Constraining the initial stages of ultrarelativistic nuclear collisions

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It is frequently supposed that quark-gluon plasma created in heavy-ion collisions undergoes free streaming at early times. We examine this issue based on the assumption that a universal attractor dominates the dynamics already at the earliest stages, which offers a way to connect the initial state with the start of the hydrodynamic expansion in an approximate but conceptually transparent fashion. We demonstrate that the centrality dependence of the measured particle multiplicities can be used to quantitatively constrain the pressure anisotropy and find that it strongly depends on the model of the initial energy deposition. As an illustration, we compare three initial state models and show that they predict rather different early-time values of the pressure anisotropy. This suggests that while assuming free streaming prior to hydrodynamization may be compatible with some initial state models, in general, features of the prehydrodynamic flow need to be matched with the model of the initial state.

I. INTRODUCTION

Experimental studies of the dynamics of quark-gluon plasma (QGP) created in high energy nuclear collisions aim at understanding initial states over which there is little direct control. The standard approach to modeling heavy-ion collisions involves formation of QGP followed by nonequilibrium evolution until the proper time of about 1 fm/c and subsequently by hydrodynamic evolution until the local temperature drops below the confinement scale, and hadrons are formed. After a stage of hadronic cascade, the final particle distributions are measured.

Early phases of this process are usually assumed to be invariant under longitudinal boosts [1]. This approximation holds best for central events at midrapidity. When this assumption is combined with rotational and translational invariance in the plane transverse to the collision axis, it implies that observables of the system depend only on the proper time elapsed after the collision. Despite their limitations, these approximations open the door to manageable semianalytic considerations at all stages of QGP evolution. The best known consequence of this line of reasoning is the asymptotic late-time behavior of the energy density

$$\mathcal{E} \sim \frac{\Lambda^4}{(\Lambda\tau)^{4/3}}, \qquad (1)$$

where the scale Λ is the only remnant of the initial conditions. Since at such late times the system is already close to local equilibrium, one can translate this into a statement about the

entropy density, which can be directly connected to the multiplicity of observed hadrons [1,2]. The challenge is to relate such physical observables to characteristics of the initial state.

In a series of recent developments it was realized that, at least in the case of boost invariant flow, it may be reasonable to assume that the evolving plasma system behaves in a predictable way already at the earliest stages. In some simple models of equilibration it was observed that certain observables exhibit universal behavior very early on, not only at late times, when the system is very close to equilibrium [3,4]. Such far-from-equilibrium attractors were later identified also in other models of nonequilibrium dynamics, such as strongly-coupled $\mathcal{N} = 4$ supersymmetric Yang-Mills theory through the AdS/CFT correspondence [5,6] and weakly coupled models in kinetic theory [7–15].

It has been argued recently that such early-time far-from-equilibrium attractors exhibit two distinct stages, both occurring at large values of the pressure anisotropy, which precludes a purely hydrodynamic interpretation in the traditional sense. The later stage is determined by the decay of nonhydrodynamic modes, and its details depend on the microscopic theory [16–20]. It is natural to expect that this type of behavior should be generic and independent of any special symmetry assumptions. However, the earliest stage which one observes in boost-invariant models appears to be dominated by the rapid expansion successfully competing with nonhydrodynamic mode decay [12,21,22]. The basic observation underlying our work is that if this is the case, then attractor behavior should also

occur in QCD since it is at least partly a kinematic effect—a consequence of approximate boost invariance. This attractor would then provide a bridge between the initial state and the start of hydrodynamic evolution. This crucial stage, which sets initial conditions for hydrodynamics, is often taken to be free streaming. The question we address is whether this behavior is compatible with any given initial state model.

The existence of an attractor implies a specific power-law dependence of the energy density $\mathcal{E}(\tau)$ for asymptotically small proper time τ. For example, in models based on kinetic theory one finds $\mathcal{E} \sim 1/\tau$, which corresponds to free streaming: the system expands starting from an initial state where the longitudinal pressure vanishes [23]. This behavior has been explored in a number of phenomenological studies [24–27]. However, the existence of a nonequilibrium attractor does not in itself imply free streaming at early times. For instance, in hydrodynamic models of equilibration [3,4], the early time behavior of the system is determined by the transport coefficients and can therefore be tuned to the extent that those parameters can be varied.

At the level of a microscopic theory such as QCD this early time behavior is not known, so we parametrize it as $\mathcal{E} \sim \tau^{-\beta}$ in terms of a constant parameter β. The main technical point, presented in the following section, is that if an attractor exists, then the evolution of the energy density can be determined up to a single integration constant, for any value of the scaling exponent β. The significance of this is that it provides important information about the prehydrodynamic evolution of quark-gluon plasma, as we will discuss below.

The phenomenological analysis of heavy-ion collisions typically involves a model of the initial state, which supplies the initial conditions for hydrodynamic evolution. More specifically, the nuclear thickness functions T_A, T_B of the colliding nuclei are generated [28], and these are combined (using some phenomenological model of the initial state) to produce the energy-density profile at some very early proper-time τ_0. This is then evolved using free streaming, which results in a state of the energy-momentum tensor suitable for initializing hydrodynamics. Our point of departure in the observation is that instead of using free streaming for this prehydrodynamic stage of evolution one should instead try to mimic evolution along the nonequilibrium attractor.

It could be that the relevant attractor is actually free streaming at very early times [29], but a priori, this is not obvious. In this paper we point out that the prehydrodynamic evolution is not independent of the choice made for the model of the initial state. If an attractor governs early-time dynamics, then given an initial state model the centrality dependence of the expected particle multiplicities turns out to be quite sensitive to the parameter β. In this way one can link an initial-state model with the corresponding early-time behavior of the energy density, which is equivalent to identifying the early-time behavior of the boost-invariant attractor.

We find that the compatibility of free streaming with any given initial state model cannot be assumed a priori. In this exploratory study we consider three models of the initial energy deposition involving various scalings. Our goal, however, is not to study their relative merits but to see whether they favor different early-time asymptotics of the energy density and thus different attractors. In order to do this, in each case we calculate the value of the parameter β, which leads to the best fit for the observed multiplicities. While a very good fit can be found for each of these three models, the actual value of β obtained this way varies considerably. This shows that the prehydrodynamic stage of QGP evolution need not be well approximated by free streaming, and the appropriate attractor depends on the chosen initial state model.

II. ATTRACTOR BEHAVIOR

In this section we describe how to exploit attractor behavior of the pressure anisotropy to approximate the full dynamics of boost-invariant systems. We would like to emphasize that the analysis presented here does not assume any specific features of the attractor. In particular, it does *not* assume free streaming at early times.

In the case of Bjorken flow the expectation value of the energy-momentum tensor can be parametrized as

$$\langle T^\mu{}_\nu \rangle = \mathrm{diag}(-\mathcal{E}, \mathcal{P}_{||}, \mathcal{P}_\perp, \mathcal{P}_\perp)^\mu{}_\nu. \qquad (2)$$

We will focus on conformal systems, for which the energy-momentum tensor is traceless, and

$$\mathcal{P}_{||} = \frac{1}{3}\mathcal{E}\left(1 - \frac{2}{3}\mathcal{A}\right), \qquad \mathcal{P}_\perp = \frac{1}{3}\mathcal{E}\left(1 + \frac{1}{3}\mathcal{A}\right), \quad (3)$$

where $\mathcal{E}$ is the energy density, and $\mathcal{A}$ is the *pressure anisotropy*, which is a dimensionless measure of how far the system is from local thermal equilibrium.

The conservation of the energy-momentum tensor can be expressed in the form

$$\tau \partial_\tau \log \mathcal{E} = -\frac{4}{3} + \frac{2}{9}\mathcal{A}. \qquad (4)$$

It will be convenient to define the effective temperature $T \equiv \mathcal{E}^{1/4}$ up to a constant factor, which will play no role in our considerations. After introducing the dimensionless variable $w \equiv \tau T$, the conservation equation (4) is rewritten as

$$\frac{d \log T}{d \log w} = \frac{\mathcal{A} - 6}{\mathcal{A} + 12}. \qquad (5)$$

For a perfect fluid, $\mathcal{A} = 0$ and either Eq. (4) or Eq. (5) suffices to determine the solution, leading to Eq. (1). However, for dissipative systems one must also specify

$\mathcal{A}(w)$, which depends on the microscopic dynamics of the plasma as well as the initial state of the system.

If $\mathcal{A}(w)$ is given, then one can integrate Eq. (5) to solve for the effective temperature as a function of w as follows:

$$T(w) = \Phi_{\mathcal{A}}(w, w_0)T(w_0), \qquad (6)$$

for some initial condition (integration constant) $T(w_0)$. The function $\Phi_{\mathcal{A}}$ reads [29]

$$\Phi_{\mathcal{A}}(w, w_0) = \exp\left(\int_{w_0}^{w} \frac{dx}{x} \frac{\mathcal{A}(x) - 6}{\mathcal{A}(x) + 12}\right). \qquad (7)$$

The subscript $\mathcal{A}$, which appears above, indicates the dependence of this quantity on the pressure anisotropy as a function of w.

Although Eq. (6) expresses the content of the conservation of energy momentum in the Bjorken setting, one still needs to determine $\mathcal{A}(w)$ for a given solution, which will, in general, depend on additional information characterizing the initial state. Crucially, in some model systems, such as hydrodynamic models of equilibration [3,4] or kinetic theory [5,7,10,12–15,21,30,31], there is now a lot of evidence pointing to universal behavior of the pressure anisotropy setting in very early on, when the system is still very far from equilibrium. By this we mean that for a given range of initial conditions, apart from an initial transient, the function $\mathcal{A}(w)$ quickly approaches a universal attractor $\mathcal{A}_\star(w)$, which is determined by the microscopic theory under consideration (perhaps numerically or through some sort of a "slow roll" approximation [3]). We assume that the physically interesting range of initial conditions is in the basin of attraction of this unique attractor. This suggests that it should be a good approximation to replace the form of the pressure anisotropy $\mathcal{A}(w)$, as it appears in Eq. (5), by the attractor $\mathcal{A}_\star(w)$ as follows:

$$T(w) \approx \Phi_{\mathcal{A}_\star}(w, w_0)T(w_0). \qquad (8)$$

Within such an approximation, the temperature at late times is determined by the temperature at early times *alone*: the remaining dependence on the initial state is neglected by assuming that the effective dynamics of the system is captured by its attractor, apart from a negligible initial transient.

The importance of Eq. (8) rests on the fact that it is an explicit relation between the initial and final states of expanding plasma. It implies a relationship between the initial state energy density and the entropy of the near-equilibrium system at late times. One can then estimate the multiplicities of observed hadrons by following essentially the same method as outlined in [29] for the special case of a free-streaming attractor.

III. EARLY TIME BEHAVIOR

As currently understood, the existence of a universal attractor is contingent upon there being a definite, finite, and physically distinguished behavior of the pressure anisotropy at $w = 0$ [3,4]. One can translate this into a statement about the behavior of the temperature at early times. Indeed, under the above assumptions, the conservation of energy-momentum Eq. (4) implies that for asymptotically small proper-time τ

$$\mathcal{E} \sim \frac{\mu^4}{(\mu\tau)^\beta}. \qquad (9)$$

We will focus on $0 \leq \beta < 4$. The scale μ is an integration constant, which reflects the initial conditions, and the exponent β is related to the attractor by

$$\mathcal{A}_\star(0) = 6\left(1 - \frac{3}{4}\beta\right). \qquad (10)$$

While different initial conditions will correspond to different values of the scale μ, the parameter β characterizes the attractor itself and is therefore a feature of the particular microscopic theory under consideration.

For instance, in Müller-Israel-Stewart theory the attractor is the unique stable solution which is regular at $w = 0$, where

$$\mathcal{A}_\star(0) = 6\sqrt{\frac{C_\eta}{C_{\tau\Pi}}} \Leftrightarrow \beta = \frac{4}{3}\left(1 - \sqrt{\frac{C_\eta}{C_{\tau\Pi}}}\right), \qquad (11)$$

and C_η, $C_{\tau\Pi}$ are dimensionless constants given by rescaling the transport coefficients by appropriate powers of T [3,32]. Thus, in such cases the value of β is determined by the transport coefficients.

Aside from simple models, little is known about the existence of attractors or the early time asymptotics of the energy density captured by the parameter β. However, as we will show, one can constrain the value of β by calculating certain hadronic observables sensitive to the initial state and comparing with the experiment.

IV. ENTROPY

The universal attractor determines both early and late-time behaviors of the system, and this fact makes it possible to relate final state entropy to characteristics of the initial state. To streamline the notation we will denote the value of w at very early proper time τ_0 by w_0, and its value at late times τ_∞ by w_∞.

Using Eq. (8), one finds the key relation between the entropy density per unit rapidity at late time and the initial energy density

$$s(\tau_\infty)\tau_\infty = h(\beta)(\mathcal{E}(\tau_0)\tau_0^\beta)^{\frac{2}{4-\beta}}, \qquad (12)$$

where

$$h(\beta) = \frac{4}{3}w_\infty w_0^{\frac{2\beta}{\beta-4}}\Phi_{A_\star}(w_\infty,w_0)^2. \qquad (13)$$

We emphasize that Eq. (12) is a universal, albeit approximate, statement in any model which possesses an attractor in the sense under consideration here.

An important point about the function in Eq. (13) is that it is actually *independent* of the specific values of $w_0 \ll 1$ and $w_\infty \gg 1$ appearing on the right-hand side of this equation. This ensues because the quantity in Eq. (7) diverges for small w_0 and vanishes for large w_∞ precisely in such a way that the dependence on the initial and final values of w drops out in the implied asymptotic limits, leaving a finite and nonzero result. This can be shown, in general, based on the asymptotic behaviors of the pressure anisotropy.

V. CONNECTION TO THE EXPERIMENT

The estimate of entropy density, Eq. (12), can be translated into a statement about centrality dependence of observed particle multiplicities. Given the entropy density in Eq. (12), the charged particle multiplicity of a specific event can be expressed as

$$\frac{dN}{dy} = A\tau_0^{\frac{2\beta}{4-\beta}}h(\beta)\int d^2\mathbf{x}_\perp \mathcal{E}(\tau_0,\mathbf{x}_\perp)^{\frac{2}{4-\beta}}, \qquad (14)$$

where A is a constant, whose value will not be relevant to our considerations. The new element here is allowing for a nontrivial dependence of the initial energy density on the location in the plane transverse to the collision axis. This brings in dependence on the impact parameter of a given event. The underlying assumptions and applicability of this procedure are discussed in Ref. [29]. We will then use formula (14) to estimate the expected multiplicity by averaging over Monte Carlo generated events. For a given event, the calculation of the energy density requires a model of the initial state. In this work we consider three such models, all of which are formulated in terms of a nuclear thickness function $T(\mathbf{x}_\perp)$, which is obtained by Monte Carlo sampling [28].

In the first model, which we will refer to as model I, the energy density for a given event is given by [29,33–37]

$$\mathcal{E}^{(I)}(\tau_0,\mathbf{x}_\perp) = CT^<(\mathbf{x}_\perp)\sqrt{T^>(\mathbf{x}_\perp)}, \qquad (15)$$

where the constant C is independent of the impact parameter $\mathbf{b}$, which enters only through

$$T^<(\mathbf{x}_\perp) = \min(T(\mathbf{x}_\perp+\mathbf{b}/2), T(\mathbf{x}_\perp-\mathbf{b}/2)) \quad (16)$$

with an analogous formula holding for $T^>(\mathbf{x}_\perp)$. The second model is the $p = -1$ case of the reduced thickness event-by-event nuclear topology family of models [38]. In this case,

$$\mathcal{E}^{(II)}(\tau_0,\mathbf{x}_\perp) = C\frac{T(\mathbf{x}_\perp+\mathbf{b}/2)T(\mathbf{x}_\perp-\mathbf{b}/2)}{T(\mathbf{x}_\perp+\mathbf{b}/2)+T(\mathbf{x}_\perp-\mathbf{b}/2)}. \quad (17)$$

The third model is defined by [39,40]

$$\mathcal{E}^{(III)}(\tau_0,\mathbf{x}_\perp) = CT(\mathbf{x}_\perp+\mathbf{b}/2)T(\mathbf{x}_\perp-\mathbf{b}/2). \qquad (18)$$

Here C is again a normalization factor independent of the impact parameter.

We now focus our discussion on the case of PbPb collisions at $\sqrt{s} = 2.76$ TeV at the LHC and confront our findings with the ALICE data, which quotes the multiplicity of charged particles in each of 9 centrality classes [41]. We will use our formula, Eq. (14), to calculate the corresponding prediction in each centrality class. The centrality c is connected to the impact parameter b by the relation $c = \pi b^2/\sigma$, where $\sigma = 797$ fm^2 is the total inelastic nucleus-nucleus cross-section at $\sqrt{s} = 2.76$ TeV.

We define the following ratios of multiplicities at different centralities:

$$Q(c,c') \equiv \frac{\langle dN/dy\rangle_c}{\langle dN/dy\rangle_{c'}}, \qquad (19)$$

where the angle-brackets denote the mean value over events in the specified centrality class. These quantities are independent of the normalization factors C entering Eqs. (15), (17), and (18); they are also independent of the factor $h(\beta)$, which contains the details of the presumptive attractor. However, they retain dependence on the parameter β itself, which is related to the attractor by Eq. (10). In this way, for any value of β, we obtain a set of numbers $Q(c,c')$ which can be directly compared to published experimental results. The best fit for each of the three models is found to be

$$\beta^{(I)} = 1.12, \qquad \beta^{(II)} = 1.96, \qquad \beta^{(III)} = 0.44, \quad (20)$$

with statistical errors not exceeding 0.02. The corresponding longitudinal and transverse pressures are related to β by

$$\mathcal{P}_\parallel = (-1+\beta)\mathcal{E}, \qquad (21)$$

$$\mathcal{P}_\perp = (1-\beta/2)\mathcal{E}. \qquad (22)$$

The beta values for the three models differ by a factor of almost 4.5, which is significant even though we are not attempting to account fully and quantitatively for the errors arising from the various approximations we have made.
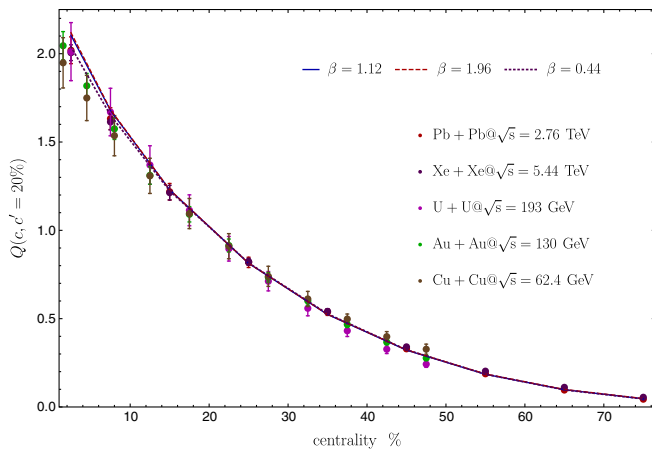
FIG. 1. Universal centrality dependence of $Q(c, c' = 20\%)$, i.e., the number of produced charged particles normalized to 20% centrality for each of the three models we consider. Experimental data shown for different collision systems: Xe + Xe [42], Pb + Pb [41], Au + Au [43], U + U [44], Cu + Cu [43].

This shows that if indeed an attractor determines early time behavior, it is strongly connected to the initial state model.

We can also normalize all the multiplicities to 20% centrality—this choice gives a result closest to what is obtained from the fit to all data points. This allows us to compare the prediction of our analysis to many independent measurements. In Fig. 1 we show our prediction for the ratio $Q(c, c' = 20\%)$ for different collision systems as a function of centrality, normalized to 20% centrality. This is plotted together with data taken from various experiments [41–44].

VI. CONCLUSIONS AND OUTLOOK

An early-time attractor should be viewed as a bridge between an assumed initial state model and the stage of hydrodynamical evolution. By considering three models of the initial energy deposition we have shown that the centrality dependence of measured particle multiplicities can be used to constrain the early-time behavior of the energy-momentum tensor implied by a given model of the initial state.

It is often assumed that irrespective of the initial state model, the prehydrodynamic evolution of QGP can be approximated by free streaming, with the longitudinal pressure vanishing at asymptotically small times. In the three models which we have considered here, this is not the case. Instead, the longitudinal pressure varies significantly between them. Furthermore, none are really close to free streaming: even in the closest case the predicted pressure anisotropy is about 30% lower, relative to what is expected for free streaming.

In conclusion, we have argued that the existence of an early-time attractor couples the preequilibrium evolution of the system to the initial state model. This suggests, in particular, that depending on the model of the initial energy deposition it may not be appropriate to assume a stage of free streaming governing the interval between QGP formation and hydrodynamization. This conclusion was reached using some simplifying assumptions and relies on the most basic observables, but its impact could be even more significant when some symmetry requirements are relaxed, so that more sophisticated observables can be explored.

Our study suggests that it would be worthwhile to look for other options for the prehydrodynamic evolution. This should be particularly relevant for Bayesian studies [27,45–47], which scan over families of initial state models, but assume free-streaming prehydrodynamic evolution for all of them.

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