones from the medium in such a way as to produce a hadron by means of recombination. This process has been referred to as shower and thermal parton recombination, in Ref. [22]. In this way, the loss of identity is related to the fact that when this kind of partons form a hadron, the energy loss should stop being described in terms of parton degrees of freedom and start being described in terms of hadron degrees of freedom, thereby effectively producing the parton to disappear from the description.

Typical time scales involved in hard-soft parton recombination are of order $\tau_{\text{recomb}} \sim 1.5$ fm, despite the low momentum transfers involved, since, as argued for instance in Ref. [22], this recombination needs not be local and it can be mediated by a QCD string. Such time scale is well within the life-time of the QCD medium, for central
collisions and the largest nuclei, where it is estimated to be of order 5 fm.

When the medium length shortens and falls below \( \tau_{\text{recomb}} \), hadronization is more likely to happen outside the medium. This means in particular that for peripheral collisions or collisions of smaller systems, the energy loss description in terms of partonic degrees of freedom is appropriate. This is accounted for in our description when we take \( \eta^{-1} > L \), for which the energy loss with and without the use of the parameter \( \eta \) coincide. Indeed, a change in the energy loss of intermediate momentum hadrons should exist as a function of centrality and system size. In order to quantify such change it is necessary to quantitatively estimate the energy loss of a hadron within a QCD medium. This calculation is for the moment outside the scope of the present work.

In the context of recombination, the use of the parameter \( \eta \) should be introduced into a statistical scenario that also incorporates the evolution of the colliding system with energy density, as well as using a realistic geometry, including adequate probability profiles to produce jets [21] and the effects of an expanding medium. All this is for future.

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V. APPENDIX

The functions defined in Eq. (21) are explicitly given by
\[ J_0 = \int \frac{d\Omega}{4\pi} \left| \frac{\sin[(\omega - v \cdot q + i\eta)L/2v]}{\omega - v \cdot q + i\eta} \right|^2 \]
\[ = -\frac{1}{8qv\eta} \left[ i \cosh(\eta L/v) \left( \sum_{l,l' = \pm 1} \text{sgn}(l)\text{sgn}(l')Ci[(\omega - lqv - l'i\eta)L/v] \right) \right. \]
\[ + 2i \left. \left[ \arctan \left( \frac{qv + \omega}{\eta} \right) + \arctan \left( \frac{qv - \omega}{\eta} \right) \right] \right] \]
\[ + \sinh(\eta L/v) \left( \sum_{l,l' = \pm 1} \text{sgn}(l)Si[(\omega - lqv - l'i\eta)L/v] \right) \]
\[ J_1 = \int \frac{d\Omega}{4\pi} \left| \frac{\sin[(\omega - v \cdot q + i\eta)L/2v]}{\omega - v \cdot q + i\eta} \right|^2 (\omega - v \cdot q) \]
\[ = \frac{1}{8qv} \left[ \cosh(\eta L/v) \left( \sum_{l,l' = \pm 1} \text{sgn}(l)Ci[(\omega - lqv - l'i\eta)L/v] \right) \right. \]
\[ + \ln \left( \frac{(\omega + qv)^2 + \eta^2}{(\omega - qv)^2 + \eta^2} \right) \]
\[ - i \sinh(\eta L/v) \left( \sum_{l,l' = \pm 1} \text{sgn}(l)\text{sgn}(l')Si[(\omega - lqv - l'i\eta)L/v] \right) \],
\[ J_2 = \int \frac{d\Omega}{4\pi} \left| \frac{\sin[(\omega - v \cdot q + i\eta)L/2v]}{\omega - v \cdot q + i\eta} \right|^2 (\omega - v \cdot q)^2 \]
\[ = \frac{1}{8qv} \left[ \eta \cosh(\eta L/v) \left( i \sum_{l,l' = \pm 1} \text{sgn}(l)\text{sgn}(l')Ci[(\omega - lqv - l'i\eta)L/v] \right) \right. \]
\[ - 2 \left. \left[ \arctan \left( \frac{qv + \omega}{\eta} \right) + \arctan \left( \frac{qv - \omega}{\eta} \right) \right] + \frac{4qv}{\eta} \right] \]
\[ + \eta \sinh(\eta L/v) \left( \sum_{l,l' = \pm 1} \text{sgn}(l)Si[(\omega - lqv - l'i\eta)L/v] \right) \]
\[ - \left( \frac{4v}{L} \right) \cos(L\omega/v) \sin(Lq) \right] , \quad (27) \]

where \text{sgn} is the sign function and \(Ci\) and \(Si\) are the cosine and sine integrals, respectively. Despite their appearance, the above functions are all real for real values of \(\omega, q, v, L\) and \(\eta\).

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