Local energy-momentum conservation and initial conditions for quark-gluon plasma evolution in nucleus-nucleus collisions at SPS and RHIC BES energies

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Abstract

We investigate consequences of local energy-momentum conservation for initial eccentricity coefficients in heavy ion collisions at not too high energies relevant for CERN SPS and RHIC BES. Different models of energy density available for pion production are considered. We study dependence of eccentricity coefficients on space-time rapidity for different impact parameters. The naive formula how to define eccentricities breaks with our initial conditions and must be corrected. We predict considerable eccentricities for \( \epsilon_{1,2,3,4} \) and specific dependence on space-time rapidity as well as on impact parameter for lower energies. The effect becomes smaller at larger energies when restricting to narrow rapidity interval around zero. Our predictions are in principle input for further hydrodynamical evolution but it is not clear whether they can be easily used. Our initial condition suggest a strong preequilibrium phase which is difficult for modelling.

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I. INTRODUCTION

The initial conditions for hydrodynamical evolution of quark-gluon plasma is of crucial importance for final observables such as rapidity distributions of the produced pions or flow parameters and their dependence of particle (pseudo)rapidity and collision centrality. So far mostly Glauber-collision or color-glass-condensate initial conditions were most often used in this context [1–4]. They lead to more-or-less correct behaviour of the flow at midrapidities which is of special interest for plasma studies.

In Ref.[5] we suggested that the conservation of energy and momentum in local longitudinal streaks, dependent on the position in the transverse momentum plane \(\vec{\rho} = (x, y)\), provides strong constraints on initial conditions for plasma evolution in heavy ion collisions. In this model the produced plasma in peripheral collisions is not at rest and its movement depends on the position in the impact parameter space \(\vec{\rho} = (x, y)\). In particular, the parts of plasma that are close to spectators move with velocities only slightly smaller than velocities of spectators while those for \(\vec{\rho}(0, 0)\) are at rest on average. Such initial conditions were used recently in studies of electromagnetic effects caused by strong fields generated by fast moving spectators and a good description of its Feynman-\(x_F\) dependence was obtained [6] in agreement with the NA49 data on the \(\pi^+ / \pi^-\) ratio at low pion transverse momenta [7]. Similar idea was studied recently in Ref.[8] where so-called directed flow of pions and protons was studied. The two approaches differ in defining energy available for particle production which will be discussed here in detail.

The important ingredient of the modelling of pion emission in [5] was fragmentation function, assumed universal for different impact parameters. There the fragmentation function parametrizes both nuclear transparency, evolution of plasma and fluid-to-hadron transition in an effective way, so it must be collision energy dependent. The fragmentation function is then adjusted to experimental data. In [8] the transparency or intial plasma distribution is parametrized and the fluid-to-hadron transition is a matter of a procedure applied after termination of the hydrodynamical evolution, not explained in extenso in [8]. The latter may be not well under control for the specific initial conditions for plasma evolution. It is worth to mention in this context that the production of forward/backward hadrons via fragmentation function is not well understood even for proton-proton collisions [9]. It becomes even more difficult for heavy hadron production [10].

Here we wish to investigate how the local (in \((x,y)\)) energy-momentum conservation as proposed in [5] influences/determines the initial eccentricity parameters of plasma created during a peripheral collision as a function of space-time rapidity.

II. LOCAL ENERGY-MOMENTUM CONSERVATION, ENERGY AVAILABLE FOR PARTICLE EMISSION AND ECCENTRICITIES OF THE PRODUCED PLASMA

For peripheral collisions of identical nuclei (Pb + Pb (SPS) or Au + Au (RHIC)) the full stop of nuclear matter is not always possible and energy-momentum conservation implies some local collective longitudinal flow of parts of the plasma. As discussed in [5] the tips of the almond-like initial fire-ball move very forward or very backward with large velocities, so the initial almond-like shape changes quickly with time and after a while does not remain the almond shape. Such a picture, very different from previously used initial conditions (Glauber or KLN) must have consequences for rapidity-dependent
observables such as energy dependence on rapidity and Fourier coefficients of azimuthal correlations.

Let us consider for a moment full stopping of nuclear matter i.e. no nuclear transparency. The full stopping for peripheral collisions does not mean that the plasma is at rest. The rapidity of the fully stopped plasma at the \((x, y)\) point for the collision at impact parameter \(b\) can be calculated as

\[
y_{\text{stop}}(x, y; b) = 2 \arctanh \left( \frac{T_A(x, y; b) - T_B(x, y; b)}{T_A(x, y; b) + T_B(x, y; b)} \tanh(y_{\text{beam}}) \right).
\]

The so-called thickness functions \(T_A\) and \(T_B\) are calculated here based on nucleon (proton and neutron) distributions obtained from Hartree-Fock-Bogoliubov method \([12]\) which includes neutron skin effects. \(y_{\text{stop}}\) above can be interpreted as rapidity of the plasma contracted to a point in \(z\) direction, of course different for different \((x, y)\). This is a textbook example of fully inelastic collision. In this approach the plasma moves with different rapidities for different position in the \((x, y)\) space and different impact parameter \(b\).

The distribution of plasma in space-time rapidity with respect to \(y_{\text{stop}}(x, y; b)\), rapidity of the fully stopped (no nuclear transparency) plasma at the point \((x, y)\) cannot be calculated at present and must be parametrized. \(^1\)

As a first option we consider a Gaussian parametrization of the plasma-like not equilibrated matter:

\[
f(\eta_s; x, y, b) = \frac{1}{\sqrt{2\pi\sigma_\eta^2}} \exp \left( \frac{(\eta_s - y_{\text{stop}}(x, y; b))^2}{2\sigma_\eta^2} \right). \tag{2.2}\]

A rectangular step-like function with the width \(d_\eta\) could be another option:

\[
f(\eta_s; x, y, b) = \text{const} = 1 / d_\eta \tag{2.3}\]

for \(\eta_s > y_{\text{stop}}(x, y; b) - d_\eta/2\) and \(\eta_s < y_{\text{stop}}(x, y; b) + d_\eta/2\) and zero otherwise. The \(d_\eta\) is a free parameter for the second parametrization. For simplicity we have assumed universal \((x,y;b)\)-independent functions centered at \(y_{\text{stop}}(x, y; b)\). At even higher energies (RHIC) a double-Gaussian shape can be considered \([11]\).

In the following we shall call \(f(\eta_s; x, y, b)\) “transparency function” for brevity, in spite of the fact that this name may be slightly confusing, and in order to distinguish from “fragmentation function”, also somewhat confusing, used in \([5]\).

The authors of Ref. \([8]\) proposed that the (total) energy density distribution for a given \((x, y)\) transverse position point can be calculated as:

\[
E(x, y; b) = (T_A(x, y; b) + T_B(x, y; b)) m_N \cosh(y_{\text{beam}}) \Big|_C.
\]

Please note the letter \(C\) in the formula above which means an extra condition which must be imposed to eliminate spectator matter from the participant matter \(^2\). The extra collision condition \(C\) written above can be formulated e.g. as:

\[
M_A(x, y) > m_0 \text{ and } M_B(x, y) > m_0, \tag{2.5}\]

---

\(^1\) There are first trials to calculate baryon stopping, see e.g. \([14]\). The relevant experimental data were obtained some time ago by e.g. the BRAHMS collaboration at RHIC \([15]\).

\(^2\) It is not clear to us whether it was included in \([8]\).
where \(M_A\) and \(M_B\) are cold mass densities in nucleus \(A\) and \(B\), respectively and \(m_0\) is a minimal (mass) density required for the creation of plasma (a free parameter of the model here). Such a condition assures that there must be some minimal nuclear matter portion in both nuclei to create plasma. Ignoring such a condition would mean “including spectators” as emitters of particles. A precise formulation of the collision condition \(\mathcal{C}\) is not easy close to spectators where the division into spectator and participant zone takes place. This condition has sizeable effect on initial eccentricities as will be discussed in the following.

For comparison the cold mass taking part in the collision can be written as:

\[
M(x, y; b) = (T_A(x, y; b) + T_B(x, y; b)) \, m_N|\mathcal{C} .
\]  

(2.6)

How the energy is distributed in space-time rapidity requires extra modelling. The extra distribution in space-time rapidity can be formally written as

\[
\frac{dE}{d\eta_s} = f(\eta_s; x, y, b)E(x, y)
\]

(2.7)

using the phenomenological functions \(f(\eta_s; x, y, b) = f(\eta_s; y_{stop})\) (Gaussian, step-like) discussed above.

The space-time rapidity distribution of energy contained in plasma for a given impact parameter \(b\) can be calculated as:

\[
\frac{dE}{d\eta_s}(\eta_s; b) = \int \frac{dE}{d\eta_s dxdy}(x, y; \eta_s, b) \, dxdy .
\]

(2.8)

In our opinion not the whole energy density as is written above is available for particle production.

In [5] we suggested to subtract from the total energy density \(E(x, y; b)\) at least the cold mass density \(M(x, y; b)\). The cold mass must reappear in the form of produced baryons due to baryon-number conservation so corresponding energy cannot be used to production of mesons (mostly pions). Such a correction obviously depends on collision energy and is larger percentage-wise for lowest energies where plasma is produced. Is this correction enough?

A sizeable part of the energy which is not available to particle production is a kinetic energy of parts of plasma moving in forward \((\eta_s > 0)\) or backward \((\eta_s < 0)\) directions. How to calculate the unavailable part of energy to be subtracted from the total energy for different bins in \(x\) and \(y\) and a given space-time rapidity \(\eta_s\) \((b\) is treated here as a parameter) \? A big part of it is contained in final moving(!) baryons. Including such a correction is straightforward. The available energy can be then obtained as:

\[
E_{avail}(x, y; b) = E(x, y; b) - M(x, y; b) \cdot \cosh y_{stop}(x, y; b)|\mathcal{C} > 0 .
\]

(2.9)

The \((x, y)\) profile of the newly defined quantity \(E_{avail}(x, y; b)\) is obviously different than that for total energy density \(E(x, y; b)\). In Ref. [8] the subtraction term above was not included. It may be crucial which energy is used as initial condition for plasma hydrodynamical evolution. We shall discuss consequences of the correction (not included in [8]) both on particle production and on eccentricities in the following.
To illustrate the global situation related to energy available in the collision or available for the particle production we define the integrated energy densities:

\[ E(b) = \int dx dy E(x, y; b), \]
\[ E_{\text{avail}}(b) = \int dx dy E_{\text{avail}}(x, y; b). \quad (2.10) \]

Very interesting is the energy distribution per space-time rapidity unit for a given impact parameter (centrality) which is strongly related to particle rapidity distribution. Formally it can be defined as:

\[ \frac{dE}{d\eta_s}(\eta_s; b) = \frac{\int dE(x, y; b) d\eta_s}{d\eta_s} dx dy, \quad (2.11) \]
\[ \frac{dE_{\text{avail}}}{d\eta_s}(\eta_s) = \frac{\int dE_{\text{avail}}(x, y; b) d\eta_s}{d\eta_s} dx dy. \quad (2.12) \]

The eccentricities of produced plasma, defined usually for midrapidities, read (see e.g. \[16\]):

\[ \epsilon_n = -\frac{\int \exp(+i\phi) \rho^n E(\rho, \phi) d^2\rho}{\int \rho^n E(\rho, \phi) d^2\rho}, \quad (2.13) \]

where \( \rho = \sqrt{x^2 + y^2} \) and \( \phi \) is azimuthal angle with respect to the reaction plane. Is this definition universal and valid also for forward/backward space-time rapidities? We shall return to the issue in the next section when discussing energy density as a function of \( (x, y, \eta_s) \).

In the next section we shall compare our results for \( \epsilon_n \) for the two proposed parametrizations of energy distribution in \( \eta_s \) and different ways of weighting the energy available for meson production.

### III. NUMERICAL RESULTS AND DISCUSSION

We wish to start our discussion by showing the collective longitudinal motion of plasma created in the high-energy collision. In Fig.1 we show dependence of rapidity of the stopped plasma as a function of \( x \) for fixed \( y = 0 \) for two different values of impact parameter \( b = 2, 8 \text{ fm} \). We wish to note the difference of our result (solid line) compared to that from \[8\] (dashed line). We predict that the pieces of plasma that are close to spectators have almost the same space-time rapidities as spectators. This is universal feature independent of collision energy.

The corresponding \( (x, y) \)-dependent energy density is shown in Fig.2. We show the results obtained using the naive formula (2.4) (dashed line) and the energy corrected for the collective longitudinal energy which cannot be used up for particle production. There is a significant effect for positions in \( x \) close to spectators. This means a reduced production of particles (pions) from this region of configuration space. This means in a hidden way also earlier creation of mesons (pions) from these parts of plasma. This must have direct consequences for electromagnetic effect on pions generated by fast moving spectators. The earlier produced pions feel the EM field of spectators somewhat earlier.
FIG. 1: Rapidity of the stopped plasma $y_{\text{stop}}$ as a function of $x$ for $y = 0$ and two different values of impact parameter $b = 2$ (black lines), 8 (red lines) fm. The solid lines are in agreement with [5] while the dashed lines are obtained according to [8].

FIG. 2: Energy density calculated for the same geometrical situation as in the previous figure for $b = 8$ fm collision. The dashed line represents result obtained from Eq.(2.4) while the solid line takes into account the subtraction of the collective longitudinal energy, which is useless for particle production.
FIG. 3: Dependence of integrated energies defined by (2.10) on impact parameter for $\sqrt{s_{NN}} = 17.3$ GeV. The $E(b)$ is shown as a the dashed line, while $E_{\text{avail}}(b)$ as the solid line and their difference as the dotted line.

Now we wish to present influence of the longitudinal flow on integrated quantities defined in (2.10). We observe that the energy available for particle production (solid line) is somewhat smaller, about 10 %, than the total energy in the collision (dashed line). The relative global effect only weakly depends on impact parameter. So the effect for minimum bias collisions is rather small.

The distribution of energy per unit of space-time rapidity is shown in Fig.4 for two different impact parameters $b = 2$ fm (left panel) and $b = 8$ fm (right panel). We show distribution of $dE/d\eta_s$ (dashed line) and $dE_{\text{avail}}/d\eta_s$ (solid line). The shape of the distribution should remind the shape of the rapidity distribution of pions which is known for $\sqrt{s_{NN}} = 17.3$ GeV. As in Ref. [5] we get slightly broader distribution in $\eta_s$ for peripheral collisions. The $\sigma_\eta = 1.0$ give more or less realistic distributions. $\sigma_\eta = 0.3$ from [8] gives in our opinion too narrow distribution and should be increased.

Now we will proceed to eccentricity coefficients. In Fig.5 we collected the dependences of the lowest eccentricities as a function of space-time rapidity for different impact parameters for $Pb + Pb$ collisions at $\sqrt{s_{NN}} = 17.3$ GeV relevant for the NA49 experiment. Here we have used $\sigma_\eta = 1.0$ as adjusted roughly to experimental data on pion rapidity distributions. The eccentricity coefficients calculated according to formula (2.13) depend strongly on rapidity. The even coefficients depend also strongly on the impact parameter. The so-calculated eccentricity coefficients contain spurious shift of the forward/backward moving matter from the (0,0) point.

In Fig.6 we present energy available for particle emission as a function of $(x, y)$ for the Gaussian distribution in $\eta_s$ (see Eq. (2.2)), for fixed values of $\eta_s = -2, -1, 0, 1, 2$ with the $\eta_s$-window width $\pm 0.1$. The shapes for larger $\eta_s$ become more eccentric and naturally should have the whole spectrum of eccentricity coefficients. Pieces of plasma with large $|\eta_s|$ have
funny shapes in \((x,y)\). Equilibration of plasma with such strange shapes cannot be fast, or does not occur at all as the remote parts do not communicate between themselves due to causality. Here naturally preequilibrium evolution must be at work.

We clearly see that \(x=0, y=0\) point is inadequate for defining eccentricity coefficients for forward or backward going plasma. In this case one should find more adequate reference point for shifted \(x\), centered at the middle of the plasma at a given space-time rapidity.

In our approach the situations shown in upper left panel and upper right panel as well as in lower left and lower right panels are correlated by the geometry. The parts of plasma going in forward and backward directions are correlated by geometry in the configuration space. This correlation translates (via preequilibrium and/or hydro evolution) to correlation in the momentum space which leads to forward-backward azimuthal correlations. Such an effect was discussed first in [21] in a different approach. We do not know if this was observed experimentally. This is also related to forward-backward multiplicity correlations which in our case is strongly related to the energy available for particle emission. Such correlations have been observed experimentally [17, 18]. However, in this case fluctuations are crucial.

The standard (hidden) reference point used for midrapidities \(x = 0, y = 0\) (see Eq.(2.13)) must be modified to \(x = x_{\text{shift}}, y = 0\) for nonzero \(\eta_s\).  

\[
x_{\text{shift}}(\eta_s; b) = \frac{\int E(x,y;\eta_s,b) x \ dx \ dy}{\int E(x,y;\eta_s,b) \ dx \ dy}.
\]  

Above \(x_{\text{shift}}\) is weighted by energy density \((E\) or \(E_{\text{avail}}\)) at a given space-time rapidity.  

\[3\] We use similar procedure as proposed in [8] but our weights seem different.
FIG. 5: Initial eccentricity parameters as a function of space-time rapidity calculated from formulae (2.13) for $\sqrt{s_{NN}} = 17.3$ GeV for different values of impact parameter $b = 2, 4, 6, 8$ fm.

slice. Then also $\rho$ and $\phi$ in (2.13) must be modified as:

$$\rho \rightarrow \tilde{\rho} = \sqrt{(x - x_{\text{shift}})^2 + y^2}, \quad (3.2)$$
$$\phi \rightarrow \tilde{\phi}, \quad (3.3)$$

where $\tilde{\phi}$ is a new azimuthal angle in the shifted frame of reference placed at $(x_{\text{shift}}, 0)$. $x_{\text{shift}}$ is of course defined with the same weight ($E$ or $E_{\text{avail}}$) as used in the definition of $\epsilon_n$.

We shall consider also another option trying to subtract the unavoidable local longitudinal flow as discussed in the previous section.

The shift in $x$ is shown in Fig. 7 as a function of space-time rapidity for different values of the impact parameter. The actual dependence is a consequence of the parametrization (2.2) of energy density as a function of $\eta_s$ which is not known from first principles. The shift for large $\eta_s$ is really large, independently of the parametrization of the transparency function. The transparency function related to bayon-stopping mechanism is therefore crucial for further results. In general, the function could be adjusted to reproduce rapidity dependent observables. Some of them will be discussed in the following.
FIG. 6: Density of energy available for particle production as a function of the position in the transverse space for five different intervals of space-time rapidity: $\eta_s = -1$ (upper left panel), $\eta_s = 0$ (upper middle panel), and $\eta_s = 1$ (upper right panel) and for $\eta_s = \mp 2$ (lower panels). Here $\sqrt{s_{NN}} = 17.3$ GeV and $b = 6$ fm and the Gaussian transparency function was used.

Now, when calculating eccentricity coefficients, one could replace $x$ by $x - x_{cm}(\eta_s, b)$ in Eq. (2.13). The so-defined eccentricities are shown in Fig. 8 as a function of space-time rapidity. We observe that the even eccentricity coefficients are relatively large, while odd eccentricity coefficients are rather small which is a consequence of the specific, in our opinion realistic, model considered here. We observe large reduction of odd eccentricity coefficients compared to the naive formula (2.13). The modulus of even coefficients grow with rapidity which is a bit different than in other models in the literature. It would be pedagogical to compare eccentricity coefficient as calculated here to other models used in the literature but we do not feel it is our task at present. Most of the models describe the elliptic flow coefficients. The directed flow was much less often studied, see e.g. [7, 8, 19].

The initial eccentricity coefficients $\epsilon_n$’s are transformed to finally experimentally observed flow coefficient $v_n$’s in the preequilibrium [20] and hydrodynamical [22, 23] evolution phases to be compare to experimental data [13, 24]. This part is to complicated to be discussed here where we wish to concentrate exclusively on the initial conditions consistent with local energy-momentum conservation and baryon number conservation. Relatively simple relation was found for $v_2 \approx 0.25\epsilon_2$ at midrapidity [16] in some simplified version of hydrodynamical evolution. Similar studies could be done for different eccentricity coefficients from our model in a broad range of rapidity of pions using a version of hydrodynamical evolution dedicated to the “low”-energy scattering where plasma is not
FIG. 7: $x_{\text{shift}}$ defined in (3.1) for $\sqrt{s_{NN}} = 17.3 \text{ GeV}$ as a function of $\eta_s$ for five different impact parameters $b = 2$ (solid), 4 (dashed), 6 (dotted), 8 (dash-dotted) for Gaussian transparency functions with $\sigma_\eta = 1$. flavour symmetric and $\mu_B \neq 0$. Not all available codes are able to correctly handle such situations, especially those for our specific model.
FIG. 8: Initial eccentricity parameters as a function of space-time rapidity calculated from “corrected” formulae for $\sqrt{s_{NN}} = 17.3$ GeV for different values of impact parameter $b = 2$ (solid), 4 (dashed), 6 (dotted), and 8 (dash-dotted) fm.

IV. CONCLUSIONS

We have studied consequences of local transverse position (x,y)-dependent energy-momentum conservation for initial eccentricities of the plasma created in the noncentral collisions of symmetric nuclei (Au+Au or Pb+Pb). These studies are of interest for collisions performed by the NA51 experiment at SPS and low-energy scan project at the RHIC collider.

The following general picture of transformation of initial kinetic energy to heat stands behind our model. The first stage of the thermalization happens in the longitudinal direction which is related to a fast partial stopping of two pieces of relativistic nuclear matter. Only then the transverse expansion may start to develop, after some dissipation of initial kinetic energy. We wish to point out here that our dynamical fire-streak model formulated in [5] naturally explains torqued initial conditions [21] that lead to forward-backward azimuthal correlations.

The proposed approach leads to initial eccentricities of the plasma. The eccentricity coefficients strongly depend on space-time rapidity of the plasma especially that for $\epsilon_1$
or $\epsilon_3$. The initial eccentricities are potentially important ingredients for final particle flow generated at the freeze-out space-time moment. Here we have concentrated on initial conditions and do not consider hydrodynamical evolution with such conditions. The specialized groups having codes dedicated to adequate hydrodynamical evolution are welcome to continue the studies.

We have shown that naive calculation of eccentricities with our initial conditions leads to no physical results of eccentricity parameters as a function of $\eta_s$. We have proposed a method how to improve the calculation the eccentricity coefficients in the case of $(x, y)$-dependent longitudinal flow being a consequence of initial collision geometry.

We have calculated $\epsilon_{1.2.3,4}$ as a function of space-time rapidity for different initial impact parameters. The results strongly depend on the impact parameter as in the earlier studies with different initial conditions. In our model of initial longitudinal dynamics the odd eccentricity coefficients strongly depend on the space-time rapidity. We wish to note large $\epsilon_2$ and $\epsilon_4$ everywhere and nonnegligible $\epsilon_1$ and $\epsilon_3$ for large space-time rapidities. The larger the space-time rapidity the larger the $\epsilon_1$ eccentricity parameter. The even initial eccentricity coefficients, ($\epsilon_2$, $\epsilon_4$), are large, different than in other models of initial conditions. They are rather weakly dependent on space-time rapidity. Our initial conditions generate also triangular eccentricities which are usually attributed to event-by-event fluctuations and not to initial geometry. The large eccentricity coefficients in our model are a direct consequence of the energy-momentum conservation broken in many other approaches in the literature.

We have analysed the initial shapes of the matter created right after the (peripheral) collisions. We have shown that the shapes in $(x, y)$ for more peripheral collisions are rather exotic which is reflected in large eccentricity parameters. In our opinion those shapes must evolve very fast already in the preequilibrium phase. The exotic shapes with large gradient of velocity in $z$-direction cause that the fluctuations in the preequilibrium stage are probably crucial but very difficult to put into a mathematical formulation.

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