Transverse Energy Density Fluctuations in the Color Glass Condensate Model

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We calculate the transverse correlation of fluctuations of the deposited energy density in nuclear collisions in the framework of the Gaussian color glass condensate model.

I. EVENT-BY-EVENT FLUCTUATIONS

The event-by-event fluctuations of the transverse emission pattern of hadrons in high-energy collisions of identical heavy nuclei have recently attracted much interest experimentally [1–7] and theoretically [8–22]. When averaged over collision events, the azimuthal angular distribution of emitted hadrons around the beam axis is symmetric with respect to the plane perpendicular to the impact parameter vector \( \mathbf{b} \). When averaged over collision events, the azimuthal angular distribution is peaked at zero; for odd \( n \), \( \psi_n \) is peaked around zero; for odd \( n \), \( \psi_n \) is randomly distributed. The dominant coefficient, \( v_2 \), is called elliptic flow.

Owing to quantum fluctuations in the density distributions of the colliding nuclei and finite particle number effects on the distribution of emitted particles, the left-right symmetry is broken in individual collision events. The angular distribution can then be written in the form

\[
\frac{dN}{d^2p_T} = \frac{dN}{\pi p_T^2} \left( 1 + \sum_{n=1}^{\infty} v_n(p_T) \cos(\phi_p) \right),
\]

(1)

where \( \phi_p \) is the angle between \( p_T \) and \( \mathbf{b} \), and \( \phi_p \) is therefore completely characterized by the even Fourier coefficients \( v_n \). The event averaged angular distribution

\[
\frac{dN}{d^2p_T} = \frac{dN}{\pi p_T^2} \left( 1 + \sum_{n=1}^{\infty} v_n(p_T) \cos(\phi_p + \psi_n) \right),
\]

(2)

where \( \psi_n \) describes the tilt angle of the “event plane” for each Fourier coefficient with respect to the reaction plane defined by the vector \( \mathbf{b} \). For even \( n \), \( \psi_n \) is peaked around zero; for odd \( n \), \( \psi_n \) is randomly distributed. The dominant odd coefficient, \( v_3 \), is known as triangular flow.

II. ENERGY DENSITY FLUCTUATIONS

In the Gaussian random source approximation to the color glass condensate model of small-\( x \) gluon structure of atomic nuclei [24, 25] the probability distribution of color charge density \( \rho^a(x) \) in the transverse plane is assumed to be of the form

\[
P[\rho^a] = \exp \left( -\frac{1}{g^2 \mu^2} \int d^2x \, \rho^a(x)^2 \right),
\]

(3)
Here $\mu^2$ represents the area density of color charges in the colliding nuclei, and $Q_s = g \mu$ is called the saturation scale, because it represents the scale at which the small–$x$ evolution of the gluon density becomes nonlinear due to saturation effects \[23\] \[27\]. Owing to the independent contributions of several nucleons to the color field, the Gaussian approximation is expected to provide a good description to the color source distribution in colliding nuclei at small $x$ \[23\]. In the light-cone gauge, the Gaussian color charge distribution translates into a Gaussian distribution of transverse gauge field strengths. Here we will follow the work of Lappi \[29\].

To calculate the initial state density fluctuations

$$\langle \varepsilon(x)\varepsilon(y) \rangle - \langle \varepsilon(x) \rangle \langle \varepsilon(y) \rangle,$$  

(4)

where $x, y$ denote vectors in the transverse plane, we start from the expression for the deposited energy density with the IR cut-off parameter $\Lambda$. It is convenient to decompose the momentum quadrupole tensor as follows:

$$\varepsilon(x) = \frac{1}{4} F_{ij}^c(x) F_{ij}^c(x) + 2 A^{\alpha c}(x) A^{\alpha c}(x)$$  

(5)

with transverse vector indices $i, j, m, n, ... = 1, 2$. Following the collision, the field strength tensor in the region between the receding nuclei only receives contributions from the mixed terms, as the color field of each individual nucleus is a pure gauge and the field strength tensor of each individual nucleus is thus zero outside the nuclear volume:

$$F_{ij}^{\alpha c}(x) = g f_{abc} \left( A_0^a(1;x) A_j^b(2;x) + A_0^b(2;x) A_j^a(1;x) \right)$$  

(6)

$$A^{\alpha c}(x) A^{\alpha c}(x) = \frac{g^2}{4} f_{abc} f_{\alpha' \beta' \gamma'} c$$  

(7)

Here “1” and “2” denote the gauge fields carried by nucleus 1 and 2, respectively. The field correlator in the color glass condensate model is given by

$$\langle A_i^a(n;x) A_j^b(m;y) \rangle = \frac{1}{2} \left( \langle A_i^a(n;x) A_j^b(m;y) \rangle + \langle A_j^b(m;y) A_i^a(n;x) \rangle \right) = \delta_{mn} \delta_{ab} \int \frac{d^2 p}{(2\pi)^2} \cos[p \cdot (x - y)] \frac{p_i p_j}{p^2} G(|p|)$$  

(8)

where $G(|p|)$ is the Fourier transform of the function

$$G(|x|) = g^2 N |x|^2 \left[ 1 - \exp \left( \frac{g^2 N g^2 \mu^2}{8\pi} \right) \right] \Theta(1 - \Lambda |x|)$$  

(9)

with the IR cut-off parameter $\Lambda$. It is convenient to decompose the momentum quadrupole tensor as follows:

$$p_i p_j = \frac{p_1^2 + p_2^2}{2} \delta_{ij} + \frac{p_1^2 - p_2^2}{2} \sigma^3_{ij} + p_1 p_2 \sigma^1_{ij},$$  

(10)

where $\sigma^1, \sigma^3$ are the familiar Pauli matrices. We thus obtain

$$\langle A_i^a(n;x) A_j^b(m;y) \rangle = \frac{1}{2} \delta_{mn} \delta_{ab} \int \frac{d^2 p}{(2\pi)^2} \left( \cos[p_1(x_1 - y_1)] \cos[p_2(x_2 - y_2)] \delta_{ij} G(|p|) + \cos[p_1(x_1 - y_1)] \sin[p_2(x_2 - y_2)] \sigma^3_{ij} \frac{p_1^2 - p_2^2}{p^2} G(|p|) - \sin[p_1(x_1 - y_1)] \sin[p_2(x_2 - y_2)] \sigma^1_{ij} \frac{p_1 p_2}{p^2} G(|p|) \right)$$

$$= \frac{1}{2} \delta_{mn} \delta_{ab} \left( \delta_{ij} D(x - y) + \sigma^3_{ij} E(x - y) - \sigma^1_{ij} F(x - y) \right) = \delta_{mn} \delta_{ab} S_{ij}(x - y).$$  

(11)

For later use, we will note the values of the individual correlation functions $D, E, F$ at the origin:

$$D(0) = \int \frac{d^2 p}{(2\pi)^2} G(|p|) = \lim_{|x| \to 0} G(|x|); \quad E(0) = F(0) = 0.$$  

(12)

The expression for $D(0)$ diverges logarithmically for the function $G(|x|)$ given in \[29\], if the gauge coupling $g$ is taken as a constant. However, as pointed out by Kovchegov and Weigert \[30\], the infrared divergence can be removed by including effects from the running of the coupling constant by means of the substitution

$$g^4 \longrightarrow g^2(\mu^2) g^2(1/|x|^2)$$  

(13)
in the exponent of \[ 9 \]. The specific structure of this substitution, sometimes called the “triumvirate” structure of the running coupling, is motivated by the form of next-to-leading order corrections to the small–\( x \) evolution of the BFKL kernel in the color dipole approach to parton saturation \[ 31 \].

We first evaluate the expectation value of the deposited energy density:

\[
\langle \varepsilon(x) \rangle = \frac{g^2}{2} f_{abc} f_{a'b'c'} \left\langle A^a_i(1; x) A^b_j(2; x) A^{a'}_i(1; x) A^{b'}_j(2; x) + A^a_i(1; x) A^b_j(2; x) A^{a'}_i(1; x) A^{b'}_j(2; x) \right\rangle \\
+ \frac{g^2}{2} f_{abc} f_{a'b'c'} \left\langle A^a_i(1; x) A^b_j(2; x) A^{a'}_i(1; x) A^{b'}_j(2; x) \right\rangle \\
= \frac{g^2}{2} N (N^2 - 1) D^2(0),
\]

recovering the result given in Eq. (14) of ref. \[ 29 \].

Next we evaluate the two-point correlator of the energy density:

\[
\langle \varepsilon(x) \rangle \langle \varepsilon(y) \rangle = \frac{g^4}{4} f_{abc} f_{a'b'c'} f_{efd} f_{e'f'd'} \left\langle \left( A^a_i(1; x) A^b_j(2; x) A^{a'}_i(1; x) A^{b'}_j(2; x) \right) + A^a_i(1; x) A^b_j(2; x) A^{a'}_i(1; x) A^{b'}_j(2; x) \right\rangle \\
x \left( A^a_m(1; y) A^b_n(2; y) A^{a'}_m(1; y) A^{b'}_n(2; y) \right) + A^a_m(1; y) A^b_n(2; y) A^{a'}_m(1; y) A^{b'}_n(2; y) \right\rangle \\
= \frac{g^4}{4} N (N^2 - 1) D^2(0, 1),
\]

We again make use of the fact that only correlators among fields in the same nucleus are non-zero, which allows us to suppress the labels 1 and 2:

\[
\langle \varepsilon(x) \rangle \langle \varepsilon(y) \rangle = \frac{g^4}{4} f_{abc} f_{a'b'c'} f_{efd} f_{e'f'd'} \left\langle \left( A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right) + A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right\rangle \\
+ \left\langle A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right\rangle \left\langle A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right\rangle \\
+ \left\langle A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right\rangle \left\langle A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right\rangle \\
+ \left\langle A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right\rangle \left\langle A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right\rangle \\
+ \left\langle A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right\rangle \left\langle A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right\rangle \\
= \frac{g^4}{4} N (N^2 - 1) D^2(0, 1),
\]

In the spirit of the Gaussian approximation we now factorize the correlators of four gauge fields into products of correlators among two gauge fields \[ 35 \], e. g.:

\[
\left\langle A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right\rangle = \left\langle A^a_i(x) A^b_j(x) \right\rangle \left\langle A^{a'}_m(y) A^{b'}_m(y) \right\rangle + \left\langle A^a_i(x) A^{a'}_m(y) \right\rangle \left\langle A^b_j(x) A^{b'}_m(y) \right\rangle \\
+ \left\langle A^a_i(x) A^{a'}_m(y) \right\rangle \left\langle A^b_j(x) A^{b'}_m(y) \right\rangle
\]

To proceed further we use the symmetry with respect to the color indices \( e' \) and \( f' \) to combine, e. g., the second and third term in the large brackets of \[ 16 \]:

\[
M_2 + M_3 = \left\langle A^a_i(x) A^b_j(x) A^{a'}_m(y) A^{b'}_m(y) \right\rangle \left\langle A^a_i(x) A^b_j(x) \left( A^{a'}_m(y) A^{b'}_m(y) - A^{b'}_m(y) A^{a'}_m(y) \right) \right\rangle
\]
The second factor is easily shown to vanish:

\[
\langle \cdots \rangle = \left\langle A^f_i(x)A^g_n(y)\right\rangle \left\langle A^h_j(x)A^i_m(y)\right\rangle \left\langle A^j_k(x)A^k_n(y)\right\rangle - (m \leftrightarrow n)
\]

\[
= \delta_{ef}\delta_{bd}'S_{jm}(x-y)S_{jm}(x-y) + \delta_{ef'}\delta_{bd}S_{jm}(x-y)S_{jm}(x-y) - (m \leftrightarrow n) = 0
\]

(19)

The same holds true for the forth and seventh term, \(M_4\) and \(M_7\). After considerable algebra, the fifth, sixth, eighth and ninth terms combine to

\[
M_5 + M_6 + M_8 + M_9 = \frac{g^4}{4}f_{abc}f_{a'd'b}f_{efd}f_{ef'd'} \left\langle A^a_i(x)A^b_j(x)A^c_n(y)A^d_n(y)\right\rangle
\]

\[
\times \left\langle A^i_j(x)A^j(x) - A^i_j(x)A^j(x)\right\rangle \left\langle A^f_i(x)A^g_n(y) - A^f_i(x)A^g_n(y)\right\rangle
\]

\[
= \frac{g^4}{4}f_{abc}f_{a'd'b}f_{efd}f_{ef'd'} \left( \delta_{ae}\delta_{a'd}'S_{im}S_{jm} + \delta_{af}\delta_{a'd}S_{im}S_{jm} \right)
\]

\[
\times \left( \delta_{bf}\delta_{b'd}2[S_{im}S_{jm}S_{im}S_{jm} + \delta_{be}\delta_{b'd}'2[S_{im}S_{jm} - S_{im}S_{jm}]] \right)
\]

\[
= \frac{g^4}{16}N^2(N^2 - 1) \left[ D(x-y)^2 + E(x-y)^2 + F(x-y)^2 \right]^2,
\]

(20)

where we made use of the relation which follows from the Jacobi identity.

\[
f_{abc}f_{a'd'b}f_{d'bd} = \frac{N}{2}f_{ab'd}
\]

(21)

Finally we evaluate the first term:

\[
M_1 = \frac{g^4}{4}f_{abc}f_{a'd'b}f_{efd}f_{ef'd'} \left\langle A^a_i(x)A^b_j(x)A^c_n(y)A^d_n(y)\right\rangle
\]

\[
\times \left\langle A^i_j(x)A^j(x)A^j(x)A^d_n(y)A^f_i(x)A^g_n(y)\right\rangle
\]

\[
= \frac{g^4}{4}N^2(N^2 - 1)^2D(0)^4
\]

\[
+ \frac{g^4}{2}N^2(N^2 - 1)D(0)^2 \left[ D(x-y)^2 + E(x-y)^2 + F(x-y)^2 \right]
\]

\[
+ \frac{3g^4}{8}N^2(N^2 - 1) \left[ D(x-y)^2 + E(x-y)^2 + F(x-y)^2 \right]^2.
\]

(22)

Combining these equations we finally obtain

\[
\langle \varepsilon(x)\varepsilon(y) \rangle = \langle \varepsilon(x) \rangle \langle \varepsilon(y) \rangle
\]

\[
= \frac{g^4}{2}N^2(N^2 - 1)D(0)^2 \left[ D(x-y)^2 + E(x-y)^2 + F(x-y)^2 \right]
\]

\[
+ \frac{7g^4}{16}N^2(N^2 - 1) \left[ D(x-y)^2 + E(x-y)^2 + F(x-y)^2 \right]^2.
\]

(23)

Since the functions \(D, E, F\) always appear in the same combination, it makes sense to introduce the abbreviation

\[
K(x-y) = D(x-y)^2 + E(x-y)^2 + F(x-y)^2,
\]

(24)

In terms of which the average deposited energy density and its fluctuation can be expressed as:

\[
\varepsilon_0 = \langle \varepsilon \rangle = \frac{g^2}{2}N(N^2 - 1)K(0);
\]

\[
\Delta \varepsilon(x-y)^2 = \langle \varepsilon(x)\varepsilon(y) \rangle - \langle \varepsilon(x) \rangle \langle \varepsilon(y) \rangle = \frac{g^4}{2}N^2(N^2 - 1) \left[ K(0)K(x-y) + \frac{7}{8}K(x-y)^2 \right].
\]

(26)
III. EVALUATION OF INTEGRALS

Next we simplify the integrals $D(x - y)$, $E(x - y)$ and $F(x - y)$. We abbreviate $z = x - y$ and $z = |z|$. We begin with $D(z)$.

$$D(z) = \int_0^\infty \frac{dp}{4\pi^2} G(p) \int_0^{2\pi} d\phi \cos(pz_1 \cos \phi) \cos(pz_2 \sin \phi)$$

$$= \int_0^\infty \frac{dp}{4\pi^2} G(p) \int_0^{2\pi} d\phi \left( \cos[p(z_1 \cos \phi + z_2 \sin \phi)] + \cos[p(z_1 \cos \phi - z_2 \sin \phi)] \right)$$

We substitute $\phi \to \pi - \phi$ in the last term and introduce the notation $z_1 = z \cos \psi$, $z_2 = z \sin \psi$:

$$D(z) = \int_0^\infty \frac{dp}{4\pi^2} G(p) \int_0^{2\pi} d\phi \cos[pz \cos(\phi - \psi)]$$

$$= \int_0^\infty \frac{dp}{4\pi^2} G(p) J_0(pz) = \int_0^\infty \frac{dp}{4\pi^2} G(p) e^{ipz} \equiv G(z),$$

where we have used Eq. (3.715.18) from [32]. Similarly we obtain

$$E(z) = \int_0^\infty \frac{dp}{4\pi^2} G(p) \int_0^{2\pi} d\phi \cos(pz_1 \cos \phi) \cos(pz_2 \sin \phi) \cos^2 \phi - \sin^2 \phi$$

$$= \int_0^\infty \frac{dp}{4\pi^2} G(p) \int_0^{2\pi} d\phi \left( \cos[p(z_1 \cos \phi + z_2 \sin \phi)] + \cos[p(z_1 \cos \phi - z_2 \sin \phi)] \right) \cos(2\phi).$$

Using the same substitutions we find:

$$E(z) = -\cos(2\psi) \int_0^\infty \frac{dp}{4\pi^2} G(p) J_2(pz)$$

where we used Eqs. (3.715.18) and (3.715.7) from [32]. Finally, a similar calculation yields:

$$F(z) = \int_0^\infty \frac{dp}{4\pi^2} G(p) \int_0^{2\pi} d\phi \sin(pz_1 \cos \phi) \sin(pz_2 \sin \phi) \cos \phi \sin \phi$$

$$= \sin(2\psi) \int_0^\infty \frac{dp}{2\pi} G(p) J_2(pz).$$

We conclude that the function $K(z)$ only depends on the distance $z$ between the points $x$ and $y$.

We can express $D(z)$ and $E(z)^2 + F(z)^2$ in terms of the following integrals:

$$C_n(z) = \int_0^\infty \frac{dp}{2\pi} G(p) J_n(pz),$$

for $n = 0, 2$, namely:

$$D(z) = C_0(z);$$

$$E(z)^2 + F(z)^2 = C_2(z)^2.$$ We have to evaluate integrals of the type

$$B_n(x, z) = \int_0^\infty \frac{dp}{2\pi} J_0(px) J_n(pz).$$

This integral can be evaluated for $n = 0$ using formula (6.633.2) in [32]:

$$\int_0^\infty \frac{dp}{2\pi} e^{-c^2 p^2} J_0(px) J_0(pz) = \frac{e^{-x^2 + z^2/4c^2}}{2c^2} I_0 \left( \frac{xz}{2c^2} \right).$$

We are interested in the limit $c \to 0$, which means that we can apply the limit of $I_0(z)$ for large arguments:

$$I_0(z) \to \frac{e^z}{\sqrt{2\pi z}}.$$
This yields

\[ B_0(x, z) = \lim_{c \to 0} \frac{1}{4\pi c} \frac{1}{\sqrt{\pi xz}} \exp \left( -\frac{(x-z)^2}{4c^2} \right) \]

\[ = \frac{1}{2\pi z} \delta(x-z). \]  \hspace{1cm} (39)

For \( n = 2 \) we use the recursion relation for Bessel functions:

\[ J_2(z) = \frac{2}{z} J_1(z) - J_0(z), \]  \hspace{1cm} (40)

and apply formula (6.512.3) from [32]:

\[ \int_0^\infty dp J_0(px) J_1(pz) = \frac{1}{z} \theta(z-x), \]  \hspace{1cm} (41)

with the convention \( \theta(0) = 1/2 \). This implies:

\[ B_2(x, z) = \frac{1}{\pi z^2} \theta(z-x) - \frac{1}{2\pi z} \delta(x-z). \]  \hspace{1cm} (42)

When we insert these results into the desired integrals, we find:

\[ C_0(z) = G(z); \]  \hspace{1cm} (43)

\[ C_2(z) = \frac{2}{z^2} \int_0^z x \, dx \, G(x) - G(z). \]  \hspace{1cm} (44)

**IV. NUMERICAL RESULTS**

We now evaluate the average energy density and its fluctuations for a choice of the parameters that is motivated by the initial conditions at which thermal QCD matter is formed in heavy ion collisions at RHIC and LHC:

\[ Q_s^2 = (g^2\mu)^2 = 2 \text{ GeV}^2; \]

\[ g^2(\mu^2) = 3.785; \]

\[ g^2(1/x^2) = \frac{16\pi^2}{9\ln(1/(\Lambda^2 x^2))}. \]  \hspace{1cm} (45)

Note that the result is independent of the value of \( \Lambda \).

The functions \( C_0(z) \) and \( C_2(z) \) are shown in Figs. 1 and 2 for these parameter values.

For these parameters, the initial value of the deposited energy density is \( \varepsilon_0 \approx 240 \text{ GeV/fm}^3 \). This very large energy density quickly decreases due to the longitudinal expansion and reaches much smaller values by the time of thermalization. What matters for us is not the absolute value of the initial energy density, but the relative size and spatial correlation of its fluctuations, \( \Delta \varepsilon(x-y)/\varepsilon_0 \).

This function is shown in Fig. 3 for the parameters listed above.

As the figure shows, the fluctuations of the initial energy density are locally of similar magnitude as the energy density itself and fall over distances of the inverse saturation scale, here assumed as \( Q_s^{-1} \approx 0.14 \text{ fm} \). This result is in accord with the intuitive picture of the field configuration immediately after the collision in the color glass condensate model, as a bundle of longitudinally stretching random color flux tubes with characteristic transverse width \( 1/Q_s \).
V. SUMMARY

We have calculated the initial energy density fluctuations in high-energy heavy ion collisions within the Gaussian color glass condensate model. These turn out to be very large with a transverse profile determined by the saturation scale $Q_s$. A finite result is only obtained when the “triumvirate” running coupling is used, giving additional support to the correctness of Eq. (13). The fluctuation probabilities thus derived can serve as input for any calculation aiming at the investigation of early fluctuations, in particular for calculations which study the fate of such fluctuations during thermalization. For example, it is possible to investigate the problem of event-by-event fluctuations in heavy ion collisions within the AdS/CFT paradigm using methods similar to those employed in [33].

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