Radiative Parton Energy Loss and Baryon Stopping in $AA$ Collisions

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We study the radiative energy loss contribution to proton stopping in $AA$ collisions. The analyses is performed within the light-cone path integral approach to induced gluon emission. We have found that the radiative correction can fill in part the midrapidity dip in the net proton rapidity distribution in $AA$ collisions at $\sqrt{s} \sim 10$ GeV. We argue that at $\sqrt{s} \sim 10$ GeV the net proton fluctuations at midrapidity may be dominated by the initial fluctuations of the proton flow, which, to a good accuracy, should be binomial.

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I. The baryon stopping in hadron and nucleus collisions has attracted much attention for a long time. However, there is no consensus yet on the mechanism of the baryon number transfer over a large rapidity interval (which is also closely related to the mechanism of $B\bar{B}$ annihilation). Presently, there is no answer to the most basic question about the baryon production: whether the baryon number carriers are quarks. It was proposed long ago [1] that a purely gluonic object—the so-called string junction (SJ) may play role of the baryon number carrier. In [1] it was suggested to describe the processes with baryons in the topological expansion (TE) scheme [2, 3] by treating the baryon as a Y shaped string configuration with the SJ. In this picture, the baryon number transfer is associated with the transfer of the SJ. Say, $B\bar{B}$ annihilation corresponds to SJ–$\overline{\text{SJ}}$ annihilation.

In the standard quark–gluon string model (QGSM) [4–6], based on the TE scheme [2], it is assumed that the baryon can be treated as a quark–diquark system with a point like diquark in the $\{3\}$ color state. Within the SJ picture [1], it means that the SJ belongs to diquark (we denote this mechanism by D-SJ). In this approximation, diquarks/antidiquarks are playing the role of the baryon number carriers in processes with baryons. In the QGSM the net proton ($\Delta p = p - \bar{p}$) midrapidity density in $pp$ collisions is related to the diquark distribution in the proton at small fractional momentum $x$, where it is $\propto x^{\alpha_{R}(0) - 2\alpha_{s}(0)}$ with $\alpha_{R}(0) \approx 0.5$ and $\alpha_{s}(0) \approx -0.5$ intercepts of the meson and nucleon Regge trajectories [7]. This leads to the energy dependence of the net proton midrapidity density $dN_{\Delta p}/dy \propto 1/x^{2.25}$. This disagrees strongly with the experimentally observed at ISR energies [8] $s$-dependence $dN_{\Delta p}/dy \propto 1/s^{n}$ with $n \approx 0.25$. The model with the diquarks in the antitriplet color state also underestimates the midrapidity net baryon production in $AA$ collisions [9, 10].

In [11], it was proposed that the baryon number transfer over a large rapidity interval can be related to the transition of the string configuration from $\{3\}$ to $\{6\}$ color state after one gluon $t$-channel exchange. The transition $D(\bar{\tau}) \rightarrow D(\bar{\epsilon})$ should lead to creation of the string configuration with the SJ located near the valence quark (we denote this mechanism as q-SJ). Hadronization of such string configurations leads naturally to the energy dependence of the midrapidity net proton density $dN/dy \propto 1/s^{\alpha_{q}(0)/2}$ that agrees with the data. Contribution of the q-SJ mechanism of the baryon number transfer over a large rapidity interval should be enhanced in $AA$ collisions due to an increase in the probability of the $D(\bar{\tau}) \rightarrow D(\bar{\epsilon})$ transition in the nuclear matter. Similar to $pp$ collisions, in $AA$ collisions the diquark breaking leads to the net proton midrapidity distribution $\propto 1/s^{\alpha_{q}(0)/2}$ [12].

The contribution from the q-SJ mechanism to the net proton midrapidity density in $AA$ collisions becomes of the order of that from the ordinary D-SJ mechanism at $\sqrt{s} \sim 20$ GeV. At energies $\sqrt{s} \sim 10$ GeV, the contribution of the q-SJ mechanism is relatively small. One more effect of the nuclear matter that can increase the baryon number flow to the midrapidity, which can potentially be important at $\sqrt{s} \lesssim 10$ GeV, is the diquark radiative energy loss. At high energies the radiative energy loss (which is not very large) cannot increase considerably the contribution of the D-SJ mechanism, and the radiative correction to the D-SJ
mechanism cannot compete with the q-SJ mechanism. However, at $\sqrt{s} \approx 10$ GeV, where the contribution of the q-SJ mechanism becomes small, the radiative mechanism might be important, and it is of great interest to understand how large the radiative contribution to the baryon stopping can be.

The question of how the baryon number flows to the midrapidity, which is interesting in itself, at energies $\sqrt{s} \approx 10$ GeV is also of great importance in connection with the future experiments at NICA and the beam energy scan (BES) program at RHIC aimed to search for the QCD critical point. One of the important signals of the critical point may be the non-monotonic variation of the moments of the net-baryon distribution in the central rapidity region [13]. However, in the event-by-event fluctuations of the net-baryon yield a considerable contribution may come from the initial state fluctuations [14, 15], for the D-SJ mechanism the variance for the net-baryon number should be bigger by a factor of ~3 than in the picture with the baryon number associated with quarks.

In this work, we calculate the diquark and quark radiative energy loss in $\AA$ collisions within the light-cone path integral (LCPI) approach [16–18] to the induced gluon emission. We study the effect of the diquark radiative energy loss in the nuclear matter on the net proton rapidity distribution for the D-SJ mechanism. We perform comparison with the NA49 Collaboration data [19] on the proton distribution in Pb + Pb collisions at $\sqrt{s} = 8.76$ GeV.

2. The scalar $ud$ diquark, which dominates in the nucleon wavefunction [20–22], is a rather compact object. It has a form factor $F(Q^2) = 1/[1 + Q^2/Q_0^2]$ with $Q_0^2 = 10$ GeV$^2$ [20]. The typical virtuality scale for the induced gluon emission in the nuclear matter ($-m_g^2 \lesssim 0.5$ GeV$^2$) is much smaller than $Q_0^2$. For this reason, the induced gluon radiation from scalar diquarks can be calculated as for a point like particle. This approximation should be reasonable even for the less compact vector $ud$ and $uu$ diquarks, for which $Q_0^2 \approx 2$ GeV$^2$ [20].

We will treat nuclei as uniform spheres. We consider $\AA$ collisions in the rest frame of one (target) of the colliding nuclei. To a good approximation, the induced gluon emission from each diquark/quark in the projectile nucleus can be evaluated similarly to $pA$ collisions, i.e., ignoring the presence of other nucleons. Indeed, the typical transverse size for the induced gluon emission is $-1/m_g^2$. The typical number of nucleons in the tube of radius $r \sim 1/m_g$ in the projectile nucleus is $\mu \sim A/R_A^2 m_g^2$. For $m_g \sim 400–800$ MeV it gives $\mu \sim 0.4–1.5$ for heavy nuclei with $A \sim 200$. This estimate shows that the possible collective effects in the projectile nucleus should not be strong.

In the LCPI approach the $x$-spectrum of the induced $a \to bc$ transition in the fractional longitudinal momentum $x = x_b = E_b/E_a$ can be written as (we take the $z$-axis along the projectile momentum) [16, 18]

$$
\frac{dP}{dx} = 2\text{Re} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} \exp[-i(z_b - z_i)/L_f] \\
\times \hat{g} \left[ \mathcal{H}(\rho_z, z_2, \rho_z, z_i) - \mathcal{H}(\rho_z, z_2, \rho_z, z_i) \right]_{\rho_i = 0}.
$$

(1)

Here, $\mathcal{H}$ is the Green’s function for the Schrödinger equation with the Hamiltonian

$$
H = -\frac{1}{2M} \left( \frac{\partial}{\partial \rho} \right)^2 + v(\rho, z),
$$

(2)

$$
\hat{v}(\rho, z) = -i \frac{\langle n(\rho)\sigma_3(\rho, x) \rangle}{2},
$$

(3)

where $M(x) = E_a x(1 - x)$, $L_f = 2M / [m_q^2 x_c + m_t^2 x_2 - m_q^2 x_c x_2]$, $n$ is the medium density, and $\sigma_3(\rho, x)$ is the cross section for the interaction of the $bc\bar{q}$ system with a nucleon, $\mathcal{H}_0$ is the vacuum Green’s function for $v = 0$. In Eq. (1), $\hat{g}$ is the vertex factor given by the expression

$$
\hat{g} = \alpha_s P_{gq} \frac{\partial}{\partial \rho} \frac{\partial}{\partial \rho},
$$

(4)

where $P_{gq}$ is the ordinary splitting function (for $q \to gq$, $P_{gq} = C_F[(1 + (1 - x)^2]/x$, and for $D \to gD$, $P_{gD} = 2C_F(1 - x)/x$). The three-body cross section, entering the imaginary potential (3), for the $D(q) \to gD(q)$ processes can be expressed in terms of the dipole cross section for the color singlet $q\bar{q}$ pair [23]

$$
\sigma_3(\rho, x) = \frac{9}{8} [\sigma_2(\rho) + \sigma_2((1 - x)p)] - \frac{1}{8} \sigma_2(xp).
$$

(5)

In the limit of $L_f \gg L$, the radiation rate is dominated by the configurations with large negative $z_i$ and large positive $z_2$ ($|z_2| \gg L$ (L is thickness of the target, for a nucleus with the radius $R_A (L) = 4R_A/3$). In this regime, the spectrum (1) may be calculated treating the transverse parton positions during traversing the target as frozen. The integrals over $z_{3,2}$ outside the target region on the right hand side of Eq. (1) can be expressed via the vacuum light-cone wavefunction for the $a \to bc$ [18], and the Green’ function in the medium is reduced to the ordinary Glauber attenuation...


de for the $bc\bar{a}$ system. This leads to the formula \[ \frac{dP_r}{dx} = 2 \int |\Psi(x,\mathbf{p})|^2 \left[ 1 - \exp \left( -\frac{nL\sigma_s(\mathbf{p},x)}{2} \right) \right], \] where $\Psi$ is the light-cone wavefunction of the $bc$ system in the $(x,\mathbf{p})$ representation. The frozen-size approximation is widely used in physics of $AA$ collisions; e.g., it was used for evaluation of the gluon spectrum in [25]. This approximation should be good for emission of gluons with energies $\omega \gg 10-20$ GeV (in the target nucleus rest frame). However, for $AA$ collision at $\sqrt{s} \sim 10$ GeV, interesting to us here, the typical gluon energies turn out not to high enough for its applicability.

We will perform calculation for the quadratic approximation $\sigma_2 = C_2 \rho^2$ ($C_2$ may be written in terms of the well-known transport coefficient $\hat{q}$ [26] as $C_2 = \hat{q} C_F / 2 n C_A$). In this case, the Hamiltonian (2) in the medium takes an oscillator form with the complex frequency given by $\Omega = \sqrt{in C_3 / M}$ with

$$ C_3 = C_2 \left[ \frac{9}{8} \left[ 1 + (1 - x)^2 \right] - \frac{x^2}{8} \right]. $$

In the oscillator approximation, the Green’s function can be written in the form

$$ \mathcal{G}(\mathbf{p}, z) \approx \frac{\gamma}{2\pi i} \exp \left( i(\alpha_2 \mathbf{p}^2 + \beta \mathbf{p} \cdot \mathbf{z} - \gamma \mathbf{p} \cdot \mathbf{z}) \right). $$

Only the parameter $\gamma$ is important in calculating the $x$-spectrum. From Eqs. (1) and (7), one can obtain

$$ \frac{dP}{dx} = \frac{\alpha_2 P_{\text{opt}}(x)}{\pi M^2} \Re \mathcal{G} \left( \int_{z_1} \int_{z_2} \exp \left( -i(x_2 - x_1) / L \right) (\gamma_0^2 - \gamma^2) \right). $$

Here, $\gamma_0 = M / (z_2 - z_1)$ is the parameter $\gamma$ in (7) for vacuum when $\Omega = 0$, and

$$ \gamma = M \Omega \left[ \begin{array}{c} \sin^{-1}(\Omega(z_2 - z_1)) \\ \cos(\Omega(L - z_2))(\tan(\Omega(L - z_1)) + \Omega(z_2 - L)) \\ \cos(\Omega z_2)\left[ \tan(\Omega z_2) + \Omega z_2 \right] \\ \cos(\Omega L)\left[ \tan(\Omega L - z_2 - L) + \tan(\Omega L)(1 - \Omega^2(z_2 - L)z_2) \right] \cos(\Omega L) \right]. $$

3. For numerical calculations we take $\alpha_2 = 0.5$. We use the value $\hat{q} = 0.01$ GeV$^3$, supported by calculations of the coefficient $C_2$ using the double gluon formula for the dipole cross section [17]. For the quark and diquark masses we take $m_q = m_\rho = 300$ MeV. However, the radiative energy loss is only weakly dependent on the specific choice of the mass of the initial fast particle. However, its dependence on the gluon mass, which plays the role of the infrared cutoff, turns out to be rather strong.\footnote{Note that the situation with the infrared sensitivity of the radiative energy loss for a parton incident on the medium from outside is very different from that for a parton produced inside the medium. In the former case, the fast parton approaches the target with a formed gluon cloud with the transverse size $\sim m_\rho$. While in the latter case, the fast parton is produced without the formed gluon cloud. After gluon emission, the transverse size of the two-parton system grows $\sim \sqrt{L/\omega}$ [27]. As a result, the spectrum for gluons with $L_f \gtrsim L$ depends weakly on the gluon mass. This leads to a weak infrared sensitivity of the total parton energy loss [27].}

We perform numerical computations for $m_g = 750$ MeV and $m_g = 400$ MeV. The former value was obtained in the analysis within the dipole BFKL equation [28] of the data on the low-$x$ proton structure function $F_2$. The values of $m_g$ in the range 400–800 MeV have been predicted in the non-perturbative calculations using the Dyson–Schwinger equations [29–31]. We take for Pb nucleus $r_A = 6.49$ fm, which corresponds to the typical parton path length in the nuclear matter $\langle L \rangle \approx 8.65$ fm.

In Fig. 1, we present the result for the total radiative energy loss for diquarks and quarks for Pb nucleus

$$ \Delta E = E \int_{x_{\text{min}}}^{x_{\text{max}}} dx \frac{dP}{dx}. $$

We take $x_{\text{min}} = m_g / E$ and $x_{\text{max}} = 1 - m_g / E$. To illustrate the effect of the parton transverse motion we show in Fig. 1 also the prediction obtained with the spectrum calculated in the frozen-size approximation (6). One can see that the frozen-size approximation underestimates the energy loss by a factor of 1.6–1.7(1.2–1.3) at $m_g = 400$ (750) MeV. Thus, the frozen-size approximation turns out to be rather crude. Figure 1 shows that the ratio $\Delta E / E$ flattens at $E \approx 50–100$ GeV, and at $E \approx 1000$ GeV for diquark $\Delta E / E \approx 0.18(0.062)$ at $m_g = 400$ (750) MeV. Note that some violation of the $1/m_g^2$ scaling for $\Delta E$ is due to the Landau–Pomeranchuk–Migdal suppression that is stronger for smaller $m_g$.

For evaluation of the net baryon spectrum within the QGSM one needs only the contribution of the diquark fragmentations to baryons, i.e., the contribution from color strings with the valence diquarks. In $pp$
and $AA$ collisions the contributions from the projectile nucleons (diquark–quark strings) and from the target nucleons (quark–diquarks strings) can be evaluated independently. The valence quarks and sea quarks/diquarks are irrelevant. The net baryon $x$-distribution for each projectile nucleon in the QGSM can be written as \[9, 10\]

\[
\frac{dN_{AA}^{BB}}{dx} = \int \frac{dz}{z} \rho_D(z) D^B_D(x/z),
\]

where $\rho_D$ is the diquark distribution in the projectile nucleon, $D^B_D$ is the $D \rightarrow B$ fragmentation function. We use for the diquark distribution the parametrization of the form used in the \[9, 10\]

\[
\rho_D(x) = C_x \alpha_s(0)^{-2 \alpha_s(0)} (1-x)^{-\alpha_s(0)-1+k},
\]

where $C = \left[ \int_0^1 dx \rho_D(x) \right]^{-1}$ is the normalization constant. The parameter $k$ controls the energy degradation due to creation of additional sea $gQ^2$ pairs in the nucleon, that lead to formation of additional color strings in the final state. Within the QGSM there is no rigorous theoretical method for its determination. It is input by hand in order to reproduce experimental data. In the case of $pp$ collisions, data on the charged particle multiplicities can be reasonably described using the quasieikonal formulas \[32\]. However, the model fails to describe the $AA$ data \[32\]. In \[9, 10, 12\] the value of $k$ was set to $v = 2N_{col}/N_{part}$ (here, $N_{col}$ and $N_{part}$ are the number of the binary collisions and the number of participant nucleons evaluated in the Glaber model). In terms of the wounded nucleon model \[33, 34\] $v$ is simply the average number of inelastic collisions per participant nucleon. This prescription gives reasonable agreement with the data on charged particle multiplicity in $AA$ collisions \[9\]. In the present analysis we use a slightly modified formula $k = v + \langle k_N \rangle - 1$, where $\langle k_N \rangle$ is the average number of cut Pomerons in the quasieikonal formulas for $pp$ collisions ($\langle k_N \rangle \approx 1.65$ at $E = 40$ GeV). This modification practically does not change the charged particle multiplicity for $AA$ collisions, but it somewhat improves description of the data on the baryon stopping in $pp$ and $AA$ collisions. In addition, this prescription ensures matching of the predictions for very peripheral $AA$ collisions with that for $pp$ collisions.

The induced gluon emission in $AA$ collisions softens the diquark and quark distributions due to the radiative energy shift. However, if one neglects the diquark transitions to the $\{6\}$ color state, the radiative change of the quark distribution is irrelevant for the baryon stopping. However, the radiative diquark energy loss should increase the baryon stopping as compared to the predictions of the QGSM. In the presence of the induced gluon emission, we write the diquark distribution as (we denote it by $\rho_D^{\text{eff}}$)

\[
\rho_D^{\text{eff}}(x) = \rho_D(x) + \Delta \rho_D(x),
\]

where $\Delta \rho_D(x)$ is the radiative correction. In terms of the induced gluon spectrum (1), the radiative correction to the QGSM diquark distribution $\rho_D$ can be written as

\[
\Delta \rho_D(x) = \int_{z_{\min}}^{1-x} dz \frac{Q P}{z} \left[ \frac{\rho_D(x/(1-z))}{1-z} - \rho_D(x) \right].
\]

Here, the second term in the square brackets is due to reduction of the probability to find the diquark without gluon emission. Note that due to this term the radiative softening of the diquark distribution is insensitive to the $z$-dependence of the gluon spectrum at $z \rightarrow 0$.

In Fig. 2a, we confront the results of our calculations with and without the induced gluon emission of the net proton rapidity spectrum (in the center of mass frame) with the data from NA49 Collaboration for Pb + Pb collisions at $E = 40$ GeV ($\sqrt{s} = 8.76$ GeV) \[19\]. The theoretical curves are obtained by summing...
the contributions of the valence proton flow from the projectile and the target nuclei. The calculations are performed with the fragmentation function \( D \rightarrow p \) from [35] \( D_0^i(z) = a \gamma^{1.5} \). The normalization parameter \( a \) depends on the strange suppression factor \( S/L \) \((2L + S = 1)\) introduced in [35], for which we take \( S/L = 0.367 \). This value was adjusted to reproduce the experimental midrapidity ratio \( p/\Lambda = 0.4 \) for Pb + Pb collisions at \( E = 40 \) GeV [19, 36]. From Fig. 2a, one sees that the radiative correction partly fills the minimum at \( y = 0 \). It increases the spectrum at \( y = 0 \) by a factor of about 1.35 and 1.12 for \( m_g = 400 \) and 750 MeV, respectively. Thus, we conclude that for \( m_g = 750 \) MeV the radiative effect is relatively weak. For \( m_g = 400 \) MeV, the radiative effect is quite strong, and the theoretical spectrum overshoots the data at \( y \sim 0 \). Of course, one can improve agreement with the data by changing the QGSM parton distributions, which have not stringent theoretical constraints in the QGSM. However, our purpose in this preliminary study is to understand if the radiative energy loss may be important and not as good as possible fitting to the data.

Our results shown in Fig. 2a do not include the possible baryon diffusion in the hot QCD matter produced after interaction of the colliding nuclei at the proper time \( \tau_0 = 2R_\gamma/\gamma \) (\( \gamma \) is the nucleus Lorentz factor in the center of mass frame). For \( E = 40 \) GeV \((\gamma = 4.7)\) we have \( \tau_0 \approx 2.8 \) fm. The change of the baryon rapidity is related to its random walk in the longitudinal \( z \) direction in the comoving frame of the QCD matter. For the diffusion mean squared rapidity fluctuations one can easily obtain

\[
\langle \Delta y_d^2(\tau) \rangle^{1/2} \approx \int_{\tau_0}^\infty \frac{d\langle \Delta z(\tau) \rangle^{1/2}}{d\tau} \frac{d\tau}{\tau}.
\]  

(15)

Making use of the random walk formula \( \langle \Delta z(\tau) \rangle = (\tau - \tau_0)\gamma \bar{v}l_c/3 \) here, \( \bar{v} \) is the mean proton velocity and \( l_c \) is the proton mean free path) from (15) one obtains

\[
\langle \Delta y_d^2(\tau) \rangle^{1/2} \approx \left[ \frac{\bar{v}l_c \pi}{3\tau_0} \right]^{1/2} \arcsin \left[ \sqrt{\frac{\tau_0}{\tau_f}} \right].
\]  

(16)

where \( \tau_f \) is the freeze out proper time. From the relation for the initial entropy density \( s_0 = C \frac{dN_{ch}}{\tau_0 \pi R_i^4 \frac{d\eta}{d\eta}} \) \((C = ds/dy/dN_{ch}/d\eta = 7.67 \) is the entropy/multiplicity ratio [37]), using the midrapidity charged particle density \( dN_{ch}/d\eta = 280 \) [19, 38] and the matter equation of state of lattice QCD [39], we obtain \( T_0 = 170 - 175 \) MeV for \( \tau_0 = 2.8 \) fm. Using the hadron gas model for the relevant range of the temperature \( T \sim 140 - 175 \) MeV from (16) with \( \tau_f = 13 \) fm [40] one obtains \( \langle \Delta y_d^2(\tau) \rangle^{1/2} \approx 0.3 \). To illustrate the magnitude of the diffusion correction, in Fig. 2b we show the prediction obtained with the Gaussian smearing with the width \( \sigma_d = 0.3 \). One sees that the smearing partly fills the minimum.

It is seen from Fig. 2 that the diffusion effect should not affect strongly the net proton rapidity distribution. However, one can expect that for \( \sigma_d \sim 0.3 \) the diffusion may be important for the event-by-event net proton fluctuations for the rapidity window \( \Delta y \) about \((2-3)\sigma_d \). In this regime, the diffusion can modify somewhat the primordial net proton yield fluctuations. However, at the same time, it is clear that for \( \Delta y/\sigma_d \sim 3 \) this modification cannot be strong. For this reason, the net baryon charge fluctuations cannot be described by the grand canonical ensemble formulas, which require \( \sigma_d \gg \Delta y \), and one cannot expect to observe a real critical regime. In the light of this, the absence of a clear signal of the critical point in the net
proton fluctuations at $\sqrt{s} \sim 10$ GeV for $|y| < 0.5$ in Au + Au collisions from STAR [41] is not surprising. For the diquark mechanism of the baryon flow, to a good accuracy, these fluctuations should be binomial. This agrees with the STAR observation [41] that the net proton fluctuations are close to binomial/Poissonian at $\sqrt{s} \sim 10$ GeV, where the antiproton yield becomes very small. In the scenario of [42], in which the carriers of baryon number are quarks, the binomial/Poissonian distribution occurs for the net quark fluctuations. In this case $〈(N_{\text{pp}} - 〈N_{\text{pp}}〉)^2〉 \approx 〈N_{\text{pp}}〉/3$ (cf. [15]), which contradicts the STAR measurement [41]. Note that even without calculating the diffusion width $\sigma_{\text{d}}$, the existence of the dip in the experimental net proton distribution at $y = 0$ from NA49 [19] says that $\sigma_{\text{d}}$ is considerably smaller than ~1. Because, numerical calculations show that for $\sigma_{\text{d}} \sim 1$ the diffusion should completely wash out the dip. So, one can say, that we have an experimental evidence that at $\sqrt{s} \sim 10$ GeV the net proton fluctuations for $\Delta y \sim 1$ cannot be close the critical point regime.

One remark is in order regarding the status of the parton radiative energy loss in the nuclear matter in the string model. One might wonder whether the adding of the radiative energy loss leads to the double counting, because the QGSM already includes the energy degradation of the ends of the strings for multiple Pomeron exchanges. However, in the TE/QGSM the energy degradation is due to creation of the parallel hadronic states, say, $h_{\pi} \rightarrow h_{\pi}h_{\pi}$ for two Pomeron exchanges. In the TE [2], based on the $1/N$ expansion ($N = N_{\pi}, N_{\pi}/N_{\pi} = \text{const}$), such parallel energy degradation is due to the quark loops that survive at $N \rightarrow \infty$, only because $N_{\pi} \rightarrow \infty$ together with $N_{\pi}$. On the contrary, the induced gluon emission is a purely gluonic effect, which does not depend on $N_{\pi}$ at all. Contrary to the parallel energy degradation in the TE, this mechanism does not increase the number of the color strings (as in the LUND model [43] the radiated gluon may be treated as a kink on the available string attached to the diquark). Thus, it is clear the radiative mechanism is absent in the ordinary QGSM [4–6]. Within the TE scheme, it is a subleading (in $1/N$) effect. Our calculations show that its contribution to the midrapidity net proton density at $\sqrt{s} \sim 10$ GeV may be quite large.

4. In summary, we have studied, for the first time, the radiative energy loss contribution to the proton stopping in AA collisions in the energy region $\sqrt{s} \sim 10$ GeV, which is of interest in connection with the future experiments at NICA and the BES program at RHIC. The analysis is based on the LCPI approach [16, 17] to the induced gluon emission. Calculations are performed beyond the soft gluon approximation. We have compared our results with the data on the proton rapidity distribution from NA49 [19] for central Pb + Pb collisions at $E = 40$ GeV. We have found that the radiative correction can fill in partly the midrapidity dip in the net proton rapidity distribution. For $m_{\pi} \sim 400$ MeV, the radiative effect turns out to be rather strong, and its contribution to the midrapidity net proton density is comparable to the prediction of the ordinary QGSM.

We argue that the net proton midrapidity fluctuations at $\sqrt{s} \sim 10$ GeV may be dominated by the initial fluctuations of the proton flow to the midrapidity region. The fact the net proton distribution from STAR [41] at $\sqrt{s} \sim 10$ GeV is close to the Poissonian/binomial distribution provides evidence for the diquark mechanism of the baryon flow, while disfavors the picture where quarks (without SJ) transmit the baryon number [42]. The dominance of the initial state fluctuations makes questionable observation of the QCD critical point via the data on the net proton fluctuations in a rapidity window $\Delta y \sim 1$.

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REFERENCES

1. G. C. Rossi and G. Veneziano, Nucl. Phys. B 123, 507 (1977).
2. G. Veneziano, Phys. Lett. B 52, 220 (1974).
3. G. Veneziano, Nucl. Phys. B 117, 519 (1976).
4. G. Cohen-Tannoudji, A. E. Hassouni, J. Kalinowski, and R. B. Peschanski, Phys. Rev. D 19, 3397 (1979).
5. A. Capella and J. Tran Thanh Van, Phys. Lett. B 114, 450 (1982).
6. A. B. Kaidalov, Phys. Lett. B 116, 459 (1982).
7. A. B. Kaidalov and O. I. Piskunova, Z. Phys. C 30, 145 (1986).
8. B. Alper, H. Boggild, P. Booth, et al., Nucl. Phys. B 100, 237 (1975).
9. A. Capella, A. Kaidalov, A. K. Akil, C. Merino, and J. Tran Thanh Van, Z. Phys. C 70, 507 (1996); hep-ph/9507250.
10. A. Capella and C. A. Salgado, Phys. Rev. C 60, 054906 (1999); hep-ph/9903414.
11. B. Z. Kopeliovich and B. G. Zakharov, Z. Phys. C 43, 241 (1989).
12. A. Capella and B. Z. Kopeliovich, Phys. Lett. B 381, 325 (1996); hep-ph/9603279.
13. M. A. Stephanov, Phys. Rev. Lett. 102, 032301 (2009); arXiv:0809.3450.
14. M. Asakawa and M. Kitazawa, Prog. Part. Nucl. Phys. 90, 299 (2016); arXiv:1512.05038.
15. M. Asakawa, U. W. Heinz, and B. Muller, Phys. Rev. Lett. 85, 2072 (2000); hep-ph/0003169.
16. B. G. Zakharov, JETP Lett. 63, 952 (1996); hep-ph/9607440.
17. B. G. Zakharov, JETP Lett. 65, 615 (1997); hep-ph/9704255.
18. B. G. Zakharov, Phys. Atom. Nucl. 61, 838 (1998); hep-ph/9807540.
19. T. Anticic, B. Baatar, D. Barna, et al. (NA49 Collab.), Phys. Rev. C 83, 014901 (2011); arXiv:1009.1747.
20. M. Anselmino, E. Predazzi, S. Ekelin, S. Fredriksson, and D. B. Lichtenberg, Rev. Mod. Phys. 65, 1199 (1993).
21. V. T. Kim, Mod. Phys. Lett. A 3, 909 (1988).
22. C. Chen, B. El-Bennich, C. D. Roberts, S. M. Schmidt, J. Segovia, and S. Wan, Phys. Rev. D 97, 034016 (2018); arXiv:1711.03142.
23. N. N. Nikolaev, B. G. Zakharov, and V. R. Zoller, JETP Lett. 59, 6 (1994); hep-ph/9412268.
24. N. N. Nikolaev, G. Piller, and B. G. Zakharov, JETP 81, 851 (1995); hep-ph/9412344.
25. Y. V. Kovchegov and A. H. Mueller, Nucl. Phys. B 529, 451 (1998); hep-ph/9802440.
26. R. Baier, Y. L. Dokshitzer, A. H. Mueller, S. Peigne, and D. Schiff, Nucl. Phys. B 483, 291 (1997); hep-ph/9607355.
27. B. G. Zakharov, JETP Lett. 73, 49 (2001); hep-ph/0012360.
28. N. N. Nikolaev and B. G. Zakharov, Phys. Lett. B 327, 149 (1994); hep-ph/9402209.
29. Si-xue Qin, L. Chang, Yu-xin Liu, C. D. Roberts, and D. J. Wilson, Phys. Rev. C 84, 042202 (2011); arXiv:1108.0603.
30. J. Papavassiliou, arXiv:1112.0174.
31. A. C. Aguilar, D. Binosi, J. Papavassiliou, and J. Rodriguez-Quintero, Phys. Rev. D 80, 085018 (2009); arXiv:0906.2633.
32. A. Capella and E. G. Ferreiro, Eur. Phys. J. C 72, 1936 (2012); arXiv:1110.6839.
33. A. Bialas, M. Bleszynski, and W. Czyz, Nucl. Phys. B 111, 461 (1976).