Medium Effects for Hadron-Tagged Jets in Proton–Proton Collisions

B. G. Zakharov*

Landau Institute for Theoretical Physics, Russian Academy of Sciences, Moscow, 117334 Russia

*e-mail: bgz@itp.ac.ru

Received August 17, 2022; revised August 17, 2022; accepted August 22, 2022

We study the medium modification factor $I_{pp}$ within the light-cone path integral approach to induced gluon emission. We use parametrization of the running coupling $\alpha_s(Q, T)$ which has a plateau around $Q \sim \kappa T$. We calculate $I_{pp}$ with no free parameters using $\kappa$ fitted to the LHC data on the nuclear modification factor $R_{AA}$.

We find that the theoretical multiplicity dependence of the ratio $I_{pp}/I_{pp}$ for 5.02 TeV $pp$ collisions agrees reasonably with the recent preliminary ALICE data [1].

DOI: 10.1134/S0021364022601713

INTRODUCTION

The observations of the transverse flow effects and the strong suppression of high- $p_T$ hadron spectra (jet quenching) in $AA$ collisions at RHIC and the LHC are the most compelling arguments in favor of the quark–gluon plasma (QGP) formation in $AA$ collisions (for reviews see, e.g., [2–4]). The results of hydrodynamic analyses of the flow effects support that the QGP is formed at the proper time $\tau_0 \sim 0.5–1$ fm [2, 3]. Such a picture of the QGP formation allows one to describe the RHIC and LHC heavy ion data on the nuclear modification factor $R_{AA}$ (see, e.g., [5–7]) in the pQCD jet quenching models with dominating contribution to the jet modification from induced gluon emission [8–13]. The observation of the ridge effect [14, 15] in $pp$ collisions at the LHC energies, suggests that the QGP can be created in $pp$ collisions as well. The mini QGP (mQGP) formation in $pp$ collisions is also supported by the steep growth of the midrapidity strange particle production at charged multiplicity $dN_{ch}/d\eta \sim 5$ [16]. This agrees with the onset of the QGP regime at $dN_{ch}/d\eta \sim 6$, found in [17] from behavior of $\langle p_T \rangle$ as a function of multiplicity, employing Van Hove’s arguments [18] that the phase transition should lead to an anomalous multiplicity dependence of $\langle p_T \rangle$. These estimates of the critical multiplicity density for the onset of the mQGP formation regime are smaller than the typical midrapidity charged multiplicity of the soft (underlying-event (UE)) hadrons for jet events in $pp$ collisions at the LHC energies—$dN_{ch}/d\eta \sim 10–15$ (it is bigger than the ordinary minimum bias multiplicity by a factor of $\sim 2–2.5$ [19]).

The typical size of the mQGP fireball created in $pp$ collisions should be $\sim 2–3$ fm. The Knudsen number, $Kn \sim \lambda/L$, for the mQGP may be estimated (as in [20]) with the help of the Drude formula and the lattice electrical conductivity of the QGP from lattice simulations [21]. In this way for the mQGP in the minimum bias $pp$ UEs at the proper time $\tau \sim 0.5(1)$ fm one can obtain for quarks quite small values $Kn \sim 0.2(0.1)$. For gluons the Knudsen number should be smaller by a factor of $\sim C_A/C_F = 9/4$. Such small values of $Kn$ support the collective hydrodynamic behavior of the mQGP fireball in $pp$ collisions. Note that it is possible (see, e.g., [22]) that the Knudsen number criterion $Kn \leq 1$ is too crude and underestimates the applicability of hydrodynamics. In [23] it was argued that the lower bound for the size of the mQGP with hydrodynamic behavior may be as small as $\sim 0.15$ fm. There it was demonstrated that neither local near-equilibrium nor near-isotropy are required for successful description of $pp$ data within viscous hydrodynamics. From the analysis of the data on $v_2$ it was concluded that in $pp$ collisions the hydrodynamics works at $dN_{ch}/d\eta \gtrsim 2$. This agrees well with the results of [24].

The mQGP formation in $pp$ collisions should lead to some jet modification. However, one can expect that the quenching effects in $pp$ collisions should be significantly smaller than in heavy ion collisions due to lower temperature of the mQGP and due to strong reduction of the induced gluon emission for a small size fireball. The latter is closely related to the anoma-
lously strong $L$-dependence of the radiative parton energy loss, $\Delta E_{\gamma}$, in a finite-size QCD matter [9] (as compared to predictions of the Bethe–Heitler formula). Fixed coupling calculations without the Coulomb effects within the BDMPS approach [9] give $\Delta E_{\gamma} \propto L^2$ for a static QGP, and $\Delta E_{\gamma} \propto L$ [25] for an expanding QGP with entropy density $s \propto 1/\tau$ (as in the Bjorken model [26] with purely longitudinal expansion of the QGP). The linear $L$-dependence of $\Delta E_{\gamma}$ for an expanding QGP remains approximately valid also for calculations with accurate treatment of the Coulomb effects with running $\alpha_s$ [27]. Calculations of the medium modification factor $R_{pp}$ (which is not directly observable quantity) within the light-cone path integral (LCPI) approach [10] with accurate treatment of the Coulomb effects and running $\alpha_s$ give a small deviation of $R_{pp}$ from unity at the LHC energies [7] ($R_{pp} \sim 0.8$ at $p_T \sim 10$ GeV). For this reason, observation of jet quenching in $pp$ collisions via a weak modification of the $p_T$-dependence of hadron spectra is practically impossible. A promising observable for quenching effects in $pp$ collisions is the variation with the UE activity of the medium modification factor $I_{pp}$ for the photon-tagged jet fragmentation functions (FFs) [28]. However, this measurement requires high statistics due to a very small cross section. This problem is absent for the modification factor $I_{pp}$ for the hadron-tagged jets. The medium modification factor $I_{pp}$ for the di-hadron production in $pp$ collisions can be written similarly to $AA$ collisions [29]

$$I_{pp}(p_T^a, \eta^a) = \frac{Y_{pp}^a(p_T^a, \eta^a, y^a)}{Y_{pp}^m(p_T^a, \eta^a, y^a)}$$  \hspace{1cm} (1)

where $p_T^a$ and $\eta^a$ are the transverse momenta and rapidities of the trigger ($h^\prime$) and the associated ($h^\prime \prime$) hadrons, $Y_{pp}^m$ is the per-trigger yield accounting the medium effects, and $Y_{pp}^m$ is the per-trigger yield calculated ignoring the medium effects. The per-trigger yields (similarly to $AA$ collisions [30, 29]) can be written in terms of the di-hadron (back-to-back) and one-hadron inclusive cross sections as

$$Y_{pp}^m(p_T^a, \eta^a, y^a, y^b) = \frac{d^4\sigma_{\text{m}}}{dp_T^a d\eta^a dy^a dy^b}/\frac{d^4\sigma_{\text{m}}}{dp_T^a d\eta^a dy^a dy^b} \left/ \frac{d^4\sigma_{\text{m}}}{dp_T^a d\eta^a dy^a dy^b} \right.$$ \hspace{1cm} (2)

Of course, the denominator in (1) is unobservable. But one can study the UE multiplicity dependence of $I_{pp}$, say, by using the ratio of the per-trigger yield to its minimum bias value. Because we can reasonably expect that the UE multiplicity dependence of the numerator and the denominator of (2) for $Y_{pp}^m$ is very similar, and consequently $Y_{pp}^m/Y_{pp}^m$ is approximately equal to 1. Then we have

$$\frac{I_{pp}(p_T^a, \eta^a, y^a, y^b)}{I_{pp}(p_T^a, \eta^a, y^a, y^b)} = \frac{Y_{pp}^m(p_T^a, \eta^a, y^a, y^b)}{Y_{pp}^m(p_T^a, \eta^a, y^a, y^b)} \left/ \frac{d^4\sigma_{\text{m}}}{dp_T^a d\eta^a dy^a dy^b} \right.$$ \hspace{1cm} (3)

This relation allows one to study the multiplicity dependence of $I_{pp}$ by measuring the per-trigger yield (which corresponds to the theoretical $Y_{pp}^m$). Recently, this method has been used by ALICE [1] in the first measurement of the variation of $I_{pp}$ with the UE multiplicity for the hadron-tagged jets for 5.02 TeV $pp$ collisions (for the trigger hadron momentum $8 < p_T^a < 15$ GeV, and the associated away side hadron momentum $4 < p_T^a < 6$ GeV). It was found that $I_{pp}$ decreases monotonically about 15% with increase in the UE activity in the range $5 \lesssim n_{ch}/d\eta \lesssim 20$ (we use $dN_{ch}/d\eta$ for the whole range of the azimuthal angle $\phi$ and the transverse momentum which is bigger by a factor of $\sim 4.4$ than the transverse side charged multiplicity $N_{ch}$ [1] for the kinematic region $\pi/3 \leq |\theta| \leq 2\pi/3, |\eta| < 0.8$, and $p_T > 0.5$ GeV). Such a decrease in $I_{pp}$ agrees qualitatively with the quenching effect obtained in [28] for the jet energy $E = 25$ GeV (which is of the order of the jet energy for the ALICE trigger particle momentum region [1]). For drawing a more definitive conclusion on whether the ALICE data [1] on $I_{pp}$ may be consistent with jet quenching in the mQGP, it is of course highly desirable to perform calculations of $I_{pp}$ for hadron-tagged jets accounting for the jet energy fluctuations and the quenching effects for both the back-to-back jets. In the present paper, we carry out such calculations of $I_{pp}$ for hadron-tagged jets for conditions of the ALICE experiment [1] within the LCPI approach [10] to the induced gluon emission. We calculate the medium modified FFs using the LCPI scheme with temperature dependent $\alpha_s$ [31] used in our recent global analysis [7] of the data on the nuclear modification factor $R_{AA}$.

In the present analysis we ignore the possible effect on $I_{pp}$ of the double parton scattering (DPS) with production the two nearly collinear di-jets. The production of such di-jet pairs potentially can lead to some dependence of $I_{pp}$ on the UE multiplicity due to correlations between the UE activity and intensity of the DPS [19]. However, one can easily understand that the multiplicity dependence of $I_{pp}$ observed in [1] cannot be due to the DPS because its effect should be of the opposite sign. Indeed, for the ALICE measurement we have $p_T^a/p_T^m \sim 2$, therefore, due to the steep decrease in the jet cross sections with the transverse momenta, the DPS contribution to $I_{pp}$ should be
dominated by the production of two di-jets with different energies, when the di-jet with the smallest energy contributes mostly to the production of the associated hadrons. Since it is reasonable to assume that the probability of the DPS should be bigger for events with high UE multiplicity, we conclude that the DPS mechanism should lead to a growth of \( I_{pp} \) with \( dN_{ch}/d\eta \); i.e., the effect should be of opposite sign to that observed in [1]. Our estimates show that the magnitude of the DPS contribution to \( I_{pp} \) should not exceed \( \sim 1-2\% \). This conclusion is supported by the recent simulations of [32] for conditions of the ALICE measurement [1] within the PYTHIA Monte Carlo model, which includes the DPS, where it was found that \( I_{pp} \) in the away region is consistent with unity, and independent of the UE multiplicity.

### THEORETICAL FRAMEWORK

#### FOR CALCULATION OF DI-HADRON AND ONE-HADRON CROSS SECTIONS

We calculate the di-hadron and one-hadron cross sections neglecting the internal parton transverse momenta in the colliding protons. In this case, in the approximation of independent hadron fragmentation for the back-to-back jets, the cross sections that we need can be written as

\[
\frac{d^2\sigma_{m,x}}{dp_{t1}^a dp_{t2}^a dy^a dy^j} = \int \frac{d^2\sigma_{ij}}{dp_{t1}^a dp_{t2}^a dy^a dy^j} D_{h_{jj}^a}^{m,x}(z^j, p_T) \times D_{h_{jj}^a}^{m,y}(z^j, p_T) \frac{d^2\sigma_{ij}}{dp_{t1}^a dp_{t2}^a dy^a dy^j},
\]

(4)

\[
\frac{d^2\sigma_{m,x}}{dp_{t1}^a dy^j} = \int \frac{d^2\sigma_{ij}}{dp_{t1}^a dy^j} D_{h_{ij}^a}^{m,x}(z^j, p_T) \frac{d^2\sigma_{ij}}{dp_{t1}^a dy^j},
\]

(5)

where \( \frac{d^2\sigma_{ij}}{dp_{t1}^a dp_{t2}^a dy^a dy^j} \) is the two-parton cross section of \( p + p \rightarrow i + j + X \) process for \( y_i = y_j', y_j = y^a \), \( p_T = p_T^i = p_T^j / z^j \) and \( z^a = z^a p_T^i / p_T^j \) \( \frac{d^2\sigma_{ij}}{dp_{t1}^a dy^j} \) is the one-parton cross section for \( p + p \rightarrow i + X \) process, and \( D_{h_{ij}^a}, D_{h_{jj}^a} \) are the \( i \rightarrow h^i \) and \( j \rightarrow h^j \) FFs. The averaging over the jet trajectories is implicit in Eqs. (4) and (5). As in our calculations of \( R_{pp} \) in [7], we perform calculations of \( I_{pp} \) for an effective azimuthally symmetric mQGP fireball for the whole range of the impact parameter. This approximation is reasonable because we calculate an azimuthally averaged quantity, and anyway the jet production is dominated by \( pp \) collisions with small impact parameters. As in [7], the distribution in the transverse plane of the di-jet production points is evaluated using the MIT bag model quark density (for simplicity we assume that gluons have the same distribution in the transverse plane as quarks).

For calculation of the medium-modified FFs \( D_{h_{ij}^a}(z, Q) \) we use the same method as in [7]. We write \( D_{h_{ij}^a} \) as the triple \( z \)-convolution

\[
D_{h_{ij}^a}(Q) = D_{h_{ij}^a}(Q_o) \otimes D_{h_{ik}^a}^{DGLAP}(Q),
\]

(6)

where \( D_{h_{ij}^a}^{DGLAP} \) is the DGLAP FF for \( i \rightarrow k \) transition, \( D_{h_{ik}^a} \) is the in-medium \( j \rightarrow k \) FF, and \( D_{h_{ik}^a} \) describe vacuum transition of the parton \( j \) to hadrons \( h^j \). As in [7], we calculate the vacuum FFs \( D_{h_{ij}^a}(z, Q) \) dropping the induced stage FFs \( D_{h_{ik}^a}^{DGLAP} \) in Eq. (6). To calculate the DGLAP FFs \( D_{h_{ik}^a}^{DGLAP} \) we use the PYTHIA event generator [33]. The in-medium FFs \( D_{h_{ik}^a} \) have been obtained in the approximation of the independent gluon emission [34] using the LCPI form of the induced gluon spectrum. In calculation of \( D_{h_{ij}^a} \), we account for the collisional energy loss, that is relatively small [35, 36], by treating it as a perturbation to the radiative mechanism (see [7] for details).

For the hadronization FFs \( D_{h_{ij}^a} \), we use the KKP [37] parametrization with \( Q_o = 2 \) GeV. Note that if the trigger hadron \( h^i \) is the leading particle in the jet (as for the ALICE data [1]), then in (6) the ordinary FF \( D_{h_{jj}^a} \) should be replaced by the probability distribution for production of the leading hadron, \( D_{h_{jj}^a} = S_{h_{jj}} D_{h_{jj}^a} \). For \( z > 0.5 \) we have \( S_{h_{jj}} = 1 \), but at \( z < 0.5 \) \( S_{h_{jj}} < 1 \) [38]. The behavior of \( S_{h_{jj}} \) at \( z < 0.5 \) depends on the hadronization model. Numerical simulation with PYTHIA event generator for the LUND fragmentation gives \( S_{h_{jj}} \) that is close to unity at \( z \approx 0.3(0.2) \) for quarks(gluons), and at smaller \( z \) it falls steeply to zero. We use a step-function parametrization \( S_{h_{jj}} = \theta(z-z_0) \) with the cutoff parameter \( z_0 \) defined from the normalization condition \( \int_0^1 dz D_{h_{jj}^a}(z) = 1 \) (such a normalization is reasonable since the probability of jet hadronization without charged particle production is \( \ll 1 \)), that leads to practically the same results for \( I_{pp} \) as \( S_{h_{jj}} \) obtained using the PYTHIA event generator. Note that, in principle, the presence of the suppression factor \( S_{h_{jj}} \) for the leading trigger hadron is not very important, since it has a quite small effect on \( I_{pp} \). This occurs due to the steep fall-off of the partonic cross
sections, resulting in strong suppression of the small-\(z\) contributions in Eqs. (4) and (5).

We calculate the gluon induced spectrum, which is needed for calculation of \(D_{jk}^{\mu}A\), using for running \(\alpha_s\) the same parametrization as in [7, 31] (supported by the lattice results for the in-medium \(\alpha_s\) in the coordinate space [39])

\[
\alpha_s(Q, T) = \begin{cases} 
\frac{4\pi}{9|\log\left(\frac{Q^2}{\Lambda_{QCD}^2}\right)|} & \text{if } Q > Q_{fr}(T), \\
\frac{Q \alpha_s'(T)}{c Q_{fr}(T)} & \text{if } Q < c Q_{fr}(T),
\end{cases}
\]

with \(c = 0.8\), \(Q_{fr}(T) = \Lambda_{QCD} \exp \left[\frac{2\pi}{9\alpha_s'(T)}\right]\) (we take \(\Lambda_{QCD} = 200\) MeV) and \(Q_{fr} = \kappa T\), where \(\kappa\) is a free parameter. We take \(\kappa = 2.55\), which has been obtained by \(\chi^2\) fitting of the LHC data on \(R_{AA}\) for 2.76 and 5.02 TeV Pb + Pb, and 5.44 TeV Xe + Xe collisions for scenario with the mQGP formation in \(pp\) collisions (in this scenario the theoretical \(R_{AA}\) reads \(R_{AA} = R_{AA}^{*}/R_{pp}\), where \(R_{AA}^{*}\) is the standard nuclear modification factor calculated with the pQCD pp cross section, and \(R_{pp}\) accounts for the medium jet modification in the mQGP produced in \(pp\) collisions, see [7] for details).

As in [7], we calculate \(D_{jk}^{\mu}A\) for a fireball with a uniform entropy/density distribution in the transverse plane. We checked that for \(R_{pp}\) this approximation gives practically the same results as calculations with the gaussian entropy distribution (with the same mean square radius and the total entropy). It is reasonable to expect that for \(I_{pp}\) the accuracy of this approximation should be quite good as well. This is connected with the fact that for \(pp\) collisions jet quenching is dominated by the emission of gluons with the formation length which is larger or of the order of the mQGP fireball size. In this regime the induced gluon emission is mostly sensitive to the total amount of the matter traversed by fast parton, and the density profile along its trajectory is of minor importance.

It is worth noting that possible an incomplete equilibration/isotropization of the matter in the initial stage [22, 40] is practically irrelevant for the induced gluon emission, because it is practically sensitive only to the parton number density of the matter (see discussion in [27]). But of course, the collectivity of the soft partons produced in \(pp\) collision is crucial, because it guarantees that they do not escape the interaction region in the free-stream regime without interaction with the jets.

**MODEL OF THE MQGP FIREBALL**

We assume the Bjorken 1 + 1D longitudinal expansion [26] of the QGP at the proper time \(\tau > \tau_0 = 0.5\) fm, that gives the entropy density \(s = s_0(\tau_0/\tau)\) and for \(\tau < \tau_0\) we take \(s = s_0(\tau/\tau_0)\). We calculate \(I_{pp}\) for two versions of the initial entropy density \(s_0\). In the first version (A) we assume that the UE charged particles are produced from decay of a thermalized mQGP fireball and write \(s_0\) in terms of the UE charged multiplicity density \(dN_{ch}/d\eta\) as

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

where \(R_f\) is the fireball radius, and \(C = dS/dy/dN_{ch}/d\eta = 7.67\) is the entropy/multiplicity ratio [41]. As in [7], we calculate \(R_f\) using the parametrization of the fireball radius \(R_f\) obtained in the CGC picture from [42]. In the interval \(dN_{ch}/d\eta \sim 5-20\), which is of interest here, it gives a monotonic increase in \(R_f\) from \(~1.2\) to \(~1.5\) fm. For the minimum bias UE multiplicity \(dN_{pp}/d\eta \sim 12.5\) (it can be obtained by interpolating the ATLAS data [43] for \(\sqrt{s} = 0.9\) and 7 TeV assuming that \(dN_{ch}/d\eta \sim s^3\)), we obtain \(R_f = 1.48\) fm. With this \(R_f\) using (8) we obtain the average initial fireball temperature \(T_0 \sim 226\) MeV for the ideal gas model entropy density, and \(T_0 \sim 256\) MeV for the lattice entropy density [44]. In the second version (B), to account for the fact that at a sufficiently low UE multiplicity density the assumption of a complete thermalization of the mQGP fireball may become invalid, we multiply the right-hand side of (8) by a multiplicity dependent coefficient \(K_0(dN_{ch}/d\eta)\). Considering the results of [16, 17] that support the onset of the QGP regime at \(dN_{ch}/d\eta \sim 5-6\), we use the parametrization \(K_0(N) = [1 + \tanh((N - N_0)/D)]/2\) with \(N_0 = 7.5\) and \(D = 2\). It mimics the evolution of the jet quenching effect from very small (at \(dN_{ch}/d\eta \sim 5\)) to almost its ordinary magnitude for a thermalized mQGP fireball (at \(dN_{ch}/d\eta \gtrsim 10\)). It is worth noting that the results of our previous analysis [7] of the available data on \(R_{AA}\) from RHIC and the LHC also support the scenario where in \(pp\) collisions the mQGP is formed at \(dN_{ch}/d\eta \gtrsim 10\), but is absent for \(dN_{ch}/d\eta \sim 5\), because it is this very scenario that allows to get the best simultaneous description of the data at the RHIC and the LHC energies (note that the minimum bias UE \(dN_{ch}/d\eta \sim 6\) and \(\gtrsim 10\) in these cases).
The approximation of the effective fireball, used in our jet quenching calculations, ignores fluctuations of the size and density of the mQGP, which fluctuate with the impact parameter of the collisions and for each impact parameter. However, the effect of such fluctuations should be small. Indeed, for a small size QGP the energy loss grows approximately linearly with the fireball size and the QGP density [27]. The latter is due to dominance of the rescattering contribution for a small size QGP (see discussion of the pattern of induced gluon emission in a small size QGP in [27]). These facts guarantee that the results for \( I_{pp} \) should be weakly sensitive to the size and density fluctuations of the fireball, and justify the use of the model of an effective fireball with fixed size and density.

**NUMERICAL RESULTS**

Figure 1 presents the multiplicity dependence of the ratio \( I_{pp}/\langle I_{pp} \rangle \) (here, as in (3), \( \langle I_{pp} \rangle \) is the minimum bias value of \( I_{pp} \)) obtained for the versions A (solid) and B (dotted) of the initial entropy \( s_0 \). To illustrate the sensitivity to the trigger hadron fragmentation model, we show the results obtained with the ordinary FF \( D_{1jj} \) (thin lines) and with the probability distribution of the leading trigger particle \( D'_{1jj} \) (thick lines). We evaluate \( \langle I_{pp} \rangle \) using the multiplicity dependence of the UE activity from CMS [45] obtained for 7 TeV \( pp \) collisions (rescaled to \( \sqrt{s} = 5.02 \) TeV assuming the KNO scaling). As can be seen from Fig. 1, our predictions are compatible within the errors with the ALICE data [1]. However, in both the versions the theoretical curves have somewhat less steeper decrease with increasing multiplicity than the data. The discrepancy becomes smaller in the version B with an incomplete thermalization at \( dN_{ch}/d\eta \lesssim 10 \). As one can see, the difference between the results for our two trigger particle fragmentation models (thin and thick lines) is relatively small.

**SUMMARY**

We have investigated the quenching effects for hadron-tagged jets in \( pp \) collisions. We have calculated the medium modification factor \( I_{pp} \) within the LCPI approach to induced gluon emission for parametrization of the running QCD coupling \( \alpha_s(Q, T) \) which has a plateau around \( Q \sim \kappa T \) (motivated by the lattice calculations of the effective QCD coupling in the QGP [39]). We use the value of \( \kappa \) fitted to the LHC data on the nuclear modification factor \( R_{AA} \) in 2.76 and 5.02 TeV Pb + Pb, and 5.44 TeV Xe + Xe collisions. We find that the theoretical predictions with no free parameters for the multiplicity dependence of the ratio \( I_{pp}/\langle I_{pp} \rangle \) for 5.02 TeV \( pp \) collisions are in reasonable agreement with the recent preliminary data from ALICE [1]. The description of the data becomes better for the scenario with an incomplete thermalization of the matter at \( dN_{ch}/d\eta \lesssim 10 \). Our results show that the drop of the ratio \( I_{pp}/\langle I_{pp} \rangle \) with the UE multiplicity, if confirmed by further measurements, may be viewed as the first direct evidence for the jet quenching in \( pp \) collisions.

**ACKNOWLEDGMENTS**

I am grateful to S. Tripathy for useful communication on some aspects of the ALICE measurement of \( I_{pp} \) [1].

**FUNDING**

This work was supported by the Russian Science Foundation, project no. 20-12-00200.

**CONFLICT OF INTEREST**

The author declares that he has no conflicts of interest.

**OPEN ACCESS**

This article is licensed under a Creative Commons Attribution 4.0 International License, which permits use, sharing, adaptation, distribution and reproduction in any medium or
REFERENCES

1. S. Tripathy (for ALICE Collab.), in Proceedings of the 24th DAE-BRNS High Energy Physics Symposium, December 14–18, 2020, Jaitni, India; arXiv: 2103.07218.

2. R. Derradi de Souza, T. Koide, and T. Kodama, Prog. Part. Nucl. Phys. 86, 35 (2016); arXiv: 1506.03863.

3. P. Romatschke and U. Romatschke, arXiv:1712.05815.

4. M. Connors, C. Nattrass, R. Reed, and S. Salur, Rev. Part. Nucl. Phys. 37, 444 (1994); nucl-th/9306003.

5. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 681, 291 (1997); arXiv: hep-ph/9607355.

6. A. Bzdak, B. Schenke, P. Tribedy, and R. Venugopalan, J. High Energy Phys., No. 09, 172 (2011).

7. A. Amato, G. Aarts, C. Allton, P. Giudice, S. Hands, and J.-I. Skullerud, PoS Lattice 2013, 176 (2014); arXiv: 1310.7466.

8. A. Adare et al. (PHENIX Collab.), Phys. Rev. C 78, 044905 (2008); arXiv:0801.4545.

9. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 135318 (2020); arXiv: 1902.03231.

10. M. Gyulassy and X. N. Wang, Nucl. Phys. B 420, 583 (1994); nucl-th/9306003.

11. R. Baier, Y. L. Dokshitzer, A. H. Mueller, and D. Schiff, Nucl. Phys. B 483, 291 (1997); arXiv: hep-ph/9607355.

12. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

13. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

14. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

15. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

16. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

17. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

18. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

19. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

20. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

21. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

22. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

23. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

24. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

25. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

26. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

27. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

28. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

29. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

30. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

31. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

32. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

33. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.

34. A. Bazavov, N. Brambilla, P. Petreczky, A. Vairo, and D. Schiff, Nucl. Phys. B 803, 291 (1997); arXiv: hep-ph/9607355.