Jet quenching in pp and pA collisions

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We study jet quenching in pp and pA collisions in the scenario with formation of a mini quark-gluon plasma. We find a significant suppression effect. For light hadrons at \( p_T \sim 10 \) GeV we obtained the reduction of the spectra by \( \sim [20 - 30, 25 - 35, 30 - 40] \% \) in pp collisions at \( \sqrt{s} = [0.2, 2.76, 7] \) TeV. We discuss how jet quenching in pp collisions may change the predictions for the nuclear modification factors in AA collisions for light and heavy flavors. We also give predictions for modification of the photon-tagged and inclusive jet fragmentation functions in high multiplicity pp events.

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I. INTRODUCTION

One of the manifestation of the quark-gluon plasma (QGP) formation in AA collisions is the jet quenching phenomenon which is dominated by the radiative parton energy loss \([1-7]\). It leads to suppression of the high-\( p_T \) spectra, which is characterized by the nuclear modification factor \( R_{AA} \) given by the ratio of the inclusive cross section for AA collisions to the binary-scaled inclusive cross section for pp collisions

\[
R_{AA} = \frac{d\sigma(AA \to hX)/dp_Tdy}{N_{bin}\, d\sigma(pp \to hX)/dp_Tdy}.
\]

(1)

It would be extremely interesting to observe jet quenching in pp and pA collisions, since it would be a direct signal of the mini-QGP formation. The QGP formation in pp and pA collisions have been addressed in several publications recently \([8-10]\) from the viewpoint of the hydrodynamical flow effects. In recent papers \([11, 12]\) we studied the possible manifestations of jet quenching in pp collisions within the light-cone path integral approach \([3]\), which we previously used for analysis of jet quenching in AA collisions \([13-16]\). In \([11]\) we discussed the medium modification of the \( \gamma \)-tagged fragmentation functions (FFs) and in \([12]\) the medium modification factor \( R_{pp} \) and its effect on the nuclear modification factors \( R_{AA} \) and \( R_{pA} \). The medium modification factor \( R_{pp} \) characterizes the difference between the real inclusive pp cross section, accounting for the final-state jet interaction in the QGP, and the perturbative one, i.e.,

\[
d\sigma(pp \to hX)/dp_Tdy = R_{pp}\, d\sigma_{pert}(pp \to hX)/dp_Tdy.
\]

(2)

Since we cannot switch off the final state interaction in the QGP, the \( R_{pp} \) is not an observable quantity. Nevertheless, it may affect the theoretical predictions for \( R_{AA} \). Indeed, in the scenario with the QGP formation in pp collisions one should use in the denominator in (1) the real inclusive pp cross section which differs from the perturbative one. In this case one should compare with experimental \( R_{AA} \) the following quantity:

\[
R_{AA} = R_{AA}^{st}/R_{pp},
\]

(3)

where \( R_{AA}^{st} \) is the standard nuclear modification factor calculated using the pQCD predictions for the particle spectrum in pp collisions. The effect of the \( R_{pp} \) may be important for the centrality dependence of \( R_{AA} \) and the azimuthal anisotropy (simply because in the scenario with the QGP formation in pp collisions \( \alpha_s \) becomes bigger). It should also be important for the jet flavor tomography of the QGP \([15-18]\). Because the effect of \( R_{pp} \) on \( R_{AA} \) for heavy quarks should be smaller due to weaker jet quenching for heavy quarks in pp collisions. In this talk I review the results of \([11, 12]\) and extend the analysis \([12]\) to heavy flavors.

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II. MINI-QGP IN PROTON-PROTON COLLISIONS

We describe the mini-QGP fireball within 1+1D Bjorken’s model \([13]\), which gives \(T_0^3 \tau_0 = T^3 \tau\). For \(\tau < \tau_0\) we assume that the medium density \(\propto \tau\). As in our previous analyses of jet quenching in AA collisions \([13,16]\), in the basic variant we take \(\tau_0 = 0.5\) fm. For the QGP in AA collisions with the lifetime/size \(L \gg \tau_0\) the medium effects are not very sensitive to variation of \(\tau_0\). But this may be untrue for \(pp\) collisions when the plasma size is considerably smaller. To understand the sensitivity of \(R_{pp}\) to \(\tau_0\) we also perform calculations for \(\tau_0 = 0.8\) fm. To simplify the computations we neglect variation of the initial temperature \(T_0\) with the transverse coordinates. We fix \(T_0\) using the entropy/multiplicity ratio 

\[ C = \frac{dS/dy}{dN_{ch}/d\eta} \approx 7.67 \text{ obtained in [20]}. \]

The initial entropy density can be written as

\[ s_0 = \frac{C}{\tau_0 \pi R_f^2} \frac{dN_{ch}}{d\eta}, \tag{4} \]

where \(R_f\) is the fireball radius. We ignore the azimuthal anisotropy, and regard \(R_f\) as an effective mini-QGP radius, which includes \(pp\) collisions in the whole range of the impact parameter. This approximation seems to be plausible since the jet production should be dominated by the nearly head-on collisions for which the azimuthal effects should be weak.

In jet quenching calculations for the multiplicity density in \([4]\) one should use the multiplicity density of the soft (underlying-event (UE)) hadrons, which is bigger than the minimum bias multiplicity density by a factor \((K_{ue})\) of \(\sim 2\) \([21]\). Experimental studies \([21,25]\) show that the UE multiplicity grows with momentum of the leading charged jet hadron at \(p_T \lesssim 3–5\) GeV and then flattens out. The plateau region corresponds approximately to \(E_{jet} \gtrsim 15–20\) GeV. To fix the \(dN_{ch}/d\eta\) in \(\sqrt{s} = 0.2\) TeV we use the UE enhancement factor \(K_{ue}\) from PHENIX \([22]\) obtained by dihadron correlation method. Taking for minimum bias non-diffractive events \(dN_{ch}^{mb}/d\eta = 2.98 \pm 0.34\) from STAR data \([20]\), we obtained for the UEs in the plateau region \(dN_{ch}/d\eta \approx 6.5\). To evaluate the UE multiplicity at \(\sqrt{s} = 2.76\) and 5.02 TeV we use the data from ATLAS \([23]\) at \(\sqrt{s} = 0.9\) and 7 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 \(\sqrt{s} = 0.9\) TeV and 7 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 TeV, respectively. We use for \(R_f\) the values obtained in numerical simulations of \(pp\) collisions at \(\sqrt{s} = 7\) TeV performed in \([10]\) within the IP-Glasma model \([27]\). In \([10]\) it has been found that \(R_f\) grows approximately as linear function of \((dN_g/dy)^{1/3}\) and then flattens out (a convenient parametrization of \(R_f\) from \([10]\) has been given in \([28]\)). The plateau region corresponds to nearly head-on collisions where the fluctuations of multiplicity are dominated by the fluctuations of the plasma color fields \([10]\). With the help of the formula for \(R_f\) from \([28]\) for the above values of the UE multiplicity densities in the plateau regions we obtain (we take \(dN_g/dy = \kappa dN_{ch}/d\eta\) with \(\kappa = C45/2\pi^{4/3}(3) \approx 2.13\))

\[ R_f[\sqrt{s} = 0.2, 2.76, 5.02, 7\, \text{TeV}] \approx [1.3, 1.44, 1.49, 1.51] \text{ fm}. \tag{5} \]

We neglect possible variation of the \(R_f\) from RHIC to LHC since our results are not very sensitive to \(R_f\). Using \([4]\) 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

\[ T_0[\sqrt{s} = 0.2, 2.76, 5.02, 7\, \text{TeV}] \approx [199, 217, 226, 232] \text{ MeV}. \tag{6} \]

One sees that the values of \(T_0\) lie well above the deconfinement temperature \(T_c \approx 160 – 170\) MeV.

For initial temperatures \(\tau > \tau_{QGP}\) the purely plasma phase may exist up to \(\tau_{QGP} \sim 1–1.5\) fm. At \(\tau > \tau_{QGP}\) the hot QCD matter will evolve in the mixed phase up to \(\tau_{max} \sim 2R_f\) where the transverse expansion should lead to fast cooling of the fireball. For \(\tau_{QGP} < \tau < \tau_{max}\) the QGP fraction in the mixed phase is approximately \(\propto 1/\tau\) \([18]\), and for this reason we can use \(1/\tau\) 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 \(\tau > \tau_{QGP}\)).

The central question for the scenario with mini-QGP formation is the extend to which the mini-fireball created in \(pp\) collisions may be treated as a continuous macroscopic medium. This question at present is still open. The lattice studies support the idea that a collective medium may be created in \(pp\) collisions. Indeed, the macroscopic behavior of the fireball is possible when the Knudsen number \(Kn \sim \tau_c/\tau\) is small. We estimated \(Kn\) using the recent lattice results \([29]\) on the electric conductivity \(\sigma\) of the QGP.
From the Drude formula (for massless partons)

\[ \sigma \sim \frac{(e^2)^n}{3T} \]

and lattice \( \sigma \) from [29] we obtained approximately for the temperatures given in [9] \( K_n(\text{quark}) \sim 1 \) at \( \tau \sim 0.5 \text{ fm} \) and \( K_n(\text{quark}) \sim 0.25 \) at \( \tau \sim 1 \text{ fm} \). The gluon Knudsen number should be smaller by a factor of \( \sim C_F/C_A = 4/9 \). This qualitative analysis shows that the collective behavior of the mini-fireball does not seem to be unrealistic. Of course, the inequality \( K_n \ll 1 \) is just a necessary condition for the hydrodynamic behavior of the QGP. But it cannot guarantee that the QGP is produced quickly after pp collision.

III. MEDIUM INDUCED GLUON SPECTRUM AND PARAMETERS OF THE MODEL

As in [13], 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 analysis [30] 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 plasma quasiparticle masses of light quarks and gluon we take \( m_q = 300 \) and \( m_g = 400 \text{ MeV} \) supported by the analysis of the lattice data [31]. Our results are not very sensitive to \( m_g \), and practically insensitive to the value of \( m_q \). For gluon emission from a quark (or gluon) the \( x \)-spectrum may be written [32] through the light-cone wave function of the \( gg \) (or \( ggg \)) system in the coordinate \( \rho \)-representation. Its \( z \)-dependence 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 \( \sigma(\rho) \) is proportional to the QGP density/entropy times a linear combination of the dipole cross sections \( \sigma(\rho) \), \( \sigma(1-x)\rho \) and \( \sigma(x\rho) \). We perform calculations with running \( \alpha_s \) frozen at some value \( \alpha_s^f \) at low momenta. For gluon emission in vacuum a reasonable choice is \( \alpha_s^f \sim 0.7 \sim 0.8 \) [33, 34]. In plasma thermal effects can suppress \( \alpha_s^f \). However, the uncertainties of jet quenching calculations are large and the extrapolation from the vacuum gluon emission to the induced radiation may be unreliable. For this reason we treat \( \alpha_s^f \) as a free parameter of the model. In [16] we have observed that data on \( R_{AA} \) are consistent with \( \alpha_s^f \sim 0.5 \) for RHIC and \( \alpha_s^f \sim 0.4 \) for LHC. The reduction of \( \alpha_s^f \) from RHIC to LHC may be due to stronger thermal effects at LHC where the initial temperature is bigger. But the analysis [16] is performed ignoring the medium suppression in pp collisions. Accounting for \( R_{pp} \) should increase \( \alpha_s^f \). However, in [16] we used the plasma density vanishing at \( \tau < \tau_0 \), whereas now we use the QGP density \( \propto \tau \), which leads to somewhat stronger medium suppression. As a result, preferable \( \alpha_s^f \) (from the standpoint of the description of \( R_{AA} \)) remains approximately the same, or a bit larger, as obtained in [16]. If the difference between \( \alpha_s^f \) 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 [10] a reasonable window is \( \alpha_s^f \sim 0.6 \sim 0.7 \). In principle for the mini-QGP the thermal reduction of \( \alpha_s \) may be smaller than for the large-size plasma (at the same temperature). Because for the mini-QGP a considerable contribution to the induced gluon emission comes from the product 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. We perform the calculations for \( \alpha_s^f \sim 0.5, 0.6 \) and 0.7. Note that \( R_{pp} \) should be less sensitive to \( \alpha_s^f \) than \( R_{AA} \) since the typical virtualities for induced gluon emission in the mini-QGP are larger than that in the large-size QGP (see below).

The physical pattern of induced gluon emission in the mini-QGP differs somewhat from that for the large-size QGP. For the mini-QGP when the typical path length in the medium \( L \sim 1 \sim 1.5 \text{ fm} \) the energy loss is dominated by gluons with \( L_f \lesssim L \), where \( L_f \sim 2\omega/m_g^2 \) is the gluon formation length in the low density limit. In this regime the dominating contribution comes from the \( N = 1 \) rescattering, and the finite-size and Coulomb effects play a crucial role [32, 36] (see also [37]). On the contrary, for the QGP in \( AA \) collisions the induced energy loss is dominated by gluons with \( L_f \lesssim L \). Indeed, \( L_f \sim 2\omega S_{LPM}/m_g^2 \), where \( S_{LPM} \) is the LPM suppression factor. For RHIC and LHC typically \( S_{LPM} \sim 0.3 \sim 0.5 \) for \( \omega \sim 2 \text{ GeV} \), it gives \( L_f \sim 1.5 \sim 2.5 \text{ fm} \) 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 the gluon spectrum is (locally) approximately similar to that in an infinite extent matter. It is important that the induced gluon emission in the mini-QGP is more perturbative than in the large-size QGP. Indeed, from the Schrödinger diffusion relation one can obtain for the 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 $Q^2 \sim 10/\rho^2$ [39], we obtain $Q^2 \sim 5\omega/\xi$. For $\omega \sim 2$ and $\xi \sim 0.5-1$ fm it gives $Q^2 \sim 2-4$ GeV$^2$. The virtuality scale in the gluon emission vertex has a similar form but smaller by a factor of $\sim 2.5$ [39]. The $1/\xi$ dependence of $Q^2$ persists up to $\xi \sim L_f$. For the large-size QGP 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 in AA collisions.

IV. ENERGY LOSS IN THE MINI-QGP

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 $\alpha_{frs} = 0.6$ (as in [39], both the radiative and collisional contributions are defined for the lost energy smaller than half of the initial parton energy). We present results for the fireball parameters obtained for the jet energy dependent UE $dN_{ch}/d\eta$ and for that in the plateau region (details see in [12]). One can see that the energy loss for these two versions (solid and long-dashed lines) become very close to each other at $E \gtrsim 10$ GeV. Our results show that 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 is smaller than that for the large-size QGP in AA collisions obtained in [16] by a factor of $\sim 4$.

FIG. 1: Energy dependence of the energy loss of gluons (left) and light quarks (right) produced in the center of the mini-QGP fireball at $\sqrt{s} = 0.2$ TeV (upper panels) and $\sqrt{s} = 2.76$ TeV (lower panels). 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 (5) and (6). All the curves are for $\alpha_{frs} = 0.6$.

In Fig. 2 we show the 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 illustrate the difference between $pp$ and $AA$ collisions we present also predictions for radiative energy loss for $T_0 = 320$ 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. We rescaled the predictions for $AA$ collisions by the factor $(T_0(pp)/T_0(AA))^3$. One sees that at $L \geq \tau_0$ the radiative energy loss is approximately a linear function of $L$, and at $L < \tau_0$ the radiative energy loss is approximately $\propto L^3$ (since the leading $N = 1$ rescattering term to the effective Bethe-Heitler cross section is $\propto L$ [33, 36] and integration over the longitudinal coordinate of the scattering center gives additional two powers of $L$). From comparison of the radiative
energy loss for $T_0 = 199$ and 217 MeV to that for $T_0 = 320$ and 420 MeV one can see a deviation from
the $T^3$ scaling by factors of $\sim 1.5$ and $\sim 2$, respectively. This difference persists even at $L \sim 1$ fm. It
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.

V. MEDIUM MODIFICATION OF THE INCLUSIVE SPECTRA

A. 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 $D_{h/i}$

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

where $d\sigma (pp \to iX)/dp_i^T dy$ is the ordinary hard cross section, $p_i^T = p_T/z$ is the parton transverse
momentum. We write the real inclusive cross section in a similar form but with the medium modified FF
$D_{h/i}^m$

$$\frac{d\sigma (pp \to hX)}{dp_T dy} = \sum_i \int_0^1 \frac{dz}{z^2} D_{h/i}^m(z, \alpha_s cQ) \frac{d\sigma (pp \to iX)}{dp_i^T dy}.$$
for details). Note that the permutation of the DGLAP and the induced stages gives a very small effect [13].

Since we ignore the azimuthal effects, the averaging of the medium modified FFs over the geometrical variables of the hard parton process and over the impact parameter of \(pp\) collision is simply reduced to averaging over the parton path length \(L\) in the QGP. We have performed averaging over \(L\) for 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 have approximately the same distribution in the transverse spacial coordinates). We obtained that practically in the full range of the \(pp\) impact parameter the distribution in \(L\) is sharply peaked around \(L \approx \sqrt{S_{ov}}/\pi\) (here \(S_{ov}\) is the overlap area for two colliding bags). It shows that \(R_f\) at the same time gives the typical path length for fast partons. We found that, as compared to \(L = R_f\), the \(L\)-fluctuations reduce the medium modification by only \(\sim 10 - 15\%\).

We treat the collisional energy loss, which is relatively small [39], as a small perturbation to the radiative mechanism, and incorporate it simply by renormalizing the QGP temperature in calculating the medium modified FFs for the induced radiation (see [13] for details).

B. Predictions for \(R_{pp}\)

In Fig. 3 we present the results for \(R_{pp}\) of charged hadrons at \(\sqrt{s} = 0.2\), 2.76 and 7 TeV for \(\alpha_s^{fr} = 0.5\), 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. 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. One can see that for \(\tau_0 = 0.8\) fm the reduction of the suppression is not very significant. Fig. 3 shows that, as we expected, \(R_{pp}\) does not exhibits a strong dependence on \(\alpha_s^{fr}\). 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. In the left part of Fig. 4 we show a comparison between \(R_{pp}\) at \(\sqrt{s} = 7\) TeV for the minimum bias and the UE \(dN_{ch}/d\eta\). One can see that even the minimum bias \(dN_{ch}/d\eta\) gives a considerable suppression. The right part of Fig. 4 shows variation of \(R_{pp}\) between \(\sqrt{s} = 7\) and 100 TeV. One sees that the energy dependence of \(R_{pp}\) is weak.

![Fig. 3: \(R_{pp}\) of charged hadrons at \(\sqrt{s} = 0.2\) (a), 2.76 (b), 7 (c) TeV for (top to bottom) \(\alpha_s^{fr} = 0.5\), 0.6 and 0.7 for \(\tau_0 = 0.5\) (solid) and 0.8 (dashed) fm.](image-url)

To study the sensitivity of \(R_{pp}\) to the fireball radius we also performed the calculations for \(R_f\) given by [5] times 0.7 and 1.3. We found that in these two cases the medium suppression is smaller by \(\sim 3\%\) and 10\%, respectively. The weak dependence on \(R_f\) is due to a compensation between the enhancement of the energy loss caused by increase of the fireball size and its suppression due to reduction of the QGP density. Note that the stability of \(R_{pp}\) against variations of \(R_f\) shows that the variation of the plasma density in the transverse coordinates should not be very important. Indeed, the gluon spectrum is dominated by \(N = 1\) rescattering term which is a linear functional of the plasma density profile along the fast parton trajectory. Therefore the energy loss for a more realistic plasma density (with a higher density...
in the central region) can be roughly approximated by a linear superposition of that for the step density distributions with different $R_f$. And it should not change strongly $R_{pp}$ as compared to our calculations.

Fig. 3 shows the results for the typical UE multiplicity density. An accurate accounting for the fluctuations of the UE $dN_{ch}/d\eta$ is impossible since it should be done on the event-by-event basis, and requires detailed information about dynamics of the UEs. To understand how the event-by-event fluctuations of the UE $dN_{ch}/d\eta$ may change our results, we evaluated $R_{pp}$ assuming that the distribution in $dN_{ch}/d\eta$ is the same at each impact parameter and jet production point. We used the distribution in $dN_{ch}/d\eta$ from CMS [24] measured at $\sqrt{s} = 0.9$ and 7 TeV. It satisfies approximately KNO scaling similar to that in minimum bias events [44]. For this reason one can expect that it can be used for RHIC conditions as well. We observed that the fluctuating $dN_{ch}/d\eta$ suppresses $(1 - R_{pp})$ by only $\sim 5 - 6\%$ both for RHIC and LHC energies. This says that our approximation without the event-by-event fluctuations of the QGP parameters should be good.

C. Effect of $R_{pp}$ on $R_{AA}$

To illustrate the effect of the mini-QGP in pp collisions on $R_{AA}$ in Fig. 5 we compare 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 [45], and with the data for charged hadrons in central Pb + Pb collisions at $\sqrt{s} = 2.76$ TeV (b,c) from ALICE [46] and CMS [47]. We show the predictions for $R_{AA}$ defined by (3) with (red) the $1/R_{pp}$ factor in (3). The red curves are obtained with the factor $1/R_{pp}$ calculated with $\alpha_f = 0.6$. Data points are from PHENIX [45] (a), ALICE [46] (b) and CMS [47] (c). Systematic experimental errors are shown as shaded areas.
FIG. 6: $R_{AA}$ of charged particles vs $N_{\text{part}}$ for $Pb + Pb$ at $\sqrt{s} = 2.76$ TeV with (red) and without (blue) $R_{pp}$, for (top to bottom) $\alpha^{fr}_{s} = 0.4$ and $0.5$ for $\sqrt{s} = 2.76$ TeV, $R_{pp}$ is calculated at $\alpha^{fr}_{s} = 0.6$. Data points are from ALICE [52].

factor, and for $R_{AA}^{st}$ without (blue) this factor. We use the $R_{pp}$ for $\alpha^{fr}_{s} = 0.6$. We calculated $R_{AA}^{st}$ for $\alpha^{fr}_{s} = 0.5$ and $0.6$ at $\sqrt{s} = 0.2$ TeV, and for $\alpha^{fr}_{s} = 0.4$ and $0.5$ at $\sqrt{s} = 2.76$ TeV. Because these values give better agreement with the data. We accounted for the nuclear modification of the PDFs with the EKS98 correction [48]. As in [10], 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 [49] and LHC [50, 51]. 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 $\alpha^{fr}_{s} = 0.6$, and for LHC the value $\alpha^{fr}_{s} = 0.5$ seems to be preferred by the data. The agreement in the $p_T$-dependence of $R_{AA}$ is not perfect (especially for LHC). The theory somewhat underestimates the slope of the data. It seems that the regions of large $p_T$ support $\alpha^{fr}_{s} = 0.5$ and 0.4 for RHIC and LHC, respectively. The inclusion of $R_{pp}$ even reduces a little the slope of $R_{AA}$. However, it does not seem to be very dramatic since the theoretical uncertainties may be significant.

Fig. 5 shows that the effect of $R_{pp}$ on $R_{AA}$ in central $AA$ collisions can approximately be imitated by a simple reduction of $\alpha^{fr}_{s}$. However, it is clear that $R_{pp}$ may be important for the azimuthal effects and the centrality dependence of $R_{AA}$ since in the scenario with the mini-QGP formation in $pp$ collisions the values of $\alpha^{fr}_{s}$ become bigger. The effect of $R_{pp}$ on the centrality dependence of $R_{AA}$ is shown Fig. 6. $R_{pp}$ can also affect the flavor dependence of $R_{AA}$ since the suppression effect for heavy quarks in $pp$ collisions is smaller. It is illustrated in Figs. 7–9 for the $p_T$-dependence of the ratio of the $R_{AA}$ for heavy and light flavors. One sees that at $p_T \lesssim 10$ GeV $R_{pp}$ reduces the difference between the nuclear suppression of the

FIG. 7: Effect of $R_{pp}$ due to mini-QGP on ratio $R_{AA}$ for $D$-mesons to $R_{AA}$ for light charged hadrons. $\alpha^{fr}_{s} = 0.6$ for $\sqrt{s} = 0.2$ TeV and $\alpha^{fr}_{s} = 0.5$ for $\sqrt{s} = 2.76$ TeV, $R_{pp}$ for light and heavy flavors is calculated at $\alpha^{fr}_{s} = 0.6$. 
spectra for heavy and light flavors. In Fig. 10 we show the effect of $R_{pp}$ on the centrality dependence of $R_{AA}$ for $D$-mesons. One can see that $R_{pp}$ may improve somewhat agreement with the data.

FIG. 9: Effect of $R_{pp}$ due to mini-QGP on ratio $R_{AA}$ for non-photonic electrons to $R_{AA}$ for light charged hadrons. $\alpha_s^{fr} = 0.6$ for $\sqrt{s} = 0.2$ TeV and $\alpha_s^{fr} = 0.5$ for $\sqrt{s} = 2.76$ TeV, $R_{pp}$ for light and heavy flavors is calculated at $\alpha_s^{fr} = 0.6$.

FIG. 10: $R_{AA}$ of $D$-mesons vs $N_{part}$ for $Pb + Pb$ at $\sqrt{s} = 2.76$ TeV with (red) and without (blue) $R_{pp}$, for (top to bottom) $\alpha_s^{fr} = 0.4$ and $\alpha_s^{fr} = 0.5$, $R_{pp}$ is calculated at $\alpha_s^{fr} = 0.6$. Data points are from ALICE [53].
D. Jet quenching in $pA$ collisions

In the scenario with the QGP production in $pp$ collisions the correct formula for $R_{pA}$ reads $R_{pA} = R_{pA}^{\text{et}}/R_{pp}$. Evidently, the sizes and the initial temperatures of the plasma fireballs in $pp$ and $pA$ collisions should not differ strongly. For this reason for $R_{pA}$ the uncertainties related to variation of $\alpha_s$ (or the temperature dependence of the QGP density and the Debye mass) are smaller than for $R_{AA}$. The ALICE data [54] show a small deviation from unity of the bias multiplicity density $dN_{ch}/d\eta$ for $pA$ data [55]. In order to understand the acceptable range of the $dN_{ch}/d\eta$ in $pA$ collisions we calculated $R_{pA}$ for $dN_{ch}/d\eta = K_{\text{ue}} dN_{ch}^{mb}/d\eta$ for $K_{\text{ue}} = 1, 1.25, \text{ and } 1.5$.

In our calculations as a basic choice we use the parametrization of $R_f(pPb)$ vs the multiplicity given in [28] obtained from the results of simulation of the $pPb$ collisions performed in [10] within the IP-Glasma model [27]. Ref. [27] gives $R_f(pPb)$ that is close to $R_f(pp)$ where $R_f(pp) \propto (dN_2/d\eta)^{1/3}$, but $R_f(pPb)$ flattens at higher values of the gluon density. Using formula (4), we obtained for $K_{\text{ue}} = [1, 1.25, 1.5]$}

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

$$T_0(pPb) \approx [222, 229, 235] \text{ MeV}.$$ (10)

Fig. 11 shows comparison of our results with the data on $R_{pPb}$ at $\sqrt{s} = 5.02 \text{ TeV}$ from ALICE [54]. To illustrate the sensitivity to $R_f(pPb)$ we also present the results for $R_f(pPb)$ 1.2 and 1.4 times greater. We show the curves with (red) and without (blue) the $1/R_{pp}$ factor. As for $AA$ case we account for the nuclear modification of the PDFs with the EKS98 correction [48]. It gives a small deviation of $R_{pPb}$ from unity even without parton energy loss. The results for $R_{pp}$ are also shown (green). All the curves are obtained with $\alpha_f = 0.6$. However, our predictions for $R_{pPb}$ (with the $1/R_{pp}$ factor) are quite stable against variation of $\alpha_f$ since the medium effects are very similar for $pp$ and $pPb$ collisions.

![FIG. 11: (a) $R_{pPb}$ for charged hadrons at $\sqrt{s} = 5.02 \text{ TeV}$ from our calculations for $\alpha_f = 0.6$ with (red) and without (blue) the $1/R_{pp}$ factor for (top to bottom) $K_{\text{ue}} = 1, 1.25$ and 1.5 for the $R_f(pPb)$ from (10). (b,c) same as (a) but for $R_f(pPb)$ times 1.2 and 1.4. The green line shows $R_{pp}$. The dot-dashed line shows $R_{pPb}$ due to the EKS98 correction [48] to the nuclear PDFs. Data points are from ALICE [54].](image-url)

Fig. 11 shows that at $p_T > 10 \text{ GeV}$, where the Cronin effect should be small, our predictions (with $1/R_{pp}$ factor) obtained with $K_{\text{ue}} = 1$ agree qualitatively with the data. The agreement becomes better for larger $R_f(Pb)$. But just as for $R_{pp}$ the variation of $R_{pPb}$ with the fireball size is relatively weak. The curves for the higher UE multiplicities ($K_{\text{ue}} = 1.25$ and 1.5) lie below the data. Thus we see that the data from ALICE [54] 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 weakened if the size of the
fireball in $pPb$ collisions is considerably bigger than predicted in [10]. But the physical picture may change if we take into account the meson-baryon Fock component in the proton. Indeed, in $pA$ collisions the final-state interaction may be smaller due to meson-baryon Fock component in the proton. The weight of the $MB$-component may be as large as $\sim 40\%$ [54]. Contrary to $pp$ case in $pA$ collisions practically in all events meson should produce its own fireball. It means that in $\sim 40\%$ events an asymmetric two-fireball configuration may be produced (as illustrated in the left part of Fig. 12). Since jet may propagate without interaction with one of the fireball (typically it is the meson fireball as shown in the right part of Fig. 12), the final-state interaction should be weaker than for a symmetric fireball (for same $dN_{ch}/d\eta$). Note that the two-fireball state naturally generates the azimuthal flow for the soft particles as well.

FIG. 12: A cartoon of the production of a two-fireball state in $pA$ collisions from the meson-baryon Fock component of the proton (a); A cartoon of the jet quenching for the two-fireball state (b).

VI. MEDIUM MODIFICATION OF PHOTON-TAGGED AND INCLUSIVE JETS IN HIGH-MULTIPLICITY PROTON-PROTON COLLISIONS

For a direct observation of the medium effects in $pp$ collisions one can use measurement of the jet FF in $\gamma+\text{jet}$ events for different UE multiplicities. To understand the prospects of this method we evaluate the medium modification of the $\gamma$-tagged FF at $\sqrt{s} = 7$ TeV at $y = 0$. The values of the $R_f$ and $T_0$ for different values of $dN_{ch}/d\eta$ obtained using (4) are given in Table I. For $dN_{ch}/d\eta \gtrsim 40$ we obtain $T_0$ which is about that for central $Au + Au$ collisions at RHIC.

| $dN_{ch}/d\eta$ | 3 | 6 | 20 | 40 | 60 |
|------------------|---|---|----|----|----|
| $R_f$ (fm)       | 1.046 | 1.27 | 1.538 | 1.538 | 1.538 |
| $T_0$ (MeV)      | 177 | 196 | 258 | 325 | 372 |

In $\gamma+\text{jet}$ events the energy of the hard parton, $E_T$, in the direction opposite to the tagged photon is smeared around the photon energy, $E_\gamma^T$. But using the results of the NLO calculations [57] one can show that at $E_\gamma^T \gtrsim 25$ GeV and $z \lesssim 0.9$ the smearing can be safely neglected (for details, see [11]). To be conservative we present results for $z < 0.8$, where the effect of smearing is practically negligible and one can set $E_T = E_\gamma^T$. Then, as in [58], we can write the $\gamma$-tagged FF as a function of the UE multiplicity density $dN_{ch}/d\eta$ (for clarity we denote it by $N$) as

$$D_h(z, E_\gamma^T, N) = \left\langle \left\langle \sum_i r_i(E_\gamma^T) D_{h/i}^m(z, E_\gamma^T, N) \right\rangle \right\rangle,$$

where, as in [9], $D_{h/i}^m$ is the medium modified FF for $i \rightarrow h$ process, and $r_i$ is the fraction of the $\gamma + i$ parton state in the $\gamma+\text{jet}$ events, $\left\langle \left\langle \ldots \right\rangle \right\rangle$ means averaging over the transverse geometrical variables of $pp$ collision and jet production, which includes averaging over the fast parton path length $L$ in the QGP.
Just as for $R_{pp}$ 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. As compared to $L = R_f$ the $L$-fluctuations reduce the medium modification by $\sim 10 – 15\%$. In Fig. 13 we present the results for the medium modification factor (for charged hadrons)

$$I_{pp}(z, E_T, N) = \frac{D_h(z, E_T, N)}{D_h^{\text{vac}}(z, E_T)}$$

for the $\gamma$-tagged (upper panels) jets for $E_T = [25, 50, 100]$ GeV at $\sqrt{s} = 7$ TeV. For comparison we also show the results for inclusive (lower panels) jets. The smearing effect is irrelevant to inclusive jets and we show the results for the whole range of $z$. For illustration of the difference between $pp$ and $AA$ collisions we also present the curves for $\sqrt{s} = 2.76$ TeV for $L = 5$ fm and $T_0 = 420$ MeV that can be regarded as reasonable values for $Pb + Pb$ collisions (we used $\alpha_s f_r = 0.4$, which is favored by the data on $R_{AA}(p_T)$ at $p_T \gtrsim 20$ GeV). Fig. 13 shows that there is a considerable quenching effect for $dN_{ch}/d\eta > 20$. Note that the observed strong quenching of inclusive jets is qualitatively supported by the preliminary data from ALICE [59] that indicate that for the high multiplicity UEs jets undergo a softer fragmentation.

Since the vacuum FFs are unobservable, in practice, to observe the medium effect one should simply compare the FFs for different multiplicities. In Fig. 13 we show the ratio of the FFs for $N = 40$ and $N = 3$ (for inclusive jets this ratio cannot be measured, and we show it just to illustrate the difference in magnitudes of the effect for $\gamma$-tagged and inclusive jets). As for $R_{pp}$ we have investigated the sensitivity of our results to variation of $R_f$, and found that $I_{pp}$ is quite stable against variation of $R_f$.

VII. SUMMARY

Assuming that a mini-QGP fireball may be created in $pp$ collisions, we have evaluated the medium modification of high-$p_T$ particle spectra for light and heavy flavors and medium modification factors for the $\gamma$-triggered and inclusive jet FFs. 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 have studied the effect of $R_{pp}$ on the theoretical predictions for the nuclear modification factor $R_{AA}$ in $AA$ collisions at RHIC and LHC energies. We found that $R_{pp}$ 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 $\alpha_s$. But inclusion of $R_{pp}$ changes the centrality dependence of $R_{AA}$. Also, $R_{pp}$ weakens the flavor dependence of $R_{AA}$.

Our results show that the ALICE data on $R_{Pp}$ may be consistent with the scenario with the QGP formation if in $pPb$ collisions the UE multiplicity is close to the minimum bias one. But this condition may be weakened due to presence in the proton wave function of the meson-baryon Fock component. We leave analysis of its effect for future work.
We demonstrated that in \( pp \) collisions with UE multiplicity density \( dN_{ch}/d\eta \sim 20 - 40 \) the mini-QGP can suppress the \( \gamma \)-triggered FF at \( E_T \sim 25 - 100 \) GeV and \( z \sim 0.5 - 0.8 \) by \( \sim 10 - 40\% \), and for inclusive jets the effect is even stronger.

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[1] M. Gyulassy and X.N. Wang, Nucl. Phys. B420, 583 (1994).
[2] R. Baier, Y.L. Dokshitzer, A.H. Mueller, S. Peigné, and D. Schiff, Nucl. Phys. B483, 291 (1997); ibid. B484, 265 (1997).
[3] B.G. Zakharov, JETP Lett. 63, 952 (1996); ibid 65, 615 (1997); 70, 176 (1999); Phys. Atom. Nucl. 61, 838 (1998).
[4] R. Baier, D. Schiff, and B.G. Zakharov, Ann. Rev. Nucl. Part. 50, 37 (2000).
[5] U.A. Wiedemann, Nucl. Phys. A690, 731 (2001).
[6] M. Gyulassy, P. Lévai, and I. Vitev, Nucl. Phys. B594, 371 (2001).
[7] P. Arnold, G.D. Moore, and L.G. Yaffe, JHEP 0206, 030 (2002).
[8] P. Bozek, Acta Phys. Polon. B41, 837 (2010).
[9] J. Casalderrey-Solana and U.A. Wiedemann, Phys. Rev. Lett. 104, 102301 (2010).
[10] A. Bzdak, B. Schenke, P. Tribedy, and R. Venugopalan, arXiv:1304.3403.
[11] B.G. Zakharov, Phys. Rev. Lett. 112, 032301 (2014).
[12] B.G. Zakharov, J. Phys. G41, 075008 (2014).
[13] B.G. Zakharov, JETP Lett. 88, 781 (2008).
[14] B.G. Zakharov, JETP Lett. 93, 683 (2011).
[15] B.G. Zakharov, JETP Lett. 96, 616 (2013).
[16] B.G. Zakharov, J. Phys. G40, 085003 (2013).
[17] N. Armesto, M. Cacciari, A. Dainese, C.A. Salgado, and U.A. Wiedemann, Phys. Lett. B637, 362 (2006).
[18] A. Buzzatti and M. Gyulassy, Phys. Rev. Lett. 108, 022301 (2012).
[19] J.D. Bjorken, Phys. Rev. D27, 140 (1983).
[20] B. Müller and K. Rajagopal, Eur. Phys. J. C43, 15 (2005).
[21] A.A. Affolder et al. [CDF Collaboration], Phys. Rev. D65, 092002 (2002).
[22] J. Jia, for the PHENIX Collaboration, contribution to the Quark Matter 2009 Conf., March 30 - April 4, Knoxville, Tennessee; arXiv:0906.3776.
[23] G. Aad et al. [ATLAS Collaboration], Phys. Rev. D83, 112001 (2011).
[24] S. Chatrchyan et al. [CMS Collaboration], JHEP 1109, 109 (2011).
[25] B. Abelev et al. [ALICE Collaboration] JHEP 1207, 116 (2012).
[26] B.I. Abelev et al. [STAR Collaboration], Phys. Rev. C79, 034909 (2009).
[27] B. Schenke, P. Tribedy, and R. Venugopalan, Phys. Rev. Lett. 108, 252301 (2012).
[28] L. McLerran, M. Praszalowicz, and B. Schenke, arXiv:1306.2350.
[29] A. Amato, G. Aarts, C. Allton, P. Giudice, S. Hands, and J.-I. Skullerud, arXiv:1310.7466.
[30] O. Kaczmarek and F. Zantow, Phys. Rev. D71, 114510 (2005).
[31] P. Lévai and U. Heinz, Phys. Rev. C57, 1879 (1998).
[32] B.G. Zakharov, JETP Lett. 80, 617 (2004).
[33] N.N. Nikolaev and B.G. Zakharov, Phys. Lett. B327, 149 (1994).
[34] Yu.L. Dokshitzer, V.A. Khoze, and S.I. Troyan, Phys. Rev. D53, 89 (1996).
[35] B.G. Zakharov, JETP Lett. 73, 49 (2001).
[36] P. Aurenche and B.G. Zakharov, JETP Lett. 90, 237 (2009) arXiv:0907.1918.
[37] P. Arnold, Phys. Rev. D80, 025004 (2009).
[38] N.N. Nikolaev and B.G. Zakharov, Phys. Lett. B332, 184 (1994).
[39] B.G. Zakharov, JETP Lett. 86, 444 (2007).
[40] S. Kretzer, H.L. Lai, F. Olness, and W.K. Tung, Phys. Rev. D69, 114005 (2004).
[41] T. Sjostrand, L. Lonnblad, S. Mrenna, and P. Skands, arXiv:hep-ph/0308153.
[42] B.A. Kniehl, G. Kramer, and B. Potter, Nucl. Phys. B582, 514 (2000).
[43] R. Baier, Yu.L. Dokshitzer, A.H. Mueller, and D. Schiff, JHEP 0109, 033 (2001).
[44] A. Dumitru and E. Petreska, arXiv:1209.4105.
[45] A. Adare et al. [PHENIX Collaboration], arXiv:1208.2254.
[46] B. Abelev et al. [ALICE Collaboration], Phys. Lett. B720, 52 (2013).
[47] S. Chatrchyan et al. [CMS Collaboration], Eur. Phys. J. C72, 1945 (2012).
[48] K.J. Eskola, V.J. Kolhinen, and C.A. Salgado, Eur. Phys. J. C9, 61 (1999).
[49] B.I. Abelev et al. [STAR Collaboration], Phys. Rev. C79, 034909 (2009).
[50] S. Chatrchyan et al. [CMS Collaboration], JHEP 1108, 141 (2011).
[51] K. Aamodt et al. [ALICE Collaboration], Phys. Rev. Lett. 106, 032301 (2011).
[52] B. Abelev et al. [ALICE Collaboration], Phys. Lett. B720, 52 (2013).
[53] E. Bruna, for the ALICE Collaboration, contribution to 14th International Conference on Strangeness in Quark Matter (SQM2013), J. Phys. Conf. Ser. 509, 012080 (2014) arXiv:1401.1698.
[54] B. Abelev et al. [ALICE Collaboration], Phys. Rev. Lett. 110, 082302 (2013).
[55] B. Abelev et al. [ALICE Collaboration], Phys. Rev. Lett. 110, 032301 (2013).
[56] J. Speth, A.W. Thomas, Adv. Nucl. Phys. 24, 83 (1997).
[57] H. Zhang, J.F. Owens, E. Wang, and X.-N. Wang, Phys. Rev. Lett. 103, 032302 (2009).
[58] X.-N. Wang, Z. Huang, and I. Sarcevic, Phys. Rev. Lett. 77, 231 (1996).
[59] H.L. Vargas, for the ALICE Collaboration, J. Phys. Conf. Ser. 389, 012004 (2012) arXiv:1208.0940.
[60] B. Abelev et al. [ALICE Collaboration], Phys. Rev. Lett. 110, 082302 (2013).