

In-medium QCD and Cherenkov gluons

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Abstract. The equations of in-medium gluodynamics are proposed. Their classical lowest order solution is explicitly shown for a color charge moving with constant speed. For nuclear permittivity larger than 1 it describes emission of Cherenkov gluons resembling the results of classical electrodynamics. The choice of nuclear permittivity and Lorentz invariance of the problem are discussed. Effects induced by the transversely and longitudinally moving (relative to the collision axis) partons at LHC energies are described.

1 Introduction

The collective effects observed in ultrarelativistic heavy-ion collisions at SPS and RHIC [1–5] have supported the conjecture of quark–gluon plasma (QGP) having been formed in these processes. The properties and evolution of this medium are widely debated. At the simplest level it is assumed to consist of a set of current quarks and gluons. It so happens, however, that their interaction is quite strong, so that the notion of a strongly interacting quark–gluon plasma (sQGP) has been introduced. Moreover, this substance reminds one of an ideal liquid rather than a gas. Whether perturbative quantum chromodynamics (pQCD) is applicable to the description of the excitation modes of this matter is doubtful. Correspondingly, the popular theoretical approaches use either classical solutions of the in-vacuum QCD equations at the initial stage or hydrodynamics at the final stage of its evolution.

The collective excitation modes of the medium may, however, play a crucial role. One of the ways to gain more knowledge about the excitation modes is to consider the propagation of relativistic partons through this matter. Phenomenologically their impact would be described by the nuclear permittivity of the matter corresponding to its response to passing partons. Namely this approach is most successful for electro-dynamical processes in matter. Therefore, it is reasonable to modify the QCD equations by taking into account collective properties of the quark–gluon medium. For the sake of simplicity we consider here gluodynamics only. The generalization to quarks is straightforward.

The classical lowest order solution of these equations coincides with abelian electro-dynamical results up to a trivial color factor. One of the most spectacular of them is Cherenkov radiation and its properties. Now, Cherenkov gluons take the place of Cherenkov photons [6–8]. Their

emission in high-energy hadronic collisions is described by the same formulae but with the nuclear permittivity in place of the usual one. It should be properly defined. Actually, one considers them as quasiparticles, i.e. quanta of the medium excitations with properties determined by the permittivity. The interplay of the properties of the medium and the velocity of the particle is crucial for the radiation field.

Another important problem of this approach is related to the notion of the rest system of the medium. Lorentz invariance is lost if the permittivity is introduced.¹ Therefore, one has to choose the proper coordinate system, where its definition is working. While it is simple for macroscopic media in electrodynamics, one should consider partons moving in different directions with different energies in the case of heavy-ion collisions. It has a direct impact on the properties of the emitted particles. The fast evolution of the medium and its short lifetime differ from common electro-dynamical examples.

All these problems are discussed in what follows.

2 Equations of in-medium gluodynamics

To begin, let us recall the classical in-vacuum Yang–Mills equations,

$$D_\mu F^{\mu\nu} = J^\nu, \quad (1)$$

$$F^{\mu\nu} = \partial^\mu A^\nu - \partial^\nu A^\mu - ig[A^\mu, A^\nu], \quad (2)$$

where $A^\mu = iA_a^\mu T_a$; $A_a(A_a^0 \equiv \Phi_a, \mathbf{A}_a)$ are the gauge field (scalar and vector) potentials, the color matrices T_a satisfy the relation $[T_a, T_b] = if_{abc}T_c$, $D_\mu = \partial_\mu - ig[A_\mu, \cdot]$, $J^\nu(\rho, \mathbf{j})$ is a classical source current, $\hbar = c = 1$ and the metric tensor is $g^{\mu\nu} = \text{diag}(+, -, -, -)$.

¹ In principle, this deficiency is cured by the relativistic generalization of the notion of permittivity (e.g., see [9]).

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In the covariant gauge $\partial_\mu A^\mu = 0$ they are written as

$$\square A^\mu = J^\mu + ig[A_\nu, \partial^\nu A^\mu + F^{\mu\nu}], \quad (3)$$

where \square is the d'Alembertian operator. It was shown [10] (and is confirmed in what follows) that in this gauge the classical gluon field is given by the solution of the corresponding abelian problem.

The chromoelectric and chromomagnetic fields are

$$E^\mu = F^{\mu 0}, \quad (4)$$

$$B^\mu = -\frac{1}{2}\epsilon^{\mu ij}F^{ij}, \quad (5)$$

or, as functions of the gauge potentials in vector notation,

$$\mathbf{E}_a = -\text{grad } \Phi_a - \frac{\partial \mathbf{A}_a}{\partial t} + gf_{abc}\mathbf{A}_b\Phi_c, \quad (6)$$

$$\mathbf{B}_a = \text{curl } \mathbf{A}_a - \frac{1}{2}gf_{abc}[\mathbf{A}_b\mathbf{A}_c]. \quad (7)$$

The equations of motion (1) in vector form are written

$$\text{div } \mathbf{E}_a - gf_{abc}\mathbf{A}_b\mathbf{E}_c = \rho_a, \quad (8)$$

$$\text{curl } \mathbf{B}_a - \frac{\partial \mathbf{E}_a}{\partial t} - gf_{abc}(\Phi_b\mathbf{E}_c + [\mathbf{A}_b\mathbf{B}_c]) = \mathbf{j}_a. \quad (9)$$

The abelian equations of in-vacuum electrodynamics are obtained from (3) if the second term on its right-hand side is put equal to zero and color indices are omitted. The medium is accounted for if \mathbf{E} is replaced by $\mathbf{D} = \epsilon\mathbf{E}$ in $F^{\mu\nu}$, i.e. in (4).² Therefore (8) and (9) in vector form are most suitable for generalization to the in-medium case. The equations of in-medium electrodynamics differ from the in-vacuum ones by the dielectric permittivity $\epsilon \neq 1$ entering; we now have

$$\Delta \mathbf{A} - \epsilon \frac{\partial^2 \mathbf{A}}{\partial t^2} = -\mathbf{j}, \quad (10)$$

$$\epsilon \left(\Delta \Phi - \epsilon \frac{\partial^2 \Phi}{\partial t^2} \right) = -\rho. \quad (11)$$

The permittivity describes the matter response to the induced fields, which is assumed to be linear and constant in (10) and (11). It is determined by the distribution of internal current sources in the medium. Then external currents are only left on the right-hand sides of these equations.

Now, the Lorentz gauge condition is

$$\text{div } \mathbf{A} + \epsilon \frac{\partial \Phi}{\partial t} = 0. \quad (12)$$

The Lorentz invariance is broken if $\epsilon \neq 1$ in front of the second terms on the left-hand sides. Then one has to work in the coordinate system where the substance is at rest. The values of ϵ are determined just there. To cancel these requirements one must use Minkowski relations between

\mathbf{D} , \mathbf{E} , \mathbf{B} and \mathbf{H} valid for a moving medium [9]. It leads to more complicated formulae, and we do not use them in this paper.

The most important property of the solutions of these equations is that while the in-vacuum ($\epsilon = 1$) equations do not admit any radiation processes, it happens for $\epsilon \neq 1$ that there are solutions of these equations with non-zero Poynting vector.

Now we are ready to write down the equations of the in-medium gluodynamics generalizing (3) in the same way as (10) and (11) are derived in electrodynamics. We introduce the nuclear permittivity and denote it also by ϵ , since this will not lead to any confusion. After that, one should replace \mathbf{E}_a in (8) and (9) by $\epsilon\mathbf{E}_a$ and get

$$\epsilon(\text{div } \mathbf{E}_a - gf_{abc}\mathbf{A}_b\mathbf{E}_c) = \rho_a, \quad (13)$$

$$\text{curl } \mathbf{B}_a - \epsilon \frac{\partial \mathbf{E}_a}{\partial t} - gf_{abc}(\epsilon\Phi_b\mathbf{E}_c + [\mathbf{A}_b\mathbf{B}_c]) = \mathbf{j}_a. \quad (14)$$

The space-time dispersion of ϵ is neglected here.

In terms of potentials these equations are cast in the form

$$\begin{aligned} \Delta \mathbf{A}_a - \epsilon \frac{\partial^2 \mathbf{A}_a}{\partial t^2} = & \\ -\mathbf{j}_a - gf_{abc} \left(\frac{1}{2} \text{curl}[\mathbf{A}_b, \mathbf{A}_c] + \frac{\partial}{\partial t}(\mathbf{A}_b\Phi_c) + [\mathbf{A}_b \text{curl } \mathbf{A}_c] \right. & \\ \left. - \epsilon\Phi_b \frac{\partial \mathbf{A}_c}{\partial t} - \epsilon\Phi_b \text{grad } \Phi_c - \frac{1}{2}gf_{cmn}[\mathbf{A}_b[\mathbf{A}_m\mathbf{A}_n]] \right. & \\ \left. + g\epsilon f_{cmn}\Phi_b\mathbf{A}_m\Phi_n \right), & \quad (15) \end{aligned}$$

$$\begin{aligned} \Delta \Phi_a - \epsilon \frac{\partial^2 \Phi_a}{\partial t^2} & \\ = -\frac{\rho_a}{\epsilon} + gf_{abc} \left(2\mathbf{A}_c \text{grad } \Phi_b + \mathbf{A}_b \frac{\partial \mathbf{A}_c}{\partial t} + \frac{\partial \Phi_b}{\partial t} \Phi_c \right) & \\ - g^2 f_{amn}f_{nlb}\mathbf{A}_m\mathbf{A}_l\Phi_b. & \quad (16) \end{aligned}$$

If the terms with explicitly shown coupling constant g are omitted, one gets the set of abelian equations, which differ from electrodynamical equations (10) and (11) by the color index a only. Their solutions are shown in the next section. The external current is ascribed to a parton moving fast relative to the other partons ‘‘at rest’’.

The potentials are linear in g because the classical current J^μ is linear also. Therefore, the omitted terms are of the order of g^3 and can be taken into account as a perturbation. It was done in [11–14] for in-vacuum gluodynamics. Here, the general procedure is the same. After getting an explicit lowest order solution (see the next section) one exploits it together with the non-abelian current conservation condition to find the current component proportional to g^3 . Then with the help of (15) and (16) one finds the potentials up to the order g^3 . They can be represented as integrals convoluting the current with the corresponding in-medium Green function. The higher order corrections may be obtained in the same way. We postpone their consideration for further publications.

The crucial distinction between (3) and (15), (16) is that there is no radiation (the field strength is zero in the forward light-cone and no gluons are produced) in the

² ϵ denotes the dielectric permittivity of the medium. The magnetic permittivity is put equal to 1 to simplify the formulae.

lowest order solution of (3), and it is admitted for (15) and (16), because ϵ takes into account the collective response (polarization) of the nuclear matter. We have assumed that no color indices are attached to ϵ . It would correspond to the collective response of the color-neutral (on the average) medium if color exchange between the external current J^μ and medium excitations is numerous and averages to zero. The lack of knowledge of the collective excitations of the nuclear medium prevents more detailed studies. However, it seems to be justified at least for Cherenkov effects.

3 Cherenkov gluons as the classical lowest order solution of in-medium gluodynamics

Cherenkov effects are especially suited for treating them by classical approach to (15) and (16). Their unique feature is independence of the coherence of subsequent emissions on the time interval between these processes.

The problem of the coherence length for Cherenkov radiation was extensively studied [15, 16]. It was shown that the ω -component of the field of a current can be imitated by a set of oscillators with frequency ω situated along the trajectory. The waves from all oscillators add up in the direction given by the Cherenkov angle θ , independent on the length of the interval filled in by these oscillators. The lack of balance of the phase $\Delta\phi$ between emissions with frequency $\omega = k/\sqrt{\epsilon}$ separated by the time interval Δt (or the length $\Delta z = v\Delta t$) is given by

$$\Delta\phi = \omega\Delta t - k\Delta z \cos\theta = k\Delta z \left(\frac{1}{v\sqrt{\epsilon}} - \cos\theta \right) \quad (17)$$

up to terms that vanish for large distances between oscillating sources and the detector. For Cherenkov effects the angle θ is

$$\cos\theta = \frac{1}{v\sqrt{\epsilon}}. \quad (18)$$

The coherence condition $\Delta\phi = 0$ is valid independent of Δz . This is a crucial property specific for Cherenkov radiation only.³ Thus the change of color at the emission vertices is not important if one considers a particular a th component of color fields produced at the Cherenkov angle. Therefore, the fields (Φ_a, \mathbf{A}_a) and the classical current for in-medium gluodynamics can be represented by the product of their electrodynamic expressions (Φ, \mathbf{A}) and the color matrix T_a . As a result, one can neglect the ‘‘rotation’’ of color at emission vertices and use in the lowest order for Cherenkov gluons the well known formulae for Cherenkov photons just replacing α by $\alpha_s C_A$ for the gluon currents in the probabilities of their emission. Surely, there is radiation

³ The requirement for $\Delta\phi$ to be a multiple of 2π (or the weaker condition of being less than or of the order of 1) in the cases in which the Cherenkov condition is not satisfied imposes limits on the effective radiation length, as happens, e.g., for the Landau–Pomeranchuk or Ter–Mikaelyan effects.

at angles different from the Cherenkov angle (18). For such gluons one should take into account the coherence length and color rotation considering the corresponding Wilson lines [17].

Let us remind the reader of the explicit abelian solution for the current with velocity \mathbf{v} along the z -axis:

$$\mathbf{j}(\mathbf{r}, t) = \mathbf{v}\rho(\mathbf{r}, t) = 4\pi g\mathbf{v}\delta(\mathbf{r} - \mathbf{v}t). \quad (19)$$

In the lowest order the solutions for the scalar and vector potentials are related so that

$$\mathbf{A}^{(1)}(\mathbf{r}, t) = \epsilon\mathbf{v}\Phi^{(1)}(\mathbf{r}, t), \quad (20)$$

where the superscript (1) indicates the solutions of order g .

Therefore, the explicit expressions for Φ suffice. Using the Fourier transform, the lowest order solution of (11) with account of (19) can be cast in the form

$$\Phi^{(1)}(\mathbf{r}, t) = \frac{g}{2\pi^2\epsilon} \int d^3k \frac{\exp[i\mathbf{k}(\mathbf{r} - \mathbf{v}t)]}{k^2 - \epsilon(\mathbf{k}\mathbf{v})^2}. \quad (21)$$

The integration over the angle in cylindrical coordinates gives the Bessel function $J_0(k_\perp r_\perp)$. Integrating over the longitudinal component k_z with account of the poles due to the denominator⁴ and then over the transverse one k_\perp , one gets the following expression for the scalar potential [18]:

$$\Phi^{(1)}(\mathbf{r}, t) = \frac{2g}{\epsilon} \frac{\theta(vt - z - r_\perp\sqrt{\epsilon v^2 - 1})}{\sqrt{(vt - z)^2 - r_\perp^2(\epsilon v^2 - 1)}}. \quad (22)$$

Here $r_\perp = \sqrt{x^2 + y^2}$ is the cylindrical coordinate; z is the symmetry axis. The cone

$$z = vt - r_\perp\sqrt{\epsilon v^2 - 1} \quad (23)$$

determines the position of the shock wave due to the θ -function in (22). The field is localized within this cone. The Descartes components of the Poynting vector are related according to (22) and (20) by

$$S_x = -S_z \frac{(z - vt)x}{r_\perp^2}, \quad S_y = -S_z \frac{(z - vt)y}{r_\perp^2}, \quad (24)$$

so that the direction of the emitted gluons is perpendicular to the cone (23) and defined by the Cherenkov angle

$$\tan^2\theta = \frac{S_x^2 + S_y^2}{S_z^2} = \epsilon v^2 - 1, \quad (25)$$

which coincides with (18).

The higher order terms (g^3, \dots) can be calculated using (15) and (16).

The expression for the intensity of the radiation is given by the Tamm–Frank formula (up to Casimir operators)

$$\frac{dE}{dl} = 4\pi\alpha_s \int \omega d\omega \left(1 - \frac{1}{v^2\epsilon} \right). \quad (26)$$

⁴ These poles are at work only for Cherenkov radiation!

It is well known that it leads to infinity for constant ϵ . The ω -dependence of ϵ (its dispersion) usually solves the problem. For absorbing media ϵ acquires the imaginary part. The sharp front edge of the shock wave is smoothed. The angular distribution of Cherenkov radiation widens. The δ -function at the angle (18) and (25) is replaced by the Breit–Wigner shape [19] with maximum at the same angle (but $|\epsilon|$ in place of ϵ) and the width proportional to the imaginary part. Without absorption, the potential (22) is infinite on the cone. With absorption, it is finite everywhere except the cone vertex and is inverse proportional to the distance from the vertex. For low absorption, the field on the cone increases as $(\text{Im } \epsilon)^{-1/2}$ [39]. Absorption induces also longitudinal excitations (chromoplasmons), which are proportional to the imaginary part of ϵ and usually small compared to transverse excitations. The magnetic permittivity is easily taken into account replacing ϵ by $\epsilon\mu$ in the Breit–Wigner formula.

In electrodynamics the permittivity of real substances depends on ω . Moreover, it has the imaginary part determining the absorption. E.g., $\text{Re } \epsilon$ for water [20] is approximately constant in the visible light region ($\sqrt{\epsilon} \approx 1.34$), increases at low ω and becomes smaller than 1 at high energies, tending to 1 asymptotically. The absorption ($\text{Im } \epsilon$) is very small for visible light but dramatically increases in nearby regions both at low and high frequencies. Theoretically this behavior is ascribed to various collective excitations in water relevant to its response to radiation with different frequencies. Among them the resonance excitations are quite prominent (see, e.g., [21]). Even in electrodynamics, the quantitative theory of this behavior is still lacking, however.

Then, what can we say about the nuclear permittivity?

4 The nuclear permittivity

The partons constituting high-energy hadrons or nuclei interact during the collision for a very short time. Nevertheless, there are experimental indications that an intermediate state of matter (CGC, QGP, nuclear fluid, ...) is formed and evolves. Those are J/ψ -suppression, jet quenching, collective flow (v_2), Cherenkov rings of hadrons etc. They show that there is a collective response of the nuclear matter to color currents moving in it. Unfortunately, our knowledge of its internal excitation modes is very scarce, much smaller than in electrodynamics.

The attempts to calculate the nuclear permittivity from first principles are not very convincing. It can be obtained from the polarization operator. The corresponding dispersion branches have been computed in the lowest order perturbation theory [22–25]. The properties of collective excitations have been studied in the framework of the thermal field theories (for a review see, e.g., [26]). The results with an additional phenomenological ad hoc assumption about the role of resonances were used in a simplified model of scalar fields [8] to show that the nuclear permittivity can be larger than 1, which admits Cherenkov gluons. Extensive studies were performed in [27].

Let us stress the difference between these approaches and our considerations. In [22–26] the medium response to the *induced* current is analyzed. Namely, it determines the nuclear permittivity. The permittivity is an internal property of the medium. Its quantitative description poses problems even in QED. It becomes a more difficult task in QCD, where confinement is not understood. The quantitative description of the permittivity asks for solutions of à la Vlasov non-linear kinetic equations for the quark and gluon distribution functions (see, e.g., [28–30]). Therefore, we did not yet attempt to compute the nuclear permittivity and introduced it purely phenomenologically in analogy to in-medium electrodynamics. Our main goal is to study the medium response to the *external* color current. The Cherenkov effect is proportional to g^2 according to (22) if ϵ is constant or chosen purely phenomenologically. However, ϵ should tend to 1 at small g and the Cherenkov effect disappears. Thus, it is of the order of g^4 at small g . Mach waves in hydrodynamics [31] are of the same order. When the current $J^{(3)}$ is treated as an external one in the equations of in-vacuum gluodynamics [10–14], the effect is proportional to g^6 .

We prefer to use the general formulae of the scattering theory [32] to estimate the nuclear permittivity. It is related to the refractive index n of the medium:

$$\epsilon = n^2 \quad (27)$$

and the latter one is expressed [32] through the real part of the forward scattering amplitude of the refracted quanta⁵ $\text{Re } F(0^\circ, E)$ by

$$\begin{aligned} \text{Re } n(E) &= 1 + \Delta n_R = 1 + \frac{6m_\pi^3 \nu}{E^2} \text{Re } F(E) \\ &= 1 + \frac{3m_\pi^3 \nu}{4\pi E} \sigma(E) \rho(E). \end{aligned} \quad (28)$$

Here E denotes the energy, ν is the number of scatterers within a single nucleon, m_π the pion mass, $\sigma(E)$ the cross section and $\rho(E)$ the ratio of real to imaginary parts of the forward scattering amplitude $F(E)$. Thus the emission of Cherenkov gluons is possible only for processes with positive $\text{Re } F(E)$ or $\rho(E)$. Unfortunately, we are unable to calculate directly in QCD these characteristics of gluons⁶ and have to rely on analogies and our knowledge of the properties of hadrons. The only experimental facts we get for this medium are brought about by particles registered at the final stage. They have some features in common, which (one may hope!) are also relevant for gluons as the carriers of the strong forces. Those are the resonant behavior of amplitudes at rather low energies and the positive real part of the forward scattering amplitudes at very high energies for hadron–hadron and photon–hadron processes as measured from the interference of the Coulomb and hadronic parts of the amplitudes. $\text{Re } F(0^\circ, E)$ is always positive (i.e., $n > 1$)

⁵ In electrodynamics these quanta are photons, in QCD gluons.

⁶ We can only say that $\text{Re } F(E) \propto g^2$ at small g , which confirms the above estimates.

within the low-mass wings of the Breit–Wigner resonances. This shows that the necessary condition for Cherenkov effects, $n > 1$, is satisfied at least within these two energy intervals. This fact was used to describe experimental observations at SPS, RHIC and cosmic ray energies. The asymmetry of the ρ -meson shape at SPS [33] and azimuthal correlations of in-medium jets at RHIC [5, 34, 35] were explained by emission of comparatively low-energy Cherenkov gluons [36, 37]. The parton density and intensity of the radiation were estimated. In turn, cosmic ray data [38] at energies corresponding to LHC require very high-energy gluons to be emitted by the ultrarelativistic partons moving along the collision axis [6, 7]. Let us note the important difference from electrodynamics, where $n < 1$ at high frequencies. For QGP the high-energy condition $n > 1$ is a consequence of its instability.

The dispersion (ω -dependence of n) was taken into account. Otherwise the intensity of the radiation given by (26) diverges. It can easily be incorporated in (15) and (16) (more precisely, in their Fourier components). Equation (28), valid for $n - 1 \ll 1$, is generalized to the Lorenz–Lorentz expression for larger n . The imaginary part of ϵ can easily be accounted for. In principle, it may be estimated from RHIC data [39].

Up to now we did not discuss one of the most important problems of the coordinate system in which the permittivity is defined.

5 The rest system of the nuclear matter

The in-medium equations are not Lorentz-invariant. This is no problem in macroscopic electrodynamics, because the rest system of the macroscopic matter is well defined and its permittivity is considered there. For collisions of two nuclei (or hadrons) it requires special discussion.

Let us consider a particular parton that radiates in the nuclear matter. It would “feel” the surrounding medium at rest if the momenta of all other partons (or constituents of the matter), with which this parton can interact, sum to zero. In RHIC experiments the triggers, which registered the jets (created by partons) were positioned at 90° to the collision axis. Such partons should be produced by two initial forward–backward moving partons scattered at 90° . The total momentum of the other partons (medium spectators) is balanced, because for such a geometry the partons from both nuclei play the role of spectators forming the medium. Thus the center of mass system is the proper one to consider the nuclear matter at rest in this experiment. The permittivity must be defined there. The Cherenkov rings consisting of hadrons have been registered around the away-side jet, which traversed the nuclear medium. This geometry requires, however, high statistics, because the rare process of scattering at 90° has been chosen.

The forward (backward) moving partons are much more numerous and have higher energies. However, one cannot treat the radiation of such a primary parton in the c.m.s. in a similar way, because the momentum of the spectators is different from zero, i.e. the matter is not at

rest. Now the spectators (the medium) are formed from the partons of another nucleus only. Then the rest system of the medium coincides with the rest system of that nucleus and the permittivity should refer to this system. The Cherenkov radiation of such highly energetic partons must be considered there. That is what was done for interpretation of the cosmic ray event in [6, 7]. This discussion clearly shows that one must carefully define the rest system for other geometries of the experiment with triggers positioned at different angles.

Thus our conclusion is that the definition of ϵ depends on the geometry of the experiment. Its corollary is that partons moving in different directions with different energies can “feel” different states of matter in the *same* collision of two nuclei because of the dispersive dependence of the permittivity. The transversely scattered partons with comparatively low energies can analyze the matter with rather large permittivity corresponding to the resonance region, while the forward moving partons with high energies would “observe” a low permittivity in the same collision. This peculiar feature can help scan the $(\ln x, Q^2)$ -plane as it is discussed in [40]. It explains also the different values of ϵ needed for the description of the RHIC and cosmic ray data.

These conclusions can be checked at LHC, because both RHIC and cosmic ray geometry will become available there. The energy of the forward moving partons would exceed the thresholds above which $n > 1$. Then both types of experiments can be done, i.e. the 90° -trigger and non-trigger forward–backward partons experiments. The predicted results for 90° -trigger geometry are similar to those at RHIC. The non-trigger Cherenkov gluons should be emitted within the rings at polar angles of tens degrees in c.m.s. at LHC by the forward moving partons (and symmetrically by the backward ones). This idea is supported by some events observed in cosmic rays [37, 38].

6 Conclusions

Equations (15) and (16) of in-medium gluodynamics are proposed. They do remind one of the in-medium Maxwell equations with non-abelian terms added. Their lowest order classical solutions are similar (up to the trivial color factors) to those of electrodynamics (22), especially for Cherenkov gluons. The nuclear permittivity of the hadronic medium is related to the forward scattering hadronic amplitudes and its possible generalization is discussed. This definition requires the distinction between the different coordinate systems, in which the Cherenkov radiation (and nuclear permittivity) should be treated for partons moving in different directions with different energies.

This consideration has led to an explanation of several effects observed at SPS, RHIC, cosmic ray energies, and it predicts new features at LHC [40]. Some estimates of the properties of the nuclear matter formed in ultrarelativistic heavy-ion collisions have been made and are predicted.

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