The nonequilibrium evolution in a boost-invariant Bjorken flow of a hybrid viscous fluid model containing two interacting components with different viscosities, such that they represent strongly and weakly self-coupled sectors, is shown to be characterized by a hydrodynamic attractor which has an early-time behavior that is reminiscent of the so-called bottom-up thermalization scenario in heavy-ion collisions. The hydrodynamization times for the two sectors can differ strongly, with details depending on the curve realized on the two-dimensional attractor surface, which might account for different scenarios for small and large systems in nuclear collisions. The total system behaves like a single viscous fluid with a dynamically determined effective shear viscosity.

INTRODUCTION

Hydrodynamic models have enjoyed an astounding success in describing the collective flow of low-$p_T$ hadrons in heavy-ion collisions [1–4]. These models require hydrodynamics to be initialized at $\lesssim 1 \text{ fm/c}$ after the collisions when the system is far away from equilibrium and with a very low ratio of shear-viscosity to entropy density ($\eta/s \lesssim 0.2$). The studies of heavy-ion collisions in strongly interacting gauge theories using holographic methods [5, 6] were the first to provide theoretical insights not only into the low ratio of shear-viscosity to entropy density [4, 7] but also into the applicability of hydrodynamics itself in far away from equilibrium conditions. In particular, it was understood that the hydrodynamic expansion (which generically has zero radius of convergence) needs to be resummed to all orders in derivatives in order to generate a causal evolution which can be matched with the evolution in the exact microscopic theory for an arbitrary initial condition [8, 9]. Furthermore, the resummed hydrodynamic expansion can be expressed as an appropriate trans-series from which one could also extract the quasi-normal type relaxation modes of the system.

Subsequently, it was established that a wide variety of other phenomenological approaches, such as Müller-Israel-Stewart (MIS) [11, 12] and related versions of extended hydrodynamics [9, 13, 14], and kinetic theory [14–18] also demonstrated the existence of the hydrodynamic attractor, which is the evolution obtained from resumming the hydrodynamic expansion to all orders and to which the system approaches rather quickly for any initial condition. The trans-series has been also examined in such contexts, notably in [9, 13, 19–21]. The ubiquitous presence of the hydrodynamic attractor provides a stronger foundation for hydrodynamics as a causal and consistent effective theory that may be applied generally even when the system is yet to achieve local equilibration. The approach to the hydrodynamic attractor provides a more precise meaning to hydrodynamization [6, 22] of the system, referring to the feature that the energy-momentum tensor and conserved currents can be described by an effective hydrodynamic theory very accurately after a time that is commensurate with a microscopic time-scale for any arbitrary initial data (see [22] for a nice discussion). Strong coupling leads to more rapid hydrodynamization [9] with significant qualitative differences from weak-coupling scenarios [24].

A crucially important feature of quantum chromodynamics (QCD) is asymptotic freedom, which in ultrarelativistic heavy ion collisions makes it conceivable that hard quasiparticle degrees of freedom can be described to some extent by perturbative QCD (in terms of the plasma effective theory [25]) and kinetic theory. However, the gluons produced in their interactions have a coupling that is the stronger the smaller their momenta are, so that strong-coupling methods are needed for the description of the developing bath of soft gluons.

In the so-called bottom-up thermalization scenario initially proposed in [26] and refined in [27], isotropization
and eventually thermalization results from the build-up of a thermal bath of soft gluons. This is in contrast to the completely strong-coupling picture provided by gauge/gravity duality, where thermalization is instead top-down [38]. In [29, 30] it was attempted to find a transition from top-down to bottom-up scenarios as the infinite coupling limit of the standard AdS/CFT correspondence is relaxed through higher-curvature corrections [31]. However, since the quark-gluon medium produced in heavy ion collisions involves more strongly and more weakly coupled sectors simultaneously, a hybrid approach which combines weak and strong coupling features may be required.

To this end, in [32, 33] the semi-holographic approach originally developed by [34, 35] in the context of non-Fermi liquids was utilized for combining the weak-coupling glasma framework [25] with a holographic AdS/CFT description of the infrared sector of soft gluons [37]. In more general terms, a hybrid two-fluid system with couplings analogous to those used in the semi-holographic approach was introduced and studied in [35] in thermodynamic equilibrium.

In this Letter we term the first results of a study of the nonequilibrium evolution of the hybrid fluid model of [35] in a boost-invariant longitudinal expansion known as Bjorken flow [39]. The two components are assumed to have a conformal equation of state, but different amounts of shear viscosity and relaxation times, chosen so that one component corresponds to the infinite coupling limit in AdS/CFT [30] and the other has significant higher shear viscosity corresponding to a less strongly coupled fluid. The two fluids can exchange energy and momentum through mutual deformations of an effective metric that the respective subsystems are living in, while the resulting total energy-momentum tensor of the full system (which actually lives in Minkowski space) is conserved, but with nonvanishing trace caused by the interactions.

The nonequilibrium evolution of the full system turns out to involve a hydrodynamic attractor, which defines a two-dimensional hypersurface in the four-dimensional [41] phase space spanned by the degrees of freedom of the subsystems. Generically, we find that the attractor requires an evolution of the distribution of energy in the subsystems that is reminiscent of the bottom-up thermalization scenario, where the energy is dominated by the more viscous (weakly self-coupled) sector at early times and rapidly shared with the less viscous (strongly coupling) sector.

Subsequently, significant differences in the hydrodynamization times of the subsystems are observed, after which the energy-momentum tensor of the full system can be described hydrodynamically with a specific viscosity determined by which curve on the attractor hypersurface the system follows.

Besides its intrinsic interest as a tractable two-fluid model with hydrodynamic attractor, which to our knowledge is the first of its kind with many possibilities of generalizations, it is of specific interest to quark-gluon plasma physics, where the interplay of hard and soft sectors of the dynamics can be modelled in a new way. In particular, the observed dependence of the ratio of hydrodynamization times on the total energy density might account for different scenarios for small and large systems in nuclear collisions.

**HYBRID FLUID MODEL WITH BJORKEN FLOW**

In [35] a minimal coupling of the energy-momentum tensors of two subsystems was introduced in the following, purely geometrical way, inspired by the semi-holographic setup of [32, 33, 42]. The combined system is defined on one and the same spacetime with metric tensor $g_{\mu \nu}^{(B)}$, but the two energy-momentum tensors of the subsystems, denoted by $\tilde{\mu}^{\mu}$ and $\tilde{\nu}^{\nu}$, are assumed to be covariantly conserved only with respect to effective metric tensors $g_{\mu \nu}$ and $\tilde{g}_{\mu \nu}$, which differ from $g_{\mu \nu}^{(B)}$ and are determined locally by the respective other subsystem according to

$$
\begin{align*}
g_{\mu \nu} &= g_{\mu \nu}^{(B)} + \frac{\sqrt{-g}}{\sqrt{-g^{(B)}}} \left[ \gamma_{\mu \nu} \tilde{g}_{\mu \nu} + \gamma g^{(B)} \delta_{\mu \nu} \right], \\
\tilde{g}_{\mu \nu} &= g_{\mu \nu}^{(B)} + \frac{\sqrt{-g}}{\sqrt{-g^{(B)}}} \left[ \gamma_{\mu \nu} \tilde{g}_{\mu \nu} + \gamma g^{(B)} \delta_{\mu \nu} \right],
\end{align*}
$$

(1)

with two coupling constants $\gamma, \gamma' \equiv -r \gamma$ with mass dimension $-4$. (We need $\gamma > 0$ in order that the dynamics of the subsystems remains causal with respect to the physical background metric, and $r > 1$ for UV completeness [35]). These coupling rules ensure that the full system has a conserved energy-momentum tensor in the physical background, $\nabla^\nu T^\mu_\nu = 0$, with

$$
T^\mu_\nu = \frac{1}{2} \left( (\tilde{\mu}^\mu + \tilde{\nu}^\nu) \frac{\sqrt{-g}}{\sqrt{-g^{(B)}}} + (\tilde{\mu}^\nu + \tilde{\nu}^\mu) \frac{\sqrt{-g}}{\sqrt{-g^{(B)}}} \right) + \Delta K \delta^\mu_\nu =: T_{\nu, \text{int}}^\mu + T_{\nu}^\mu (\mathcal{E}_1, P_1) + T_{\nu}^\mu (\mathcal{E}_2, P_2) + T_{\nu, \text{int}}^\mu 
$$

(2)

where

$$
\Delta K = - \frac{3}{2} \left( (\tilde{\mu}^\rho \tilde{\nu}^\sigma \frac{\sqrt{-g}}{\sqrt{-g^{(B)}}}) \tilde{g}_{\alpha \beta}^{(B)} \frac{\sqrt{-g}}{\sqrt{-g^{(B)}}} \tilde{g}_{\rho}^{(B)} - \frac{3}{2} (\tilde{\mu}^\rho \tilde{\nu}^{(B)} \frac{\sqrt{-g}}{\sqrt{-g^{(B)}}}) \tilde{g}_{\alpha \beta}^{(B)} \frac{\sqrt{-g}}{\sqrt{-g^{(B)}}} \tilde{g}_{\rho}^{(B)} \right).
$$

(3)

In the Bjorken-flow case in flat Minkowski spacetime, with boost-invariant transversely homogeneous expansion in the $z$-direction, the background metric will be given in Milne coordinates $(\tau, x, y, z)$ where $g_{\mu \nu}^{(B)} = \text{diag}(-1, 1, 1, \tau^2)$ with proper time $\tau = \sqrt{t^2 - z^2}$ and spacetime rapidity $\zeta = \text{tanh}^{-1}(z/t)$. In this situation, a boost-invariant ansatz for the effective metric tensors
The dynamical system is closed by assuming relaxation of the energy-momentum tensors of the subsystem reads

\[ \mu^{\mu\nu} = \text{diag}(\frac{\varepsilon}{a^2}, \frac{P_\perp}{b^2}, \frac{P_\parallel}{b^2}, \frac{P_L}{c^2}) + \pi^{\mu\nu}, \]

\[ \pi^{\mu\nu} = \text{diag}(0, \frac{\phi}{2b^2}, \frac{\phi}{2b^2}, -\frac{\phi}{c^2}), \]

and analogously for \( \tilde{\mu}^{\mu\nu} \). Notice, however, that the full system is not conformal, but has a nonvanishing trace of the energy-momentum tensor produced by the interactions of the two subsystems, which involve the dimensionful coupling \( \gamma \).

Covariant conservation with respect to the effective metric (but not with respect to the physical background metric), \( \nabla_\mu \mu^{\mu\nu} = 0 \) yields

\[ \partial_\tau \varepsilon + \varepsilon \partial_\tau \log(b^{8/3}a^{4/3}) + \phi \partial_\tau \log b/c = 0. \]

The dynamical system is closed by assuming relaxation of \( \pi^{\mu\nu} \) by MIS equations with shear viscosities \( \eta, \tilde{\eta} \) and relaxation times \( \tau_\pi, \tilde{\tau}_\pi \) (cf. [40])

\[ \left( \tau_\pi u^\alpha \nabla_\alpha + 1 \right) \pi^{\mu\nu} = -\eta \sigma^{\mu\nu} \]

yielding

\[ \tau_\pi \partial_\tau \phi + \frac{4}{3} \eta \partial_\tau \log b/c + [a + \frac{4}{3} \tau_\pi \partial_\tau \log (b^2c)] \phi = 0, \]

and similarly for the other subsystem. In the following we parametrize \( \eta = C_\eta (\varepsilon + P)/T \) and \( \tau_\pi = C_\tau /T \) with dimensionless constants \( C_\eta \) and \( C_\tau \), where for simplicity we identify \( T = \varepsilon^{1/4} \).

In the following we choose the first system to have ten times higher specific viscosity than the second, \( C_\eta = 10 \tilde{C}_\eta \). Having in mind subsystems that are governed by kinetic theory and AdS/CFT, respectively, we set [4] \( C_\tau = 5 \tilde{C}_\eta \) and \( \tilde{C}_\eta = 1/(4\pi), \tilde{C}_\tau = (2 - \ln 2)/(2\pi) \).

**TWO-FLUID ATTRACTOR**

As found in [9], [15], a single conformal fluid described by the two equations [8] and [10] has a hydrodynamic attractor that can be characterized by the initial condition

\[ \lim_{\tau \to 0} \chi(\tau) = \sqrt{\frac{C_\eta}{C_\tau}} \varepsilon; \quad \chi := \frac{\phi}{\varepsilon + P}, \]

Generic solutions apart from the attractor solution are either singular at finite \( \tau \) or have negative \( \chi \) at early times with limiting value \(-\sigma \) at \( \tau = 0 \). Solutions which are regular for all \( \tau > 0 \) have positive \( \epsilon \) throughout, with \( \epsilon \) diverging as \( \tau \to 0 \).

When the two fluids are coupled according to [1], a common attractor arises where each fluids still has [11] as limiting value, but the behavior of the energy densities is strongly modified. With \( 0 < \sigma < \tilde{\sigma} < 1 \), as is the case for our choice of parameters (\( \sigma \approx 0.45, \tilde{\sigma} \approx 0.62 \)), it turns out that for the common attractor \( \epsilon \) always approaches the finite value \( \sqrt{(\tau - 1)/\tau \epsilon} \) and \( \tilde{\epsilon} \) vanishes as \( \tau \to 0 \). But these quantities are defined with respect to the effective metric tensors \( g \) and \( \tilde{g} \) which become singular as \( \tau \to 0 \). Viewed from the flat Minkowski-Milne background, the contribution to the total energy density from the first system is \( E_1 := (ab^2c/\tau) \epsilon \sim \tau^{4(\sigma - 1)/3} \) and thus diverging for \( \tau \to 0 \) like the single-fluid case. With \( \tilde{\sigma} > \sigma \), the contribution of the second, more viscous fluid is suppressed at early times as \( E_2/E_1 \sim \tau^{8(\tilde{\sigma} - \sigma)/3} \). Depending on the parameters \( \sigma \) and \( \tilde{\sigma} \), \( E_2 \) can diverge or go to zero as \( \tau \to 0 \); for the above choice of parameters it also diverges, but less strongly than \( E_1 \).

This initial distribution of the energy densities is reminiscent of the bottom-up scenario of thermalization in heavy-ion collisions [25, 27], where the evolution starts out with the total energy density concentrated in the more weakly coupled sector, which then gets redistributed to the more strongly coupled soft degrees of freedom.

Away from the limit \( \tau \to 0 \), the differential equations and the nonlinear algebraic coupling equations can only be solved numerically.

In Fig. 1 we display one particular attractor solution and neighboring trajectories in a plot of the anisotropy variables \( \chi \), \( \tilde{\chi} \), and \( \chi_{\text{tot}} \) of the individual subsystems and
The complete set of attractor solutions is in fact a two-dimensional manifold which can be parametrized by the dimensionless energy densities of the subsystems $\gamma \epsilon$ and $\gamma \tilde{\epsilon}$ at some nonzero reference time.

The attractor surface also depends on $r \equiv -\gamma'/\gamma > 1$. When $r$ is sufficiently close to 1, the two-fluid system exhibits a first-order phase transition (with a second-order endpoint at $r = r_c$) as analyzed in detail in [38]. Here we have chosen $r = 2$ so that there is only a cross-over behavior during the evolution of the system as is indeed the case for QCD for high temperature and small quark chemical potential.

In Fig. 2 the particular attractor solution in Fig. 1 is displayed in more detail. In the left panel, the evolution of the energy densities $E_1, E_2, E_{\text{total}}$ are shown for initial conditions which lead to $E_1 = E_2$ at $\gamma^{-1/4} \tau = 1$ and more energy in the less viscous system for a stretch of time before.

Other attractor solutions exist where $E_2$ exceeds $E_1$ for a longer time or where $E_2$ never reaches $E_1$. However, common to all is, as already explained, that at early times the initial energy is concentrated in the more viscous (“hard”) sector which during the Bjorken expansion is gradually transferred to the less viscous (“soft”) sector and to the interaction energy of the two subsystems (the latter is also displayed in Fig. 2 and as a linear plot in Fig. 3). As the system cools down towards its cross-over regime, the interaction energy rises and then switches off rapidly. Subsequently the two subsystems as well as the total system all approach perfect-fluid behavior $\epsilon \sim \tau^{-4/3}$. (Here any similarity to heavy-ion collisions of course ends, which instead ends with free-streaming of hadrons.)

In the right panel of Fig. 2 the evolution of anisotropy in the subsystems and in the full system are displayed.

While the subsystems show a monotonic decrease of anisotropy, this is not the case for the full system, where $\chi_{\text{tot}}$ interpolates the values in the subsystems in a more complicated manner.

From an analysis of the asymptotic behavior, which can be performed analytically and will be discussed in more detail elsewhere, one can show that full system behaves like a single viscous fluid with effective shear viscosity given by

$$C_{\eta}^{\text{eff}} = \lim_{\tau \to \infty} \frac{C_\eta \epsilon^{4/3}(\tau) + \tilde{C}_\eta \tilde{\epsilon}^{4/3}(\tau)}{[\epsilon(\tau) + \tilde{\epsilon}(\tau)]^{4/3}}.$$  \hspace{1cm} (12)

In both plots of Fig. 2 the results are compared with first-order hydrodynamics approximations; in the plots of $\chi$’s, also the second-order approximations are shown. This comparison allows one to define hydrodynamization times. Defining them as the point where the subsequent

FIG. 2. Energy densities and $\chi$’s of one particular attractor curve where the initially subdominant energy density in the less viscous (“soft”) system grows larger than that in the initially dominant more viscous (“hard”) system (blue and red curves, respectively; total system represented by orange curves). The green dash-dotted line gives the interaction energy between the two subsystems. In both plots, the dotted lines correspond to a first-order hydro approximation; in the $\chi$ plots also the second-order hydro approximation is shown (dashed lines).
can be entertained. For example, with smaller parame-
coefficients. Clearly, many generalizations of this model
with different amounts of self-interactions and transport
with MIS equations provides a novel interesting model for
librium properties have been analyzed, in combination
hybrid fluid model introduced in [38], where its equi-
the attractor will be the subject of a longer publication.

As one might expect, the hydrodynamization time is
longer for the hard sector than for the soft, more strongly
interacting sector. This ordering is rather generic [44],
but the ratio depends on the initial conditions.

For example, when these are such that the total energy
in the system at $\gamma^{-1/4} t = 1$ where we have set $E_1 = E_2$ is
increased until we reach the limits of the attractor hyper-
surface, the ratio is reduced to $R_{\text{hd}} \approx 1.4$. Conversely, if
the total energy at this reference point is made as small
as possible, the ratio becomes $R_{\text{hd}} \approx 18$. This could
perhaps be interpreted as a hint that the hydrodynamic
ordering of so-called small systems in nuclear collisions
(high multiplicity events in p-A and p-p collisions) [43]
may involve a markedly different hydrodynamization
scenario than the larger systems produced in heavy-ion

A more complete exploration of the possible evolutions
of the hybrid fluid system presented here as well as a
study of the details how a general solution decays onto
the attractor will be the subject of a longer publication.

To conclude, we believe to have demonstrated that the
hybrid fluid model introduced in [43], where its equi-
librium properties have been analyzed, in combination
with MIS equations provides a novel interesting model for
the nonequilibrium dynamics of a two-component system
with different amounts of self-interactions and transport
coefficients. Clearly, many generalizations of this model
can be entertained. For example, with smaller parame-
ter $r$ one can also include first-order phase transitions,
and the equation of state of the subsystems need not be
conformal.

In the special application to boost-invariant (Bjorken)
flow [41] considered here we have found an intriguing
model of bottom-up thermalization, which has been pro-
posed in the context of perturbative QCD [26, 27]. Re-
markably, this scenario turned out to be a universal fea-
ture of our model, if $\tilde{\sigma} > \sigma$, as is the case for a strongly
coupled sector with transport coefficients as given by
AdS/CFT on the one hand and a kinetic theory with
larger specific shear viscosity on the other hand.

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