**Parton energy loss in the mini quark-gluon plasma and jet quenching in proton-proton collisions**

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**Abstract.** We evaluate the medium suppression of light hadron spectra in pp collisions at RHIC and LHC energies in the scenario with formation of a mini quark-gluon plasma. We find a significant suppression effect. For $p_T \sim 10$ GeV we obtained the reduction of the spectra by $\sim [20 - 30, 25 - 35, 30 - 40] \%$ at $\sqrt{s} = [0.2, 2.76, 7]$ TeV. We also discuss how this phenomenon may change the predictions for the nuclear modification factors for AA and pA collisions.

1. Introduction

The experiments at RHIC and LHC have provided clear evidence that in AA collisions the hadroproduction goes through the formation of a fireball of hot and dense quark-gluon plasma (QGP). This follows from the observation of strong suppression of high-$p_T$ particle spectra (the so-called jet quenching phenomenon) and from the results of the hydrodynamic simulations of AA collisions. In the pQCD paradigm the jet quenching is due to radiative [1, 2, 3, 4, 5, 6, 7] and collisional [8] energy loss in the QGP which soften the parton→hadron fragmentation in AA collisions (for recent comprehensive reviews, see [9, 10]). The suppression of the high-$p_T$ particle spectra in AA collisions is characterized by the nuclear modification factor $R_{AA}$ defined as the ratio of the particle spectrum in AA collisions to the binary-scaled spectrum in pp collisions [11]

$$R_{AA} = \frac{d\sigma(AA \to hX)/dp_Tdy}{N_{bin} d\sigma(pp \to hX)/dp_Tdy}. \quad (1)$$

Presently, in theoretical calculations of the $R_{AA}$ for the inclusive cross section $d\sigma(pp \to hX)/dp_Tdy$ in the denominator in (1) predictions of the pQCD are used. However, if the QGP is produced in pp collisions as well, the real inclusive cross section differs from that calculated in pQCD by its own medium modification factor $R_{pp}$, i.e.,

$$d\sigma(pp \to hX)/dp_Tdy = R_{pp}d\sigma_{pert}(pp \to hX)/dp_Tdy. \quad (2)$$

In this scenario the theoretical quantity which should be compared with the experimental $R_{AA}$ given by (1) can be written as

$$R_{AA} = R_{AA}^{st}/R_{pp}, \quad (3)$$
where $R_{AA}^{st}$ is the standard nuclear modification factor calculated using the pQCD predictions for the particle spectrum in $pp$ collisions. Of course, the $R_{pp}$ is unobservable directly because experimentally we do not have the baseline spectrum with the final state interactions in the QGP switched off. Nevertheless, the presence of the $R_{pp}$ in $\langle 3 \rangle$ may be important for theoretical predictions for jet quenching in $AA$ collisions. For example, for the jet flavor tomography of the QGP [12, 13, 14, 15] due to different suppression of light and heavy flavors in $pp$ collisions.

Presently, it is widely believed that in $pp$ collisions in the studied energy range a hot QCD matter is not produced in the typical inelastic minimum bias events due to small energy density. But in high multiplicity (HM) $pp$ events the energy density may be comparable to that in $AA$ collisions at RHIC and LHC energies. And if the thermalization time, $\tau_0$, is small enough, say $\tau_0 \lesssim 0.5$ fm, the mini-QGP with size of $\sim 2 – 3$ fm should be formed quite likely to the large-size plasma in $AA$ collisions. In recent years the possibility that the mini-QGP may be created in HM $pp$ events has attracted increasing interest (see, for instance, Refs. [16, 17, 18, 19, 20, 21, 22]). Actually, we already have some experimental indications in favor for the formation of the mini-QGP in HM $pp$ collisions. It is possible that the ridge correlation structure in HM $pp$ events at $\sqrt{s} = 7$ TeV observed by the CMS collaboration [23] is due to the transverse flow of the QGP. In [19], employing Van Hove’s idea [24] that phase transition should lead to anomalous behavior of the mean transverse momentum $\langle p_T \rangle$ as a function of multiplicity, it has been argued that the data on $\langle p_T \rangle$ signal possible plasma formation in the domain $dN_{ch}/d\eta \sim 6 – 24$. Some intriguing similarities between the results of the femtoscopic analyses of $pp$ and $AA$ collisions at RHIC [25] and LHC [26] also may signal the formation of the collective QCD matter in the HM $pp$ events. The preliminary data from ALICE [27], indicating that for the HM $pp$ events jets undergo a softer fragmentation, also support this idea.

From the point of view of jet quenching it is important that the conditions for the QGP production in $pp$ collisions are better in events with jets, because the multiplicity of soft off-jet particles (the so-called underlying events (UE), see [28] for a review) is enhanced by a factor of $2 – 3$ [29] (below we denote this factor by $K_{ue}$) as compared to the minimum bias multiplicity. And even at RHIC energies $\sqrt{s} \sim 0.2$ TeV the UE multiplicity may be high enough for the QGP formation. In our recent work [30] we have evaluated the medium modification of the fragmentation functions (FFs) for $\gamma$-triggered and inclusive jets in HM $pp$ collisions, and have presented preliminary results for medium suppression of hadron spectra. We have found that the medium effects are surprisingly strong. In the present work we perform a detailed analysis of the medium modification of the hadron spectra in $pp$ collisions due to parton energy loss in the mini-QGP. We evaluate $R_{pp}$ of charged hadrons in the central rapidity region ($y = 0$) at RHIC ($\sqrt{s} = 0.2$ TeV) and LHC ($\sqrt{s} = 2.76$ and 7 TeV) energies. We also address the effect of $R_{pp}$ on $R_{AA}$ at RHIC and LHC energies and on $R_{pA}$ in the context of the recent data from ALICE [31] on $R_{pPb}$ at $\sqrt{s} = 5.02$ TeV. The analysis is based on the light-cone path integral (LCPI) approach [3, 4] to induced gluon emission. It treats accurately the
finite-size and Coulomb effects (which are very important for the mini-QGP), the mass effects, and is valid beyond the soft gluon approximation. We evaluate the medium modified FFs within the scheme developed previously for AA collisions [32]. It takes into account both radiative and collisional energy loss. Previously in [33,14,15] the approach has been successfully used for description of jet quenching in AA collisions.

The paper is organized as follows. In the next section we discuss the parameters of the mini-QGP for the UE pp events at RHIC and LHC. In section 3 we discuss the basic aspects of the theoretical framework. In section 4 we present the numerical results on parton energy loss in the mini-QGP and the medium modification factors for pp, AA and pA collisions, section 5 summarizes our work.

2. Mini-QGP in proton-proton collisions

We neglect the transverse expansion of the mini-QGP and use 1+1D Bjorken’s model [34], which gives \( T_3^0 \tau_0 = T_3 \tau \). For \( \tau < \tau_0 \) we take medium density \( \propto \tau \). In the basic variant we take \( \tau_0 = 0.5 \) fm. Approximately such \( \tau_0 \) is used in most studies of jet quenching in AA collisions. For the QGP produced in AA collisions with the lifetime/size \( L \gg \tau_0 \) the medium modification of hadron spectra is not very sensitive to variation of \( \tau_0 \). But this may be untrue for the mini-QGP in pp collisions when the plasma size is not very large as compared to \( \tau_0 \). To understand the sensitivity of \( R_{pp} \) to \( \tau_0 \), which is not well constrained by the hydrodynamic modeling of AA collisions [35], we also perform calculations for \( \tau_0 = 0.8 \) fm. As in our analyses of AA collisions [32,33,14], we neglect variation of the initial temperature \( T_0 \) with the transverse coordinates. To fix \( T_0 \) we use the entropy/multiplicity ratio \( C = \frac{dS}{dy}/\frac{dN_{ch}}{d\eta} \approx 7.67 \) obtained in [36]. The initial entropy density can be written as

\[
s_0 = \frac{C}{\tau_0^2 R_f^2} \frac{dN_{ch}}{d\eta},
\]

where \( R_f \) is the radius of the created mini-QGP fireball. We ignore the azimuthal anisotropy, and regard the \( R_f \) as an effective plasma radius, which includes pp collisions at all impact parameters. This approximation seems to be plausible since anyway the jet production should be dominated by the almost head-on collisions for which the azimuthal effects should be small. This is supported by calculation of the distribution of jet production cross section in the impact parameter plane using the MIT bag model which says that only 25% of jets come from pp collisions with the impact parameter larger than the bag radius. It says that typically the fireball has a relatively small eccentricity. Anyway, we are interested in \( R_{pp} \), which is averaged over the azimuthal angle, and it is practically insensitive to the fireball eccentricity.

One can expect that for pp collisions the typical radius of the fireball should be about the proton radius \( R_p \sim 1 \) fm. It agrees qualitatively with \( R_f \) obtained for pp collisions at \( \sqrt{s} = 7 \) TeV in numerical simulations performed in [21] within the IP-Glasma model [37]. The \( R_f \) from [21] grows approximately as linear function of \( (dN_g/dy)^{1/3} \) and then flatten. The flat region corresponds to almost head-on collisions. In this regime the
fluctuations of multiplicity are dominated by the fluctuations of the glasma color fields \[ \text{[21].} \] We use the \( R_f \) from \[ \text{[21]} \] parametrized in \[ \text{[38]} \] via \( dN_g/dy \) in the form

\[
R_f = 1 \text{ fm} \times f_{pp} \left( \sqrt{dN_g/dy} \right)
\]

with

\[
f_{pp}(x) = \begin{cases} 
0.387 + 0.0335x + 0.274 x^2 - 0.0542 x^3 & \text{if } x < 3.4, \\
1.538 & \text{if } x \geq 3.4.
\end{cases}
\]

We evaluate \( R_f \) taking \( dN_g/dy = \kappa dN_{ch}/d\eta \) with \( \kappa = C45/2\pi^4\xi(3) \approx 2.13 \). Possible increase of the \( R_f \) from RHIC to LHC should not be important since our results are not very sensitive to variation of \( R_f \).

The multiplicity density of the UEs grows with momentum of the leading charged jet hadron at \( p_T \lesssim 3 - 5 \text{ GeV} \) and then flatten \[ \text{[29, 39, 40, 41, 42]} \] (in terms of the jet energy the plateau region corresponds approximately to \( E_{jet} \gtrsim 15 - 20 \text{ GeV} \)). To fix the \( dN_{ch}/d\eta \) in \[ \text{[1]} \] at \( \sqrt{s} = 0.2 \text{ TeV} \) we use the UE multiplicity enhancement factor \( K_{ue} \) from PHENIX \[ \text{[39]} \] obtained by dihadron correlation method. Taking for the minimum bias non-diffractive events \( dN_{ch}^{mb}/d\eta = 2.98 \pm 0.34 \) from STAR data \[ \text{[43]} \], we obtained for the UEs in the plateau region \( dN_{ch}/d\eta \approx 6.5 \). To evaluated the UE multiplicity at \( \sqrt{s} = 2.76 \) and \( 5.02 \text{ TeV} \) we use the data from ATLAS \[ \text{[40]} \] at \( \sqrt{s} = 0.9 \) and \( 7 \text{ TeV} \) that give in the plateau region \( dN_{ch}/d\eta \approx 7.5 \) and \( 13.9 \). Assuming that \( dN_{ch}/d\eta \propto s^{\delta} \) by interpolating between \( \sqrt{s} = 0.9 \text{ TeV} \) and \( 7 \text{ TeV} \) we obtained for the UE multiplicity density in the plateau region \( dN_{ch}/d\eta \approx 10.5 \) and \( 12.6 \) at \( \sqrt{s} = 2.76 \) and \( 5.02 \text{ TeV} \), respectively. With the above values of the UE multiplicity densities in the plateau regions we obtain the following values for the fireball radii

\[
R_f[\sqrt{s} = 0.2, 2.76, 5.02, 7 \text{ TeV}] \approx [1.3, 1.44, 1.49, 1.51] \text{ fm}.
\]

With these radii, using \[ \text{[41]} \] and the ideal gas formula \( s = (32/45 + 7N_f/15)T^3 \) (with \( N_f = 2.5 \)), we obtain the initial temperatures of the QGP

\[
T_0[\sqrt{s} = 0.2, 2.76, 5.02, 7 \text{ TeV}] \approx [199, 217, 226, 232] \text{ MeV}.
\]

One can see that the values of \( T_0 \) lie well above the deconfinement temperature \( T_c \approx 160 - 170 \text{ MeV} \) \[ \text{[44, 45]} \]. For such initial temperatures the purely plasma phase may exist up to the \( \tau_{QGP} \sim 1 - 1.5 \text{ fm} \), and beyond \( \tau_{QGP} \) the hot QCD matter will evolve in the mixed phase up to \( \tau_{max} \sim 2R_f \) where the transverse expansion should lead to a fast cooling of the system. Since in the interval \( \tau_{QGP} < \tau < \tau_{max} \) the QGP fraction in the mixed phase is approximately \( \propto 1/\tau \) \[ \text{[34]} \] we can use in calculating jet quenching the \( 1/\tau \) dependence of the number density of the scattering centers in the whole range of \( \tau \) (but with the Debye mass defined for \( T \approx T_c \) at \( \tau > \tau_{QGP} \)).

Although we neglect the transverse expansion of the QCD matter, it should not lead to large errors in our predictions. As was demonstrated in \[ \text{[46]} \] the transverse

\[ \text{\footnote{In fact, if one uses the entropy from the lattice calculations \[ \text{[44, 45]} \] the fireball temperatures in \[ \text{[8]} \] will be higher by \( \sim 10 - 15\% \). We ignore this difference since for jet quenching the crucial quantity is the entropy, which we take from experimental data (see discussion below in Sec. IV).}} \]
motion does not affect strongly jet quenching in AA collisions. Physically it is due to an almost complete compensation between the enhancement of the energy loss caused by increase of the medium size and its suppression caused by reduction of the medium density. In pp collisions the effect should be even smaller since the typical formation length for induced gluon emission is of the order of $R_f$ or larger. In such a regime the parton energy loss is mostly controlled by the mean amount of the matter traversed by fast partons, and the details of the density profile along the jet trajectory are not very important. Also, in pp collisions the QCD matter spends much time in the mixed phase, where the sound velocity becomes small and the transverse expansion should be less intensive than in AA collisions.

We conclude this section with two additional remarks. First, naively one could think that for evaluation of the medium suppression of the minimum bias pp spectrum one should use the minimum bias multiplicity density in evaluating the mini-QGP parameters. But it would be wrong. Indeed, the minimum bias events include events with and without jet production, and the minimum bias high-$p_T$ spectrum is related to events with jet (at least one) production. The corresponding multiplicity density for such events is exactly the UE $dN_{ch}/d\eta$.

The second remark concerns the formula (4). It implicitly assumes that the UE multiplicity distribution $dN_{ch}/d\eta$, likewise the minimum bias multiplicity density, has a central plateau in rapidity (we assume that jet is produced at $y = 0$). For typical inelastic events the existence of the central plateau is a consequence of the approximate longitudinal boost invariance. In the glasma picture [21] it naturally appears due to boost invariance of the initial glasma color fields. However, for the UE events this invariance is broken by the jet production at $y = 0$. And in principle there may be a bump in the UE multiplicity distribution near the jet rapidity, say due to the initial state radiation. For (4) to be applicable the half-width of the bump should satisfy the inequality

$$\Delta \eta \gtrsim c_s \ln(\tau_f/\tau_0),$$

(9)

where $\tau_f$ is the freezeout time and $c_s$ is the sound velocity of the matter. This inequality ensures that in the whole interval from $\tau_0$ to $\tau_f$ the edges of the bump do not affect the $\tau$-dependence of the entropy $s \propto 1/\tau$. For the relevant temperature range $c_s \lesssim 0.5 \ [45]$. Then taking $\tau_f \sim \tau_{max}$ and $\tau_{max}/\tau_0 \sim 6$ we obtain from (9) $\Delta \eta \gtrsim 1$. The fact that data on the UE $dN_{ch}/d\eta$ from ATLAS [40] obtained for $|\eta| < 2.5$ and from ALICE [42] obtained for $|\eta| < 0.8$ agree between each other says that the above inequality is satisfied. And the formula (4) can be safely used for the UEs.

3. Sketch of the Calculations

We now turn to the jet quenching in the mini-QGP produced in pp collisions. Our treatment is similar to that for AA collisions in our previous analysis [32] to which we refer the reader for details. Here we give only a brief sketch of the calculations, focusing
3.1. Perturbative and medium modified inclusive cross sections

As usual we write the perturbative inclusive cross section in (2) in terms of the vacuum parton→hadron FF $D_{h/i}$

$$\frac{d\sigma_{\text{pert}}(pp \to hX)}{dp_T^h dy} = \sum_i \int_0^1 dz \frac{z^2 D_{h/i}(z, Q)}{z} \frac{d\sigma(pp \to iX)}{dp_{T,i}^i dy},$$  \hspace{1cm} (10)

where $d\sigma(pp \to iX)/dp_{T,i}^i dy$ is the ordinary hard cross section, $p_{T,i}^i = p_T/z$ is the parton transverse momentum. We write the real inclusive cross section, which accounts for the final state interactions in the QCD matter, in a similar form but with the medium modified FF $D_{h/i}^m$

$$\frac{d\sigma(pp \to hX)}{dp_T^h dy} = \sum_i \int_0^1 dz \frac{z^2 D_{h/i}^m(z, Q)}{z} \frac{d\sigma(pp \to iX)}{dp_{T,i}^i dy}.$$  \hspace{1cm} (11)

Here it is implicit that $D_{h/i}^m$ is averaged over the geometrical variables of the hard parton process and the impact parameter of $pp$ collision.

The formula (11) can be viewed as an analogue of the formula for the minimum bias $R_{AA}$ defined in the whole centrality (impact parameter) range. However, there is one important difference between $pp$ and $AA$ collisions. In $AA$ collisions at a given impact parameter the fluctuations of the multiplicity (and of the parameters of the fireball) are small. And this allows to relate the centrality (defined through the multiplicity) to the impact parameter. In $pp$ collisions one cannot relate the multiplicity density to the impact parameter, since for each impact parameter and jet production point in the transverse plane (which can be localized with the accuracy $\sim z/p_T$) the fluctuations of the multiplicity density are large. These fluctuations, together with the event-by-event fluctuations of the impact parameter and the jet production point, give the observable fluctuating UE $dN_{ch}/d\eta$, which can be translated into the fluctuating fireball parameters. However, the detailed dynamics of the UEs and of the multiplicity fluctuations in such events is far from being clear. In particular, we do not know whether the enhancement of the UE multiplicity is only due to the fact that the jet production is biased to more central collisions and to which extent it may be related to the increase of the soft gluon density in jet production due to the initial state radiation. Therefore an accurate accounting for the fluctuations of the parameters of the mini-QGP fireball is impossible. In the present study in evaluating $D_{h/i}^m$ we take into account (approximately) only the event-by-event variations of the geometrical parameters (see below), but ignore the fluctuations of the UE $dN_{ch}/d\eta$. And evaluate the parameters of the fireball simply using the typical UE multiplicity density, although technically the inclusion of the fluctuations of the UE $dN_{ch}/d\eta$ in our formalism is quite simple, and we do it to estimated the accuracy of our approximation (see below).
As in \[32\], we calculated the hard cross sections in the LO pQCD with the CTEQ6 \[47\] parton distribution functions (PDFs). To simulate the higher order effects in calculating the partonic cross sections we take for the virtuality scale in $\alpha_s$ the value $cQ$ with $c = 0.265$ as in the PYTHIA event generator \[48\]. This gives a fairly good description of the $p_T$-dependence of the spectra in $pp$ collisions. Of course, in principle, in the scenario with the QGP formation for a fully consistent treatment of $R_{pp}$ (and $R_{AA}$) one should use a bootstrap procedure and compare with the experimental data not the perturbative cross section \(\text{(10)}\) but the real one given by \(\text{(11)}\), and namely the latter should be adjusted (say, by varying PDFs, FFs, and $\alpha_s$) to describe experimental data. However, since the hadron spectra have very steep $p_T$-dependence (as compared to a relatively weak $p_T$-dependence of $R_{pp}$) this inconsistency may be safely ignored in calculating $R_{pp}$ (the same is true for $R_{AA}$ and $R_{pA}$).

For the hard scale $Q$ in the FFs in \(\text{(10)}, \text{(11)}\) we use $p_T/z$. We calculate the vacuum FFs $D_{h/j}$ as a convolution of the KKP \[49\] parton→hadron FFs at soft scale $Q_0 = 2$ GeV with the DGLAP parton→parton FFs $D_{j/i}^{DGLAP}$ describing the evolution from $Q$ to $Q_0$. The latter have been computed with the help of PYTHIA \[48\]. This procedure reproduces well the whole $Q$-dependence of the KKP \[49\] parametrization of the vacuum FFs. For a given fast parton path length in the QGP the medium modified FFs $D_{j/i}^m$ have been calculated in a similar way but inserting between the DGLAP parton→parton FFs and the KKP parton→hadron FFs the parton→parton FFs $D_{j/i}^{ind}$ which correspond to the induced radiation stage in the QGP. The induced radiation FFs $D_{j/i}^{ind}$ have been calculated from the medium induced gluon spectrum using Landau's method \[50\] imposing the flavor and momentum conservation (again, we refer the interested reader to \[32\] for details).

Note that, since in both the vacuum and the medium modified FFs the DGLAP evolution is accounted for in the same way, the medium effects vanish strictly at zero matter density, as it must be. The above approximation with the time ordered and independent DGLAP and induced radiation stages, suggested for the large-size plasma produced in $AA$ collisions \[32\], seems to be reasonable for the mini-QGP as well (at least in the jet energy region $\lesssim 30 – 50$ GeV where the suppression effect appears to be strongest) since the typical formation time for the most energetic DGLAP gluons is of the order of (or smaller) than the thermalization time $\tau_0$. It is worth noting that, although the time ordering of the DGLAP and induced radiation stages seems to be physically reasonable, the permutation of these stages in the above convolution gives a very small effect \[32\].

Since we do not consider the azimuthal effects, the averaging of the medium modified FFs over the geometrical variables of the hard parton process and $pp$ collisions in the impact parameter plane is simply reduced to averaging over the parton path length $L$ in the QGP. It cannot be performed accurately since the distribution of hard processes in the impact parameter plane is not known yet. But one can expect that the effect of $L$ fluctuations should be relatively small for any more or less centered distribution of energetic partons in the proton wave function. We have performed averaging over $L$
using the distribution of hard processes in the impact parameter plane obtained with the quark distribution from the MIT bag model (we assume that the valence quarks and the hard gluons radiated by the valence quarks follow approximately the same distribution in the transverse spatial coordinates). Calculations within this model show that practically in the full range of the impact parameter of $pp$ collisions the distribution in $L$ is sharply peaked around $L \approx \sqrt{S_{ov}/\pi}$, where $S_{ov}$ is the overlap area for two colliding bags. It means that our fireball radius $R_f$ (which includes all centralities) at the same time gives the typical path length for fast partons. Our calculations show that the effect of the $L$-fluctuations on $R_{pp}$ is relatively small. As compared to $L = R_f$ they reduce the medium modification by $\sim 10 - 15\%$.

As in our previous studies of jet quenching in $AA$ collisions we treat the collisional energy loss, which is relatively small [51], as a small perturbation to the radiative mechanism. We incorporate it in the above procedure simply by renormalizing the QGP temperature in calculating the medium modified FFs (see [32] for details). We assume that the collisional energy loss vanishes at $\tau < \tau_0$ in the pre-equilibrium stage which probably is populated by strong collective glasma color fields, and the concept of the collisional energy loss is hardly applicable in this region. On the contrary, it is clear that the coherent glasma fields can give some contribution to the radiative energy loss (probably rather small [52]). For this reason the use of the linearly growing plasma density at $\tau < \tau_0$ seems to be a plausible parametrization to model the transition from the glasma phase to the hydrodynamically evolving QGP, which of course cannot be abrupt.

### 3.2. Medium induced gluon spectrum and parameters of the model

As in [32] we evaluate the medium induced gluon spectrum $dP/dx$ ($x = \omega/E$ is the gluon fractional momentum) for the QGP modeled by a system of the static Debye screened color centers [1]. We use the Debye mass obtained in the lattice calculations [53] giving $\mu_D/T$ slowly decreasing with $T$ ($\mu_D/T \approx 3.2$ at $T \sim T_c, \mu_D/T \approx 2.4$ at $T \sim 4T_c$). For the quasiparticle masses of light quarks and gluon in the QGP we take $m_q = 300$ and $m_g = 400$ MeV supported by the analysis of the lattice data [54]. But the results are not very sensitive to the $m_q$, and practically insensitive to the value of $m_g$. We evaluated the induced gluon spectrum using the representation suggested in [55]. It expresses the $x$-spectrum for gluon emission from a quark (or gluon) through the light-cone wave function of the $gq\bar{q}$ (or $ggg$) system in the coordinate $\rho$-representation. The $z$-dependence of the wave function is governed by a two-dimensional Schrödinger equation with the “mass” $\mu = x(1-x)E$ ($E$ is the initial parton energy) in which the longitudinal coordinate $z$ plays the role of time and the potential $v(\rho)$ is proportional to the local plasma density/entropy times a linear combination of the dipole cross sections $\sigma(\rho), \sigma((1-x)\rho)$ and $\sigma(x\rho)$. Note that the physical pattern of induced gluon emission in the mini-QGP differs from that for the large-size QGP. For the mini-QGP when the typical path length in the medium $L \sim 1 - 1.5 \text{ fm}$ the energy loss is dominated by
gluons with $L_f \gtrsim L$, where $L_f \sim 2\omega/m_g^2$ is the gluon formation length in the low density limit. It is the diffusion regime in the terminology of [56], in which the finite-size effects play a crucial role. In this regime the dominating contribution comes from the $N = 1$ rescattering and the Coulomb effects are very important [56]. On the contrary, for the QGP in $AA$ collisions a considerable part of the induced energy loss comes from gluons with $L_f \lesssim L$. Indeed, in the bulk of the large-size QGP $L_f \sim 2\omega S_{LPM}/m_g^2$, where $S_{LPM}$ is the LPM suppression factor. For RHIC and LHC typically $S_{LPM} \sim 0.3 - 0.5$ for $\omega \sim 2$ GeV, it gives $L_f \sim 1.5 - 2.5$ fm which is smaller than the typical $L$ for the QGP in $AA$ collisions. In this regime the finite-size effects are much less important and induced gluon radiation is (locally) approximately similar to that in an infinite extent matter.

From the point of view of jet quenching in $pp$ collisions it is important that induced radiation in the mini-QGP is more perturbative than in the QGP in $AA$ collisions. Indeed, let us consider induced radiation for the mini-QGP. From the Schrödinger diffusion relation one can obtain for the typical transverse size of the three parton system $\rho^2 \sim 2\xi/\omega$, where $\xi$ is the path length after gluon emission. Then, using the fact that $\sigma(\rho)$ is dominated by the $t$-channel gluon exchanges with virtualities up to $Q^2 \sim 10/\rho^2$ [57] we obtain $Q^2 \sim 5\omega/\xi$. For $\omega \sim 2$ and $\xi \sim 0.5 - 1$ fm it gives rather large virtuality scale $Q^2 \sim 2 - 4$ GeV$^2$. The virtuality scale for $\alpha_s$ in the gluon emission vertex has a similar form but smaller by a factor of $\sim 2.5$ [51]. The $1/\xi$ dependence of $Q^2$ persists up to $\xi \sim L_f$. For the large-size QGP in the above formulas one should replace $\xi$ by the real in-medium $L_f$ (which contains $S_{LPM}$) which is by a factor of $\sim 2$ larger than the typical values of $\xi$ for the mini-QGP. It results in a factor of $\sim 2$ smaller virtualities for the QGP in $AA$ collisions. In this sense the calculations for the mini-QGP are more robust than for the large-size QGP.

As in [32, 33, 14, 15] we perform calculations of radiative and collisional energy loss with running $\alpha_s$ frozen at some value $\alpha_s^{fr}$ at low momenta. For gluon emission in vacuum a reasonable choice is $\alpha_s^{fr} \sim 0.7 - 0.8$ [58, 59]. But in plasma thermal effects can suppress $\alpha_s^{fr}$. However, in principle, the extrapolation from the vacuum gluon emission to the induced radiation is unreliable due to large theoretical uncertainties of jet quenching calculations. For this reason $\alpha_s^{fr}$ should be treated as a free parameter of the model. To evaluate the medium suppression in $pp$ collisions it is reasonable to use the information on the values of $\alpha_s^{fr}$ necessary for description of jet quenching in $AA$ collisions. Previously we have observed [15] that data on $R_{AA}$ are consistent with $\alpha_s^{fr} \sim 0.5$ for RHIC and $\alpha_s^{fr} \sim 0.4$ for LHC. The reduction of $\alpha_s^{fr}$ from RHIC to LHC may be related to stronger thermal effects in the QGP due to higher initial temperature at LHC. But the analysis [15] is performed under assumption that there is no medium suppression in $pp$ collisions. The inclusion of $R_{pp}$ should increase the values of $\alpha_s^{fr}$. However, in [15] we used the plasma density vanishing at $\tau < \tau_0$, whereas in the present work we use the QGP density $\propto \tau$ in this region which leads to stronger medium suppression. As a result, as we will see below, the values of $\alpha_s^{fr}$, which are preferable from the standpoint of the description of the data on $R_{AA}$, remain approximately the same, or a bit larger, as obtained in [15]. If the difference between the preferable values
of $\alpha_s^{fr}$ for AA collisions at RHIC and LHC is really due to the thermal effects, then for the mini-QGP with $T_0$ as given in (8) a reasonable window is $\alpha_s^{fr} \sim 0.6 - 0.7$. In principle for the mini-QGP the thermal reduction of $\alpha_s$ may be smaller than that for the large-size plasma (at the same temperature). Since at $L_f \lesssim L$, which typically holds for the mini-QGP, the dominating contribution to the induced gluon spectrum comes from configurations with interference of the emission amplitude and complex conjugate one when one of them has the gluon emission vertex outside the medium and is not affected by the medium effects at all. We perform the calculations for $\alpha_s^{fr} = 0.5, 0.6$ and 0.7. Note that, in principle $R_{pp}$ should be less sensitive to $\alpha_s^{fr}$ than $R_{AA}$ since, as we said above, the typical virtualities for induced gluon emission in the mini-QGP are larger than that in the large-size QGP. As will be seen from our numerical results, the sensitivity to $\alpha_s^{fr}$ is really quite weak.

4. Numerical Results

4.1. Energy loss in the mini-QGP

Before presenting the results for the medium modification factors it is worthwhile first to show the results for radiative and collisional energy loss that may give some insight into the magnitude of the medium effects generated by the mini-QGP in $pp$ collisions. In Fig. 1 we show the energy dependence of the total (radiative plus collisional) and collisional energy loss for partons produced in the center of the mini-QGP fireball for RHIC and LHC conditions for $\alpha_s^{fr} = 0.6$. Both the radiative and collisional contributions are calculated for the lost energy smaller than half of the initial parton energy. The fireball radius $R_f$ and the initial temperature $T_0$ have been calculated with the UE multiplicity density dependent on the jet energy $E$ using the data [39, 40]. In [39, 40] the UE activity has been measured vs the transverse momentum of the leading charged jet hadron (we denote it as $p_{T}^l$). To obtain the UE $dN_{ch}/d\eta$ as a function of the jet energy $E$ we neglect the fluctuations of $p_{T}^l$ for a given $E$ and use the rigid relation $p_{T}^l = \langle z_l \rangle E$, where $\langle z_l \rangle$ is the average fractional momentum of the leading jet hadron. For the $\langle z_l \rangle$ we take the PYTHIA predictions, which gives in the relevant energy region ($E \lesssim 10$ GeV) $\langle z_l \rangle \sim 0.26$ for gluon jets and $\langle z_l \rangle \sim 0.36$ for quark jets. The jet energy dependence of the parameters of the fireball becomes important only for partons with $E \lesssim 10 - 15$ GeV. At higher energies the UE $dN_{ch}/d\eta$ is flatten and $R_f$ and $T_0$ are very close to that given by (7) and (8). And radiative and collisional energy loss may be calculated using (7), (8). To illustrate it in Fig. 1 we presented the results for the total energy loss obtained for the fireball parameters for the UE $dN_{ch}/d\eta$ in the plateau region ($p_{T}^l \gtrsim 5$ GeV). From Fig. 1 one can see that the energy loss for these two versions of the fireball parameters (solid and long-dashed lines) become very close to each other at $E \gtrsim 10$ GeV. This says that the decrease of the UE multiplicity density at $p_{T}^l \lesssim 5$ GeV should be practically unimportant for $R_{pp}(p_T)$ already at $p_T \gtrsim 7 - 10$ GeV.

Fig. 1 shows that the parton energy loss in the mini-QGP turns out to be quite
Figure 1. Energy dependence of the energy loss of gluons (upper panels) and light quarks (lower panels) produced in the center of the mini-QGP fireball at $\sqrt{s} = 0.2$ TeV (left) and $\sqrt{s} = 2.76$ TeV (right). Solid line: total (radiative plus collisional) energy loss calculated with the fireball radius $R_f$ and the initial temperature $T_0$ obtained with the UE $dN_{ch}/d\eta$ dependent on the initial parton energy $E$; dashed line: same as solid line but for collisional energy loss; long-dashed line: same as solid line but for $R_f$ and $T_0$ obtained with the UE $dN_{ch}/d\eta$ in the plateau region as given by (7) and (8). All the curves are for $\alpha_s^f = 0.6$.

Figure 2. Left: Radiative (solid) and collisional (dashed) gluon energy loss vs the path length $L$ in the QGP with $T_0 = 199$ MeV for (bottom to top) $E = 20$ and 50 GeV. The dotted lines show radiative energy loss for $T_0 = 320$ MeV rescaled by the factor $(199/320)^3$. All curves are calculated for $\alpha_s^f = 0.6$. Right: same as in the left figure but for $T_0 = 217$ and 420 MeV and the rescaling factor $(217/420)^3$ for dotted lines.
large. At $E \sim 10 - 20$ GeV for gluons the total energy loss is $\sim 10 - 15\%$ of the initial energy. The contribution of the collisional mechanism is relatively small. The energy loss for the mini-QGP shown in Fig. 1 is smaller than that obtained in [15] for the large-size QGP in AA collisions by only a factor of $\sim 4$.

To illustrate the $L$-dependence of the parton energy loss in our model in Fig. 2 we show the results for the radiative and collisional gluon energy loss vs the path length $L$ for $E = 20$ and 50 GeV for $T_0 = 199$ and 217 MeV, corresponding to $\sqrt{s} = 0.2$ and 2.76 TeV. To show the difference between the QGP produced in $pp$ and AA collisions we present also predictions for radiative energy loss for $T_0 = 320$ and 420 MeV corresponding to central $Au + Au$ collisions at $\sqrt{s} = 0.2$ TeV, and for $T_0 = 420$ MeV corresponding to central $Pb + Pb$ collisions at $\sqrt{s} = 2.76$ TeV (the procedure that leads to these values of $T_0$ is described in [15]). To illustrate the temperature dependence better we rescaled the predictions for AA collisions by the factor $(T_0(pp)/T_0(AA))^3$. One can see that at $L \geq \tau_0$ the radiative energy loss is approximately a linear function of $L$. At $L < \tau_0$ the radiative energy loss is approximately $\propto L^3$ (since the leading $N = 1$ rescattering contribution to the effective Bethe-Heitler cross section is $\propto L$ [56, 60] and integration over the longitudinal coordinate of the scattering center gives additional two powers of $L$). The comparison of the radiative energy loss for $T_0 = 199$ and 217 MeV to that for $T_0 = 320$ and 420 MeV shows deviation from the $T^3$ scaling by factors of $\sim 1.5$ and $\sim 2$, respectively. One can see that this difference persists even at $L \sim 1$ fm. This deviation from the $T^3$ scaling comes mostly from the increase of the LPM suppression (and partly from the increase of the Debye mass) for the QGP produced in AA collisions.

4.2. Jet quenching in pp collisions

As we have seen from Fig. 1 the dependence of the UE multiplicity density from [39, 40] on the momentum of the leading jet hadron $p_T$ practically does not affect the parton energy loss at $E \gtrsim 10 - 15$ GeV, which from the standpoint of the particle spectra corresponds approximately to $p_T \gtrsim 7 - 10$ GeV. To account for the effect of the $p_T$-dependence of the UE $dN_{ch}/d\eta$ on $R_{pp}(p_T)$ at $p_T \lesssim 10$ GeV we use the rigid approximation $p_T' = \langle z_i \rangle E$ as in the above calculations of the energy loss. And in addition ignore the fluctuations of the variable $z$ in (10) (since the integrand of (10) is quite sharply peaked about $\langle z \rangle$). In this approximation we can write $p_T' = p_T/\eta$, where $p_T$ is the momentum of the observed particle in (10) and $\eta = \langle z \rangle / \langle z_i \rangle$. Jet simulation with PYTHIA [48] shows that for jets with energy $E \lesssim 10 - 15$ GeV, that can feel the energy dependence of the UE multiplicity, one can take $\eta \sim 2.1$ for $\sqrt{s} = 0.2$ TeV and $\eta \sim 1.9$ at LHC energies $\sqrt{s} \gtrsim 2.76$ TeV. The uncertainties from this prescription are restricted to the region $p_T \lesssim 7 - 10$ GeV. However, even in this region it should work on average (in the sense that the fluctuations will just smear the $p_T$-dependence of the medium suppression in this region). This problem becomes completely irrelevant for $R_{pp}$ at $p_T \gtrsim 10$ GeV.

In Fig. 3 we present the results for $R_{pp}$ of charged hadrons at $\sqrt{s} = 0.2, 2.76$ and 7
TeV for $\alpha_{fr}^s = 0.5, 0.6$ and $0.7$. To illustrate the sensitivity of the results to $\tau_0$ we show the curves for $\tau_0 = 0.5$ and $0.8$ fm. One can see that the suppression effect for the basic variant with $\tau_0 = 0.5$ fm turns out to be quite large at $p_T \lesssim 20$ GeV both for RHIC and LHC. Fig. 3 shows that for $\tau_0 = 0.8$ fm the reduction of the suppression is not very significant. One can see that, as we expected, $R_{pp}$ does not exhibit a strong dependence on $\alpha_{fr}^s$. Although the plasma density is smaller at $\sqrt{s} = 0.2$ TeV, the suppression effect is approximately similar to that at $\sqrt{s} = 2.76$ and 7 TeV. It is due to a steeper slope of the hard cross sections at $\sqrt{s} = 0.2$ TeV. The increase in the suppression from $\sqrt{s} = 2.76$ to $\sqrt{s} = 7$ TeV is relatively small.

To understand the sensitivity of $R_{pp}$ to the fireball radius we also performed the calculations for the fireball radii calculated with (5), (6) times 0.7 and 1.3. We obtained in these two cases the reduction of the medium suppression by $\sim 3\%$ and $10\%$, respectively. The weak dependence on the value of $R_f$ is due to a compensation between the enhancement of the energy loss caused by increase of the fireball size and its suppression caused by reduction of the fireball density. The stability of $R_{pp}$ against variations of $R_f$ gives a strong argument that the errors due to the neglect of the variation of the plasma density in the transverse coordinates should be small. Indeed, the dominating $N = 1$ rescattering contribution to the radiative energy loss is a linear functional of the plasma density profile along the fast parton trajectory. It means that the results for a more realistic distribution of the initial plasma density in the impact parameter plane with a higher density in the central region can be roughly approximated by a linear superposition of the results obtained for the step density distributions (with different $R_f$) that should lead to approximately the same $R_{pp}$ as our calculations. Note also that since the variation of the plasma density in our test is very large (by a factor of $\sim 3.5$) this stability at the same time indicates indirectly that the effect of the neglected

\[ \frac{\partial R_{pp}}{\partial R_f} \]

The fact that for $0.7 R_f$ and $1.3 R_f$ the variations of $R_{pp}$ have the same sign is not very surprising since we use a wide window in the the fireball size. In this situation the second order term in the Taylor expansion of $R_{pp}$ around $R_f$ may be bigger than the linear term.

\[ \frac{\partial R_{pp}}{\partial R_f} \]
hydrodynamical variation of the plasma density should be small as well.

The results shown in Fig. 3 are obtained using the typical UE multiplicity density. As we said in Sec. 3, an accurate accounting for the fluctuations of the UE $dN_{ch}/d\eta$ is impossible since it should be done on the event-by-even basis (in the sense of the impact parameter and the jet production point), and requires detailed information about dynamics of the UEs. To understand how large the theoretical uncertainties, related to the event-by-event fluctuations of the UE $dN_{ch}/d\eta$, might be we evaluated $R_{pp}$ assuming that the distribution in the UE $dN_{ch}/d\eta$ is the same at each impact parameter and jet production point. We performed the calculations using the distribution in $dN_{ch}/d\eta$ from CMS [41] measured at $\sqrt{s} = 0.9$ and 7 TeV. It obeys approximately KNO scaling low similar to that in the minimum bias events [61]. For this reason one can expect that it can be used to estimate the effect of the multiplicity fluctuations for RHIC conditions as well. Our results show that for the fluctuating $dN_{ch}/d\eta$ the magnitude of $(1 - R_{pp})$ is reduced by only $\sim 5 - 6\%$ both for RHIC and LHC energies. This says that our approximation without the event-by-event fluctuations of the fireball parameters is quite good, since it is very unlikely that an event-by-event analysis may change significantly the results obtained using the total fluctuations of the UE multiplicity density.

4.3. Jet quenching in AA collisions

Although $R_{pp}$ is unobservable quantity it can alter the results of the jet tomography of AA collisions. To illustrate the possible effect of the mini-QGP in $pp$ collisions on $R_{AA}$ we show in Fig. 4 the comparison of our results for $R_{AA}$ with the data for $\pi^0$-mesons in central Au+Au collisions at $\sqrt{s} = 0.2$ TeV (a) from PHENIX [62], and with the data for charged hadrons in central Pb+Pb collisions at $\sqrt{s} = 2.76$ TeV (b,c) from ALICE [63] and CMS [64]. We show the results obtained with (solid) the $1/R_{pp}$ factor, i.e. for $R_{AA}$ defined by (3), and the results without (dashed) this factor, i.e. for $R_{AA}^{st}$. We use the $R_{pp}$ obtained with $\alpha_{fr}^{st} = 0.6$. We present the curves for $R_{AA}^{st}$ obtained with $\alpha_{s}^{fr} = 0.5$ and 0.6 for $\sqrt{s} = 0.2$ TeV, and with $\alpha_{s}^{fr} = 0.4$ and 0.5 for $\sqrt{s} = 2.76$ TeV. Since these values give better agreement with the data of the $R_{AA}$ given by (3). In calculating the hard cross sections for AA collisions we account for the nuclear modification of the PDFs with the EKS98 correction [65]. As in [15] we take $T_0 = 320$ MeV for central Au+Au collisions at $\sqrt{s} = 0.2$ TeV, and $T_0 = 420$ MeV for central Pb+Pb collisions at $\sqrt{s} = 2.76$ TeV obtained from hadron multiplicity pseudorapidity density $dN_{ch}/d\eta$ from RHIC [66] and LHC [67, 68]. One can see that at $p_T \sim 10$ GeV for RHIC the agreement of the theoretical $R_{AA}$ (with the $1/R_{pp}$ factor) with the data is somewhat better for $\alpha_{s}^{fr} = 0.6$, and for LHC the value $\alpha_{s}^{fr} = 0.5$ seems to be preferred by the data. But the agreement in the $p_T$-dependence of $R_{AA}$ is evidently not perfect (especially for LHC). One sees that the theory somewhat underestimates the slope of the data. And the regions of large $p_T$ support $\alpha_{s}^{fr} = 0.5$ and 0.4 for RHIC and LHC, respectively. One can see that the inclusion of $R_{pp}$ even reduces a little the slope of $R_{AA}$ (since $R_{pp}$ in the denominator on the right hand side of (3) grows with $p_T$). However, this discrepancy
does not seem to be very dramatic since the theoretical uncertainties may be significant.

From Fig. 4 one can see that the effect of $R_{pp}$ on $R_{AA}$ for the central $AA$ collisions can approximately be imitated by simple reduction of the $\alpha_{fr}^{s}$. However, of course, it does not mean that all the theoretical predictions for jet quenching in $AA$ collisions are insensitive to the medium modification of high-$p_T$ spectra in $pp$ collisions. It is clear that the effect of the $R_{pp}$ should be important for $v_2$ and centrality dependence of $R_{AA}$ (simply because in the scenario with the mini-QGP formation in $pp$ collisions the values of $\alpha_{fr}^{s}$ become bigger). It should also be important for the flavor dependence of the the theoretical $R_{AA}$ since the suppression effect for heavy quarks in $pp$ collisions is smaller (by a factor of $\sim 1.5-2$ as our calculations show). In the present exploratory study we do not consider these issues, and leave them for future work.

4.4. Jet quenching in $pA$ collisions

The medium suppression factor $R_{pp}$ should also be taken into account in calculating the nuclear modification factor for $pA$ collisions. Similarly to (3) the correct formula reads $R_{pA} = R_{pA}^{st}/R_{pp}$. Comparison with data on $R_{pA}$ may be even more crucial for the scenario with the formation of the mini-QGP in $pp$ collisions since the sizes and the initial temperatures of the plasma fireballs in $pp$ and $pA$ collisions should not differ strongly. And for this reason the predictions for $R_{pA}$ should not have much uncertainties related to variation of $\alpha_s$ or the temperature dependence of the plasma density and the Debye mass. The ALICE data [31] on $R_{pPb}$ at $\sqrt{s} = 5.02$ TeV exhibit a small deviation from unity at $p_T \gtrsim 10$ GeV, where the Cronin effect should be weak. In
the scenario with the formation of the QGP in $pp$ and $pA$ collisions this is possible only if the magnitudes of the medium suppression in both the processes are very close to each other. Unfortunately, presently we have not data on the UE multiplicity in $pPb$ collisions. However, it is clear that it cannot be smaller than the minimum bias multiplicity density $dN_{ch}^{mb}/d\eta = 16.81 \pm 0.71$ \cite{69}. In principle, it is possible that in the typical minimum bias events the energy deposited in the central rapidity region is already saturated due to a large number of the nucleons which interact with the proton in each $pPb$ collision, and the enhancement of the multiplicity due to jet production is relatively small. The preliminary PHENIX data \cite{39} on the UE in $dAu$ collisions at $\sqrt{s} = 0.2$ TeV really indicate that for dominating small centralities the enhancement of the UE activity as compared to the minimum bias events is relatively small. In order to understand the restrictions on the UE multiplicity density in $pPb$ collisions in the scenario with the mini-QGP formation we simply calculate $R_{pPb}$ for $dN_{ch}/d\eta = K_{ue}dN_{ch}^{mb}/d\eta$ for $K_{ue} = 1, 1.25, \text{ and } 1.5$.

To evaluate $R_{pPb}$ we also need the fireball radius $R_f(pPb)$ which may be bigger than that in $pp$ collisions. In our calculations as a basic choice we use the parametrization of the $R_f(pPb)$ as a function of the multiplicity density given in \cite{38} obtained from the results of simulation of the $pPb$ collisions performed in \cite{21} within the IP-Glasma model \cite{37}. The $R_f(pPb)$ from \cite{37} is close to the $R_f(pp)$ in the region where $R_f(pp) \propto (dN_g/dy)^{1/3}$, but flattens at higher values of the gluon density. Using the parametrization for $R_f(pPb)$ of Ref. \cite{38} and formula (14), we obtained for our set of the enhancement factors for the UE multiplicity $K_{ue} = [1, 1.25, 1.5]$

\begin{equation}
R_f(pPb) \approx [1.63, 1.88, 1.98] \text{ fm} ,
\end{equation}

\begin{equation}
T_0(pPb) \approx [222, 229, 235] \text{ MeV} .
\end{equation}
Comparison of our results with the data on $R_{pPb}$ at $\sqrt{s} = 5.02$ TeV from ALICE [31] is shown in Fig. 5. To illustrate the sensitivity to the fireball size in Fig. 5 we also present the results for the $R_f(pPb)$ 1.2 and 1.4 times greater. As for $AA$ collisions we show the curves with (solid) and without (dashed) the $1/R_{pp}$ factor. Similarly to $R_{AA}$ we account for the nuclear modification of the PDFs with the EKS98 correction [65] (which gives a small deviation of $R_{pPb}$ from unity even without parton energy loss). The results for $R_{pp}$ are also shown (dotted). All the curves are obtained with $\alpha_{fr} = 0.6$. However, our predictions for $R_{pPb}$ (with the $1/R_{pp}$ factor) are quite stable against variation of $\alpha_{fr}$ since the medium suppression is very similar for $pp$ and $pPb$ collisions.

From Fig. 5 one can see that at $p_T \gtrsim 10$ GeV, where the Cronin effect should be small, our predictions (with $1/R_{pp}$ factor) obtained with $K_{ue} = 1$ agree qualitatively with the data. The agreement becomes better with increase of the $R_f(Pb)$. However, similarly to $R_{pp}$ the variation of $R_{pPb}$ with the fireball size is relatively weak. The curves for the higher UE multiplicities ($K_{ue} = 1.25$ and 1.5) lie below the data. Thus, Fig. 5 shows that the data from ALICE [31] may be consistent with the formation of the QGP in $pp$ and $pPb$ collisions if the UE multiplicity is close to the minimum bias one. This condition may be somewhat weakened if the size of the fireball in $pPb$ collisions is considerably bigger than predicted in [21]. But it seems to be rather unrealistic since the required increase of the fireball size is too large. Say, for a good agreement with the data for the UE multiplicity enhancement factor $K_{ue} = 1.5$ one should increase the $R_f(pPb)$ by a factor of $\sim 1.7$.

4.5. A few remarks about approximations and robustness of the results

One remark is in order about the description of the QGP in the ideal gas model in our study. Of course, the QGP temperature formally defined in this model from experimental multiplicity densities is somewhat incorrect. Say, at $T \sim 200$ MeV the ideal gas formula for the entropy underestimates the plasma temperature by $10 - 15\%$ as compared to the entropy from the lattice calculations [44, 45]. However, this fact is practically unimportant for our calculations. Indeed, the potential $v(\rho)$ in the two-dimensional Schrödinger equation, which is used for evaluation of the induced gluon $x$-spectrum in the LCPI approach [3], is proportional to the entropy. Since in our calculations in each case we use the entropy extracted directly from the experimental multiplicity densities, the temperature enters our calculations only through the Debye mass in the dipole cross section. But the latter depends weakly on the Debye mass. For this reason $10 - 15\%$ errors in the temperature can be safely ignored. An accurate definition of the temperature does not make much sense since anyway the nonperturbative effects should modify the form of the potential. Also, one should bear in mind that presently there are many other theoretical uncertainties in the jet quenching calculations, and practically the theory cannot give absolute predictions for the medium suppression. However, one can expect that it can be used to describe the variation of jet quenching from one experimental situation to another (when the parameters of the
model are already fitted to some experimental data). And we do follow this strategy in the present work. We calculate the medium suppression for the small-size plasma in \( pp \) collisions using the information about the values of \( \alpha_s^{fr} \) which are necessary for description of the data on \( R_{AA} \). Without this information one could obtain only a very crude estimate of the effect. Of course, the extrapolation from \( AA \) to \( pp \) collisions assumes that for the real QGP the potential \( v \) is approximately proportional to the entropy, as it is for the ideal QGP. But this assumption seems to be physically very reasonable. Anyway the extrapolation from \( AA \) to \( pp \) collisions should not give large errors since the difference of the plasma temperatures in these two cases is not very big.

It is worth to emphasize that for a reliable extrapolation of the theoretical predictions from \( AA \) to \( pp \) collisions the calculations should be performed with accurate treatment of the LPM suppression (which is very important for \( AA \) collisions) and finite-size and Coulomb effects (which are very important for the mini-QGP produced in \( pp \) collisions). Also, the calculations should be made with running \( \alpha_s \) since the typical virtualities for induced gluon emission in the mini-QGP are higher than that in the large-size QGP in \( AA \) collisions. The LCPI \[^3\] approach used in the present analysis satisfies all these requirements.

Note that, in principle, the assumption that the produced QCD matter exists in the form of an equilibrated QGP is not crucial for our main result that there must be a rather strong jet quenching in \( pp \) collisions. Since even if the created matter is some kind of a hadron resonance gas the parton energy loss will be approximately the same as for the QGP because for a given entropy density the intensities of multiple scattering for the hadron matter and the QGP are very similar \[^70\]. Note that, since the most important quantity, which controls induced gluon emission, is the number density of the color constituents in the medium, from the standpoint of jet quenching, it is even not very important whether the created QCD matter is equilibrated or not. Therefore one can say that in the pQCD picture of jet quenching the significant medium suppression of hadron spectra in \( pp \) collisions is an inevitable consequence of the observed UE multiplicities in \( pp \) collisions and the medium suppression of hadron spectra in \( AA \) collisions (which allows to fix free parameters).

From the point of view of the pQCD the medium suppression of the high-\( p_T \) spectra in \( pp \) collisions may be regarded as a higher twist effect. And of course it would be interesting to observe it through a deviation of the experimental spectra from predictions of the standard pQCD formulas. But it is difficult since the medium suppression should have a very smooth onset in the energy region where the regime of free streaming hadrons transforms to a relatively slow collective expansion of the fireball. Probably it could still occur at \( \sqrt{s} \sim 30 - 40 \) \( \text{GeV} \), where the UE pseudorapidity multiplicity density may be \( \sim 2 - 3 \) and \( T_0 \sim T_c \). For this reason direct observation of this effect by comparing the pQCD predictions with experimental spectra is hardly possible since it is fairly hard to differentiate it from the variations of the theoretical predictions related to small modifications of the PDFs and of the FFs or other higher twist effects not related to the mini-QGP. Also, presently the uncertainties of the pQCD predictions remain large
and the deviation of the ratio data/theory from unity at energies $\sqrt{s} \lesssim 50$ GeV \cite{1, 2, 4, 5} is often considerably bigger than the found medium effects. In this situation it is difficult to identify a relatively small effect from the mini-QGP. Nevertheless, it worth noting that the results of the most recent NLO pQCD analysis of the inclusive charged particle spectra in $pp$ and $\bar{p}p$ collisions at $\sqrt{s} = 0.2 - 7$ TeV \cite{6} show that there is some deviation of the theory from the data that seems to be qualitatively in line with the the scenario with production of the mini-QGP which is more dense at LHC energies. In \cite{6} it was found that the LHC data prefer softer gluon FFs than the lower-$\sqrt{s}$ data. But quantitatively the observed effect is considerably larger than what can be associated with the difference between $R_{pp}$ at RHIC and LHC energies found in the present analysis.

It is worth noting that in principle the preliminary data from ALICE \cite{27} for $\sqrt{s} = 7$ TeV support the existence of jet quenching in $pp$ collisions. These data clearly indicate that the jet fragmentation becomes softer with increase of the UE multiplicity. It is important that the effect is well seen already for the UE $dN_{ch}/d\eta$ smaller than the average one by a factor of $\sim 3$ (i.e., smaller than the typical UE $dN_{ch}/d\eta$ at $\sqrt{s} = 0.2$ TeV). Unfortunately, direct comparison of our calculations with the data \cite{27}) is impossible, since there the NT90 jet observable has been used, which requires more detailed information on the jet structure than our calculations can provide.

5. Summary

Assuming that a mini-QGP may be created in $pp$ collisions, we have evaluated the medium modification factor $R_{pp}$ for light hadrons at RHIC ($\sqrt{s} = 0.2$ TeV) and LHC ($\sqrt{s} = 2.76$ and 7 TeV) energies. We have found an unexpectedly large suppression effect. For $p_T \sim 10$ GeV we obtained $R_{pp} \sim [0.7 - 0.8, 0.65 - 0.75, 0.6 - 0.7]$ at $\sqrt{s} = [0.2, 2.76, 7]$ TeV. We analyzed the role of the $R_{pp}$ in the theoretical predictions for the nuclear modification factor $R_{AA}$ in central $AA$ collisions at RHIC and LHC energies. We found that the presence of $R_{pp}$ does not change dramatically the description of the data on $R_{AA}$ for light hadrons in central $AA$ collisions, and its effect may be imitated by some renormalization of $\alpha_s$. Nevertheless, the effect of the QGP formation in $pp$ collisions may be potentially important in calculating other observables in $AA$ collisions. For example, it should affect $v_2$ and the centrality dependence of $R_{AA}$, and, due to the flavor dependence of $R_{pp}$, its effect may be important for description of the flavor dependence of $R_{AA}$. We leave analysis of these effects for future work. We also calculated the nuclear modification factor $R_{pPb}$ at $\sqrt{s} = 5.02$ TeV. Comparison with the data from ALICE \cite{31} shows that the scenario with the formation of the QGP in $pp$ and $pPb$ collisions may be consistent with the data only if the UE multiplicity density in $pPb$ collisions (which is unknown yet) is close to the minimum bias one.

Acknowledgements

I am grateful to I.P. Lokhtin for discussion of the results. I also thank D.V. Perepelitsa
for useful information. I am indebted to the referee, who pointed out on the recent analysis [76]. This work is supported in part by the grant RFBR 12-02-00063-a.

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