PROCEEDINGS OF SCIENCE

High-p_T parton propagation and quenching in the medium

Konrad Tywoniuk*
Theoretical Physics Department, CERN, 1211 Geneva 23, Switzerland
E-mail: konrad.tywoniuk@cern.ch

Jet quenching stands out as one of the most striking phenomena measured in heavy-ion collisions. After being discovered at RHIC as a suppression of high-p_T hadrons, with the development of reconstruction techniques capable of dealing with large backgrounds, it has also been confirmed in a wide range of measurements involving fully reconstructed jets. The observed suppression patterns are generally understood as arising from interactions with a color deconfined medium, corroborating the role of jets as probes of the hot and dense matter created in ultrarelativistic heavy-ion collisions. In this proceeding, we report on recent theoretical developments related to how jet fragmentation inside the medium is affected by energy-loss processes. In particular, it is shown how a novel Sudakov suppression factor affects the single-inclusive jet spectrum at high-p_T. These developments provide first steps toward a unified theoretical understanding of the quenching of hadrons and jets, accompanied by substructure modifications.

12th International Workshop on High-pT Physics in the RHIC/LHC Era
2-5 October, 2017
University of Bergen, Bergen, Norway

*Speaker.

https://pos.sissa.it/

PoS(High-pT2017)011

1. Introduction

The strong suppression of high-p_T particles and jets, including heavy quarks, in heavy-ion collisions stand out as one of the hallmark measures of final-state interactions in a colour deconfined medium. While the suppression persists up to very large momenta for jets, the suppression of single-inclusive hadrons tend to gradually subside in the same variable, a feature that also holds true for the modification of intra-jet properties. This calls for a unified and comprehensive theoretical understanding of jet modifications across many observables over a large range in jet scales, mainly p_T and the reconstruction angle R.

It has already been pointed out in Monte-Carlo studies that fluctuations related to the jet fragmentation, or substructure, are extremely important for understanding experimental data, see e.g. [1, 2], but until recently a first-principle understanding of these higher-order corrections was lacking. Quenching in the medium has traditionally been applied to a single partons [3, 4], giving limited theoretical guidance to implementations of jet quenching on the level of multi-parton fragmentation. It was, however, pointed out that breaking of colour coherence plays an important role for the available phase space of emissions [5, 6, 7, 8], and recently it was also shown how this affects energy loss processes [9]. In these proceedings, we summarize the main ingredients that enter the calculation of single- and multi-parton quenching in the medium. We also show how these lead to large corrections on the level of single-inclusive spectra, and finally outline the alleys for future applications.

2. Jet quenching

Single-parton quenching An energetic, single colour charge propagating through the medium will lose energy as a result of multiple scatterings with medium constituents. In order to simplify the discussion below, let us for the moment consider a large, static medium of size L. The main source of energy degradation of the probe are inelastic processes, or so-called medium-induced emissions. In the presence of multiple soft scattering in the medium, the rate of induced soft gluon emission read [10, 11],

$$\frac{\mathrm{d}I}{\mathrm{d}\omega\mathrm{d}t} = \bar{\alpha}\sqrt{\frac{\hat{q}}{\omega^3}}\,, \tag{2.1}$$

where $\bar{\alpha} \equiv \alpha_s N_c/\pi$. Here, $\hat{q}$ is the celebrated jet transport coefficient. This spectrum is valid for $\omega \ll \omega_c$ with a cut-off at the energy $\omega_c \sim \hat{q}L^2$, which corresponds to the situation when the formation time of the gluons becomes of the order of the length of the medium, $t_{\mathrm{br}} = \sqrt{\hat{q}/\omega} \sim L$. Hence, on the level of one gluon emission, the probability of radiating energy ε off an energetic projectile becomes

$$P_1(\varepsilon) \simeq \delta(\varepsilon)\left(1 - \int_0^\infty \frac{\mathrm{d}I}{\mathrm{d}\omega}\right) - \frac{\mathrm{d}I}{\mathrm{d}\varepsilon}\,, \tag{2.2}$$

at leading order in α_s. Notice, however, that the multiplicity of gluons above a certain energy, defined as

$$N(\omega) = \int_\omega^\infty \mathrm{d}\omega' \int_0^L \mathrm{d}t\, \frac{\mathrm{d}I}{\mathrm{d}\omega'\mathrm{d}t}\,, \tag{2.3}$$

becomes large $N \gtrsim \mathcal{O}(1)$ for $\omega \lesssim \omega_s \sim \bar{\alpha}^2 \hat{q} L^2 \ll \omega_c$. In this regime, multiple induced emissions ought to be resummed. Given that their formation time is short $t_{\rm br}(\omega_s) \sim \bar{\alpha} L \ll L$, we can neglect interference effects between subsequent emissions.[1] Besides, since their emission angles, given by $\theta_{\rm br}(\omega) = (\hat{q}/\omega^3)^{1/4}$ are parametrically large, $\theta_{\rm br}(\omega_s) \sim \theta_c/\bar{\alpha}^{3/2} > R$, where $\theta_c \equiv \theta_{\rm br}(\omega_c) \sim (\hat{q}L^3)^{-1/2}$ and R is the jet cone size, they are mainly responsible for transporting energy to large angles. Keeping these simplifications in mind, the resummation of multiple induced emissions is then realised via the rate equation,

$$\frac{\partial}{\partial t} P_1(\varepsilon, t) = \int_0^\infty {\rm d}\omega \left[\frac{{\rm d}I}{{\rm d}\omega\, {\rm d}t} - \delta(\omega) \int_0^\infty {\rm d}\omega' \frac{{\rm d}I}{{\rm d}\omega'\, {\rm d}t} \right] P_1(\varepsilon - \omega, t) \,. \tag{2.4}$$

With Eq. (2.1), the solution of Eq. (2.4) is found to be

$$P_1(\varepsilon, t) = \sqrt{\frac{\omega_s}{\varepsilon^3}} {\rm e}^{-\frac{\pi \omega_s}{\varepsilon}} \,. \tag{2.5}$$

Comparing Eq. (2.2) to the resummed result in Eq. (2.4) reveals that the *typical* energy lost is strongly correlated with the scale ω_s. The simplified description we have outlined so far can be improved in many ways, for instance by imposing kinematical restrictions and by also keeping track of secondary branchings that contribute to the energy flow out of the jet cone [12, 13].

Multi-parton quenching Hard QCD splittings early in the medium create pairs of colour correlated partons that propagate through almost the entirety of the medium length. The presence of colour correlations alter the emission spectrum of partons even in vacuum, giving rise to the well-established angular ordering of subsequent gluon emissions. In the medium, this is modified because of colour exchanges with the medium that lead to colour decoherence, see e.g. [6, 8, 7]. In this section we will discuss what role decoherence effects play when considering the energy-loss off a composite system.

The quenching of a pair of partons, in the large-N_c limit, can be understood as the combined effect of the quenching of the total charge, that is related to the color charge of the parent parton, and the additional quenching related to the additional color charge generated in the splitting [9]. The latter process is however delayed due to the fact that medium only has a finite resolution, as we will describe in more detail below. Hence, in the large-N_c limit, we write

$$P_2^R(\varepsilon) = \int_0^\infty {\rm d}\varepsilon_1 \int_0^\infty {\rm d}\varepsilon_2\, P_1^R(\varepsilon_1) P_{\rm sing}(\varepsilon_2) \delta(\varepsilon - \varepsilon_1 - \varepsilon_2) \,, \tag{2.6}$$

for a generic parton in color representation R. The first term, carrying the information about this index, describes the quenching of the total colour charge. The second describes the additional energy loss of a colour-singlet dipole appearing in the splitting in the large-N_c limit.

In analogy to the steps leading up to Eq. (2.4), we resum multiple induced soft gluons by the rate equation,

$$\frac{\partial}{\partial t} P_{\rm sing}(\varepsilon, t) = \int_0^\infty {\rm d}\omega\, \Gamma_{\rm dir}(\omega, t) P_{\rm sing}(\varepsilon - \omega, t) + \int_0^\infty {\rm d}\omega\, S_2(t) \Gamma_{\rm int}(\omega, t) \delta(\varepsilon - \omega) \,, \tag{2.7}$$

[1] Note that this does not hold for hard, vacuum-like splittings where interference effects are important, see below.

where $\Gamma(\omega,t)$ is a shorthand for the regularised rates $\Gamma_i(\omega,t) \equiv \mathrm{d}I_i/(\mathrm{d}\omega\mathrm{d}t) - \delta(\omega)\int_0^\infty \mathrm{d}\omega'\, \mathrm{d}I_i/(\mathrm{d}\omega'\mathrm{d}t)$ for $i = \{\mathrm{dir}, \mathrm{int}\}$ for direct emissions and interferences. This formula contains new elements compared to the single-charge energy loss resummation in Eq. (2.4) that are related to interferences between the emitters.

The interference term involves two components. The first is a dipole suppression factor describing the survival probability of color coherence at a given time in course of the dipole propagation. It is called the decoherence parameter, and reads

$$S_2(t) = \exp\left\{-\frac{1}{12}\hat{q}\theta^2 t^3\right\}, \tag{2.8}$$

which gives rise to a characteristic time-scale for decoherence. This time-scale can easily be estimated by comparing the medium resolution scale due to multiple scattering $\lambda_\perp \sim (\hat{q}t)^{-1/2}$ with the size of the fluctuation $x_\perp \sim \theta t$ where θ is the dipole angle. The two scales become comparable at $t_\mathrm{d} \sim (\hat{q}\theta^2)^{1/3}$, which is the so-called decoherence time. We can immediately identify this scale appearing in Eq. (2.8). Finally, without going into the details, the interference spectrum can be written approximately as

$$\frac{\mathrm{d}I_\mathrm{int}}{\mathrm{d}\omega\mathrm{d}t} \approx -\bar{\alpha}\sqrt{\frac{\hat{q}}{\omega^3}}\,\Theta\big(\omega < (\theta^2 t)^{-1}\big). \tag{2.9}$$

We note first of all that it is negative and will therefore cancel the direct emissions in part of the radiation phase space. This phase space is first of all limited to large angles, as indicated by the cut-off factor in Eq. (2.9) which implies angular ordering. This becomes clear by comparing the size of the dipole $x_\perp \sim \theta t$ to the typical size of fluctuations $\lambda_\perp \sim k_\perp^{-1}$ that appear at the formation time $t_\mathrm{f} \sim \omega/k_\perp^2$, and realising that the colour-singlet dipole can only radiate quanta that resolve the individual charges, $\lambda_\perp < x_\perp$. Second of all, due to the presence of the decoherence parameter, Eq. (2.8), these contributions survive in any case only as long as the dipole pair remains colour-correlated.

The emerging picture is therefore very simple and can be summarised as follows. Initially, the colour-singlet dipole remains close together in transverse space so that the medium interactions cannot resolve its inner structure but only its (zero) total charge. Consequently, there is no induced radiation of this object. As the pair grows and ultimately becomes resolved by the medium, both constituents lose energy independently.

Before leaving this section, let us estimate the phase space for hard emissions that would give rise to the additional quenching implemented in Eq. (2.7). We demand that such emissions are formed early in the medium, yet well before the medium has had the time to resolve their structure. To be sensitive to the quenching of new dipoles they have to be resolved, and therefore $t_\mathrm{f} < t_\mathrm{d} < L$. The phase space becomes

$$\Pi^{(\mathrm{coh})} = \bar{\alpha}\int_{t_\mathrm{f}<t_\mathrm{d}<L} \frac{\mathrm{d}\theta}{\theta}\frac{\mathrm{d}\omega}{\omega} = \bar{\alpha}\ln\frac{R}{\theta_c}\left(\ln\frac{p_T}{\omega_c} + \frac{2}{3}\ln\frac{R}{\theta_c}\right), \tag{2.10}$$

in the leading-logarithmic approximation (LLA). In particular, the condition $t_\mathrm{d} < L$ immediately implies that $\theta > \theta_c$, thus cutting away the collinear divergence. Therefore, this phase space is single-logarithmic in the jet energy [14] which signals that these additional contributions should be resummed for high-p_T jets. It is worth pointing out what happens if we do not consider coherence effects in the medium, in which case the dipole constituents always lose energy independently. The resulting phase space is also enhanced, since

$$\Pi^{(\mathrm{incoh})} = \bar{\alpha} \int_{t_\mathrm{f}<L} \frac{\mathrm{d}\theta}{\theta}\frac{\mathrm{d}\omega}{\omega} = \frac{\bar{\alpha}}{2} \ln^2 p_T R^2 L \,, \tag{2.11}$$

is double-logarithmic in the jet scale. We therefore anticipate that the implementation of coherence effects for calculating quenching of multi-parton systems, as is modelled in Monte-Carlo generators, will play a significant role.

3. Sudakov suppression of jets

Let us now consider the single-inclusive jet spectrum in the presence of a medium,

$$\frac{\mathrm{d}\sigma^{\mathrm{med}}}{\mathrm{d}p_T^2 \mathrm{d}y} = \int_0^\infty \mathrm{d}\varepsilon \, P(\varepsilon) \frac{\mathrm{d}\sigma^{\mathrm{vac}}(p_T+\varepsilon)}{\mathrm{d}p_T'^2 \mathrm{d}y} \,, \tag{3.1}$$

where $\mathrm{d}\sigma^{\mathrm{vac}}$ is the vacuum spectrum while $P(\varepsilon)$ is the probability distribution related to emitting energy ε out of the jet cone, discussed in some detail above. In addition to its sensitivity to the jet quenching parameter $\hat{q}$ and the medum size L, it also depends on the jet p_T and cone size R. For the time being we will be interested in the nuclear modification factor, or quenching factor,

$$R_{\mathrm{jet}} = \left(\frac{\mathrm{d}\sigma^{\mathrm{med}}}{\mathrm{d}p_T^2 \mathrm{d}y} \right) \Bigg/ \left(\frac{\mathrm{d}\sigma^{\mathrm{vac}}}{\mathrm{d}p_T^2 \mathrm{d}y} \right) \,, \tag{3.2}$$

describing the deviation from the vacuum baseline. We note, that by approximating the spectrum as $\mathrm{d}\sigma^{\mathrm{vac}}(p_T+\varepsilon) \simeq \mathrm{d}\sigma^{\mathrm{vac}}(p_T)\mathrm{e}^{-n\varepsilon/p_T}$, the nuclear modification factor is simply related to the appropriate moment of the Laplace transform of the quenching weight, $R_{\mathrm{jet}} = Q(p_T)$, with $Q(p_T) \equiv \tilde{P}(p_T/n)$, due to the convolution with the steeply falling spectrum.

The quenching factor permits an expansion in terms of the strong coupling constant, where the zeroth order contribution corresponds to the quenching of a jet consisting of a single constituent. Remarkably, in the large-N_c approximation, all higher order terms in the quenching weight depend on the quenching of the total colour charge in a trivial way, see Eq. (2.7). We can therefore show that

$$R_{\mathrm{jet}} = Q_{\mathrm{tot}}(p_T) \times C(p_T, R) \,, \tag{3.3}$$

whose simplicity owes to the fact that all convolutions become products in Laplace space. The quenching factor of the total color charge is simply given as $Q_{\mathrm{tot}}(p_T) = \tilde{P}_1^R(p_T/n)$. The second component on the right hand side of Eq. (3.3), $C(p_T,R) \equiv Q(p_T)/Q_{\mathrm{tot}}(p_T)$, is a new Sudakov suppression factor, called the collimator function, that resums quenching of higher-order jet fluctuations to the single-inclusive jet spectrum. At first non-trivial order, $C(p_T,R) = 1 + C^{(1)}(p_T,R) + \mathcal{O}(\alpha_s^2)$, in LLA it is found to be

$$C^{(1)}(p_T,R) = \bar{\alpha} \int_{t_\mathrm{f}<t_\mathrm{d}<L} \frac{\mathrm{d}\theta}{\theta}\frac{\mathrm{d}\omega}{\omega} \left[Q(p_T)^2 - 1 \right] . \tag{3.4}$$

As seen in Eq. (2.10), this gives rise to a logarithmic enhancement, which signals that higher order contributions should be added as well. We refer to [14] for a comprehensive discussion of the resummation procedure.

4. Outlook

In these proceedings, we have reported on novel theoretical developments related how multi-parton, with internal colour correlations, systems are being affected by quenching effects in the medium. This involves a better control of the phase space restrictions for in-medium radiation, that are shown to involve the jet scale, see Eq. (2.10), and have to be resummed. One of the consequences that has recently been explored is related to higher-order corrections of single-inclusive spectra and the related nuclear modification factor [14] in Eq. (3.3). Similar ideas are also being developed for the in-medium fragmentation function [15]. Ultimately, these developments will aid in attaining a better theoretical control for Monte-Carlo implementations and will lead to a better grip on the properties of the dense QCD medium created in heavy-ion collisions.

References

[1] J. G. Milhano and K. C. Zapp, Eur. Phys. J. C **76** (2016) no.5, 288 [arXiv:1512.08107 [hep-ph]].

[2] J. Casalderrey-Solana, D. Gulhan, G. Milhano, D. Pablos and K. Rajagopal, JHEP **1703** (2017) 135 [arXiv:1609.05842 [hep-ph]].

[3] R. Baier, Y. L. Dokshitzer, A. H. Mueller and D. Schiff, JHEP **0109** (2001) 033 [hep-ph/0106347].

[4] C. A. Salgado and U. A. Wiedemann, Phys. Rev. D **68** (2003) 014008 [hep-ph/0302184].

[5] Y. Mehtar-Tani, C. A. Salgado and K. Tywoniuk, Phys. Rev. Lett. **106** (2011) 122002 [arXiv:1009.2965 [hep-ph]].

[6] Y. Mehtar-Tani, C. A. Salgado and K. Tywoniuk, Phys. Lett. B **707** (2012) 156 [arXiv:1102.4317 [hep-ph]].

[7] J. Casalderrey-Solana and E. Iancu, JHEP **1108** (2011) 015 [arXiv:1105.1760 [hep-ph]].

[8] Y. Mehtar-Tani, C. A. Salgado and K. Tywoniuk, JHEP **1210** (2012) 197 [arXiv:1205.5739 [hep-ph]].

[9] Y. Mehtar-Tani and K. Tywoniuk, arXiv:1706.06047 [hep-ph].

[10] R. Baier, Y. L. Dokshitzer, A. H. Mueller, S. Peigne and D. Schiff, Nucl. Phys. B **484** (1997) 265 [hep-ph/9608322].

[11] B. G. Zakharov, JETP Lett. **65** (1997) 615 [hep-ph/9704255].

[12] S. Caron-Huot and C. Gale, Phys. Rev. C **82** (2010) 064902 [arXiv:1006.2379 [hep-ph]].

[13] J. P. Blaizot, E. Iancu and Y. Mehtar-Tani, Phys. Rev. Lett. **111** (2013) 052001 [arXiv:1301.6102 [hep-ph]].

[14] Y. Mehtar-Tani and K. Tywoniuk, arXiv:1707.07361 [hep-ph].

[15] P. Caucal, E. Iancu, A. H. Mueller and G. Soyez, arXiv:1801.09703 [hep-ph].