

Cosmic Censorship in a Dual Collider

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We investigate cosmic censorship in anti-de Sitter space in holographic models in which the ground state is described by a good singularity. These include supersymmetric truncations of string/M theory, for which a positive-energy theorem holds. At the boundary, our solutions describe a boost-invariant fluid in which the temperature decreases monotonically with time. On the gravity side, they correspond to black-brane spacetimes with a receding horizon. In classical gravity, curvature invariants at the horizon grow without bound. In the full theory this regime may or may not be reached. In some cases it is avoided by a phase transition to a regular geometry. In others it is reached but the boundary hydrodynamic evolution can be continued, provided the equation of state at parametrically small energies is known. Both cases require the inclusion of finite- N or finite-coupling effects.

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Introduction—The formation of spacetime singularities is a prediction of general relativity (GR) and, at the same time, a signal of its own breakdown. Mathematically, the term “singularity” refers to the incompleteness of the evolution. Here we will adopt a physical definition and think of it as a region of arbitrarily large curvature. From this perspective, the content of the weak cosmic censorship conjecture (WCCC) [1] is that, if a singularity forms dynamically from smooth initial data, then it must be hidden behind an event horizon.

Examples of WCCC violation include Choptuik’s critical collapse in four dimensions [2] and black-string pinches in the evolution of the Gregory-Laflamme instability in five dimensions [3–5]. Similar pinches appear in the dynamics of five-dimensional black rings [6] and in black hole collisions in higher dimensions [7–9]. Further plausible violations in four dimensions have also been proposed [10].

A common feature of all these examples is the fact that the singularities are localized in a microscopic region of space and time. As a consequence, GR is still expected to provide a reliable description of these systems, with the possible exception of the few quanta resulting from the evaporation of the singularity [11]. Here we will present simple examples in asymptotically anti-de Sitter (AdS) spaces in different dimensions, including $D = 4$, in which arbitrarily large curvatures in a macroscopic region are generated upon classical time evolution (see Refs. [12,13] for related examples).

The essential ingredients are gravity in AdS_D plus a scalar field whose potential is negative and unbounded from below. This does not necessarily lead to instabilities [14]. For example, a tachyonic scalar field with negative mass squared above the Breitenlohner–Freedman (BF) bound [15,16] is stable. More generally, the existence of a positive-energy theorem is guaranteed provided the potential can be derived from an appropriate superpotential [17–19]. Our examples satisfy this condition. In fact, they include supersymmetric consistent truncations of string/M theory, for which positivity of the energy follows from supersymmetry. The embedding in string/M theory guarantees the existence of a complete dual description in terms of a nongravitational boundary theory in $d = D - 1$ dimensions.

The key idea is to consider models in which the ground-state geometry on the gravity side possesses a “good singularity” [20]. This geometry can be viewed as the

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low-temperature limit of a family of black-brane solutions with regular horizons. Each of these is dual to a static fluid at the boundary in thermal equilibrium. Physically, the singularity does not indicate a pathology of the system but simply that the properties of the ground state, unlike those of thermal states, are not well captured by classical GR.

Consider now the same fluid undergoing boost-invariant expansion, as in the central region of an ultrarelativistic heavy-ion collision [21] (hence our title). In this case, the energy density dilutes and the fluid temperature decreases monotonically with time. Qualitatively, we expect that the gravitational dual solution should correspond to a black-brane spacetime with a receding horizon. We will see that this is true also quantitatively. After a short initial period out of equilibrium, the fluid becomes well described by hydrodynamics. From this point onward, the gravity solution is well approximated by an adiabatic succession of black-brane geometries with a temperature given by the instantaneous fluid temperature. As the latter decreases, the gravity solution approaches the singular ground state geometry and curvature invariants near the horizon grow without bound. Once they reach the Planck or the string scales, the GR description in the bulk breaks down. At this point, the fluid energy density is parametrically small. In some microscopic models, the boundary hydrodynamic evolution can be continued beyond this point provided the equation of state (EoS) in this limit is known. In others, a phase transition to a regular geometry takes place before the high-curvature regime is reached. In both cases, a correct description requires the inclusion of finite- N or finite-coupling effects, corresponding to quantum gravity or stringy corrections in the bulk, respectively.

The gravity dual of a conformal, boost-invariant fluid at late times was constructed in [22]. In this case, conformality implies that the ground state geometry is vacuum AdS and no singularity arises. The construction was generalized to the nonconformal case in [23]. We will see that our solutions approach those in [23] at late times.

Holographic model—We will present solutions to the following theory:

$$S = \frac{2}{\ell_p^3} \int d^5x \sqrt{-G} \left(\frac{R}{4} - \frac{1}{2} \partial_M \phi \partial^M \phi - V(\phi) \right), \quad (1)$$

where G_{MN} is the bulk metric and ℓ_p is the Planck length. We assume that the potential follows from a superpotential $W(\phi)$ through the standard relation

$$V(\phi) = -\frac{4}{3} W(\phi)^2 + \frac{1}{2} \left(\frac{\partial W(\phi)}{\partial \phi} \right)^2, \quad (2)$$

and we take $W(\phi)$ to have the simple form [24]

$$W(\phi) = \frac{1}{4\gamma^2 L} \left[1 - 6\gamma^2 - \cosh(2\gamma\phi) \right]. \quad (3)$$

Here γ is a numerical constant and L is the asymptotic AdS radius. The scalar field has mass $m^2 L^2 = -3$ around the maximum [25] of V at $\phi = 0$. This means that the dual four-dimensional theory is a conformal field theory (CFT) deformed by a dimension $\Delta = 3$ scalar operator. The coefficient of this operator introduces an intrinsic mass scale, Λ , which breaks conformal invariance. The properties of the ground state geometry are determined by the leading exponential in the potential, which for $\phi \rightarrow \infty$ behaves as [provided $\gamma < 2\gamma_c$ (see next section)]

$$V(\phi) \propto -e^{4\gamma\phi}. \quad (4)$$

This model captures the essential physics we wish to describe. We focus on $D = 5$ for concreteness, but analogous results apply in other dimensions, including $D = 4$. We choose the exponential behavior (4) because it appears in supersymmetric truncations of string/M theory to $D = 4$ [27–36] and $D = 5$ [37–47], but we expect similar results for other types of unbounded potentials. In truncations with a single scalar, the latter can be directly identified with ϕ in (2). Our model reduces to [39] for $\gamma = 1/\sqrt{3}$ which, as far as we know, is the only case for which it coincides with a particular string theory truncation. In truncations with multiple scalars, ϕ parametrizes a specific direction in field space [28,29].

Thermodynamics—The EoS of the boundary theory can be reliably determined from the black-brane solutions of (1) provided their curvature obeys

$$\ell_p^2 R \ll 1, \quad \ell_s^2 R \ll 1. \quad (5)$$

The Planck length is related to the number of degrees of freedom in the boundary theory through

$$N^2 = 4\pi^2 L^3 / \ell_p^3, \quad (6)$$

meaning that the classical limit in the bulk, $\ell_p \rightarrow 0$, corresponds to the planar limit in the boundary theory, $N \rightarrow \infty$. The second condition in (5) applies in string theory models. In these cases the string length, ℓ_s , is related to a coupling constant in the boundary theory.

Under assumption (5), the energy density $\mathcal{E}$ as a function of temperature T takes the form shown in Fig. 1. At high temperatures $\mathcal{E} \propto T^4$, as expected for an ultraviolet (UV) fixed point. The behavior at low energy depends on γ . If $\gamma < \gamma_c$, with $\gamma_c = 1/\sqrt{6} \simeq 0.408$ [48], then $\mathcal{E}$ and T approach zero simultaneously. If $\gamma = \gamma_c$ then T approaches a nonzero value as $\mathcal{E} \rightarrow 0$. If $\gamma_c < \gamma \leq 2\gamma_c$ then T diverges as $\mathcal{E} \rightarrow 0$. In all these cases, the black-brane solutions on the gravity side approach the ground-state geometry as $\mathcal{E} \rightarrow 0$. We will focus on $\gamma \leq \gamma_c$ because in this case the approach takes place along a thermodynamically stable branch. If $\gamma > 2\gamma_c$ the potential develops a minimum before blowing up toward $+\infty$. As a consequence (see

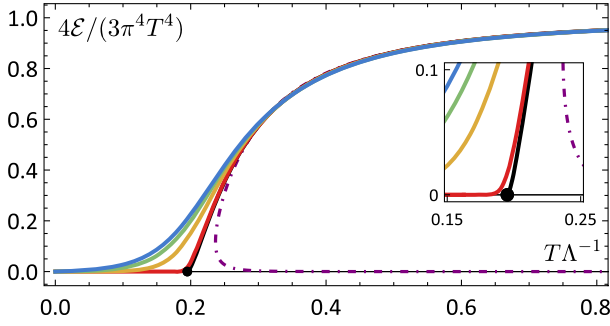


FIG. 1. Energy density as a function of temperature for different choices of $\gamma \leq \gamma_c$, with the same color coding as in Fig. 2. We also show the case $\gamma = 5/10 > \gamma_c$ (purple, dot-dashed curve). The inset plot zooms into the different $\mathcal{E} \rightarrow 0$ behaviors as γ varies from under- to supercritical.

Supplemental Material for details [50]). the theory is expected to be pathological [20].

The EoS, $\mathcal{P}_{\text{eq}}(\mathcal{E})$, is shown in Fig. 2. We see that $\mathcal{P}_{\text{eq}} = \mathcal{E}/3$ at high energies, as expected for a UV fixed point. The low-energy behavior is controlled by the leading exponential (4) and is given by $\mathcal{P}_{\text{eq}} = (1/3 - 2\gamma^2)\mathcal{E}$ [23]. This value, indicated in Fig. 2 by the small ticks on the left vertical axis, shows that the IR physics is not conformal.

The fact that the black-brane solutions approach the singular ground-state geometry as $\mathcal{E} \rightarrow 0$ means that (5) is eventually violated at sufficiently low energy. At this point the EoS above ceases to be valid and a calculation directly in the boundary theory becomes necessary. We will come back to this point below.

Time evolution—We write the boundary metric as

$$ds_{\text{B}}^2 = g_{\mu\nu} dx^\mu dx^\nu = -d\tau^2 + \tau^2 dy^2 + d\vec{x}_\perp^2, \quad (7)$$

where τ is the proper time, y is the rapidity, and $\vec{x}_\perp$ are the two transverse directions. The fact that this is just the flat Minkowski metric written in Milne coordinates means that no source is turned on for the bulk metric, for which we use the characteristic formulation [62–68]

$$ds^2 = -Ad\tau^2 + 2d\tau dr + \Sigma^2(e^{-2B}dy^2 + e^B d\vec{x}_\perp^2). \quad (8)$$

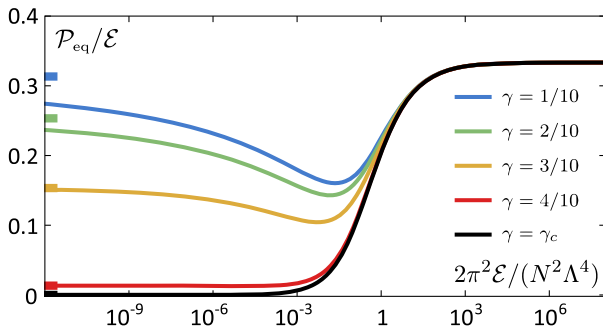


FIG. 2. Equilibrium pressure for different $\gamma \in (0, \gamma_c]$.

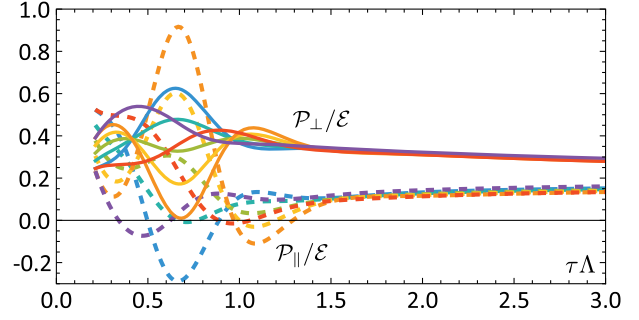


FIG. 3. Short-time evolution of the pressures for different choices of initial profiles for B and ϕ for $\gamma = 4/10$.

The metric functions A, Σ, B and the scalar ϕ depend on τ and on the holographic direction r . After solving the bulk evolution, holographic renormalization [see Refs. [69–71] and Supplemental Material (SM) [50]] can be used to read off the boundary stress tensor

$$\langle T_\nu^\mu \rangle = \text{diag}(-\mathcal{E}(\tau), \mathcal{P}_\parallel(\tau), \mathcal{P}_\perp(\tau), \mathcal{P}_\perp(\tau)), \quad (9)$$

whose conservation implies

$$\tau \mathcal{E}'(\tau) + \mathcal{E}(\tau) + \mathcal{P}_\parallel(\tau) = 0. \quad (10)$$

We set initial conditions at $\tau_i = 0.2 \Lambda^{-1}$ such that $\mathcal{E}_i = 75N^2\Lambda^4/2\pi^2$. We also specify the full radial profiles $B(\tau_i, r)$ and $\phi(\tau_i, r)$ [see Eq. (24) in SM [50]], which fix the initial pressures. In Fig. 3 we see that evolutions with different initial conditions quickly converge to one another. As in previous work [62,67,68], once this happens the evolution becomes well described by viscous hydrodynamics. This is illustrated by Fig. 4, which compares the exact pressures to their hydrodynamic values including first-order viscous corrections [see Eq. (25) in SM [50]]. We see in the inset panel of Fig. 4 that the relative difference between $\mathcal{P}$ and $\mathcal{P}^{\text{hyd}}$ is $\lesssim 10^{-4}$ from a very early time. The subsequent evolution can then be determined by solving (10) with the replacement $\mathcal{P}_\parallel(\tau) \rightarrow \mathcal{P}_\parallel^{\text{hyd}}(\tau)$. The hydrodynamization at late times can also be understood by

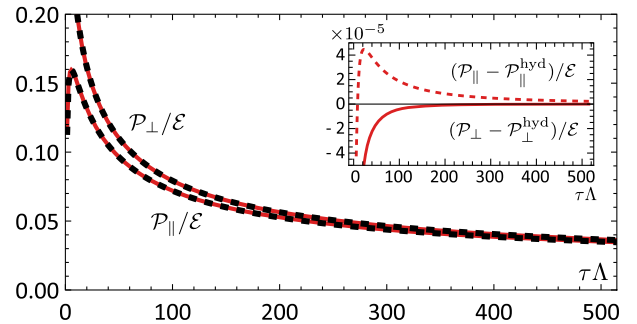


FIG. 4. Comparison between the exact pressures (solid red) and their hydrodynamic values (dashed black) for $\gamma = 4/10$.

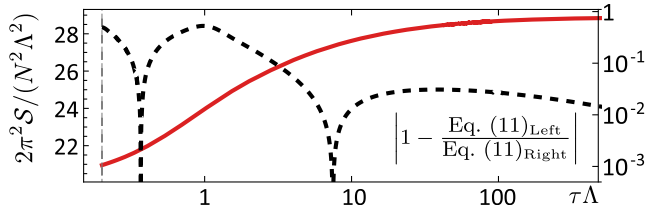


FIG. 5. Solid red: time evolution of the entropy density per unit rapidity and unit transverse area. Dashed black: relative difference between the two sides of Eq. (11).

comparing the expansion rate of the system with the relaxation rate of perturbations—see Fig. S10 in SM [50].

In Fig. 4 we see that $\mathcal{P}_{\parallel}$ and $\mathcal{P}_{\perp}$ differ at early times but converge to one another at late times. At this point, the fluid becomes well described by ideal hydrodynamics and entropy ceases to be produced, as shown in Fig. 5. In other words, the fluid has reached local thermal equilibrium and the subsequent evolution can be viewed as adiabatic. We have extracted the entropy density from the area density of the apparent horizon, but we have verified that the latter coincides with the event horizon at late times, as expected in local equilibrium. A further manifestation of this is the fact that the first law, $dE = TdS - \mathcal{P}dV$, is satisfied. In our case this takes the form

$$\partial_{\tau}(\tau\mathcal{E}(\tau)) = T(\tau)\partial_{\tau}\mathcal{S}(\tau) - \mathcal{P}(\tau), \quad (11)$$

where $\mathcal{S}$ is the entropy density per unit rapidity and unit transverse area, $\mathcal{P} \equiv (\mathcal{P}_{\parallel} + 2\mathcal{P}_{\perp})/3$, and the time evolution of the energy density is shown in Fig. 6. As we see in Fig. 5, the first law is satisfied with increasing precision as time progresses. It is interesting to appreciate the role of each term in (11). The proper volume of a fixed-rapidity region grows linearly with τ . Since the energy density decreases faster than τ^{-1} , the total energy in this region decreases with time, i.e., $dE < 0$, as expected from the dilution caused by the fluid expansion. The total entropy in this region grows monotonically with time and approaches a constant at late times. In this limit $dS \rightarrow 0$ and the first law is satisfied thanks to the $-\mathcal{P}dV$ work term, much like in an expanding FLRW universe.

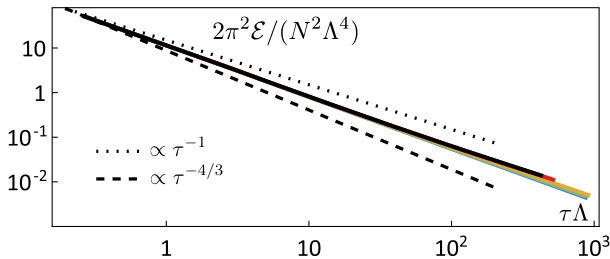


FIG. 6. Time evolution of the energy density for the same values of $\gamma \leq \gamma_c$ as in Fig. 2.

Large curvature—In Fig. 7 we see that the Ricci scalar evaluated at the horizon grows monotonically with time (the same is true for other curvature invariants). As anticipated, the reason is that the expansion of the boundary fluid dilutes the energy density, the bulk horizon recedes deeper into the bulk and the scalar at the horizon “rolls down” the unbounded potential. Our numerical simulations, indicated by the solid color curves in Fig. 7, confirm this picture. For purely numerical reasons, these simulations become challenging at late times. Fortunately, before this happens the fluid at the boundary has hydrodynamized. From that point onward, the bulk geometry can be constructed via the fluid/gravity correspondence [72]. In particular, the curvature at the horizon at a given time τ is well approximated by its value in the static black-brane solution with energy density $\mathcal{E}(\tau)$. The inset panel in Fig. 7 shows that the agreement is better than 1% at the end of the numerical simulations. Using this approximation we can predict the evolution of the curvatures at late times, as indicated by the dashed color curves. In this limit, the behavior is controlled by the leading exponential (4), which implies [23]

$$RL^2 \propto (\tau\Lambda)^{4\gamma^2}. \quad (12)$$

This limit is indicated by gray solid segments in Fig. 7. The agreement with the dashed color curves confirms the accuracy of the fluid/gravity approximation. Note, however, that this assumes that the EoS at parametrically small energy densities can be obtained from classical GR in the bulk. We will come back to this point below.

Sensitivity to initial conditions—In the “Time evolution” section we considered evolutions with different initial conditions that respect boost invariance and translation and rotational invariance in the transverse plane. The fact

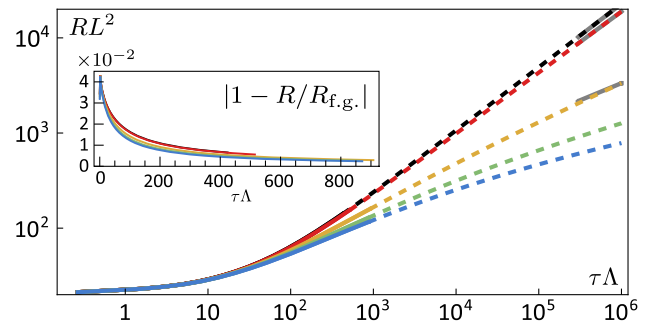


FIG. 7. Ricci scalar R evaluated at the horizon. Solid curves correspond to the numerical evolutions. Dashed curves correspond to the fluid/gravity approximation. The limiting behavior (12) for $\gamma = 3/10, 4/10$ and γ_c is indicated by the gray solid segments on the top-right corner; for the other two curves this behavior is attained at even later times. The inset panel shows the relative difference between the numerical value of the curvature and its fluid/gravity approximation, $R_{f.g.}$, for the times for which they are both available.

that these quickly converge to one another means that no fine-tuning is necessary in this sector. Consider now perturbations that break boost invariance locally but preserve it asymptotically. Hydrodynamic modes of this type are expected to be stable and decay as the fluid expands; see, e.g., [73]. In order to check the stability against nonhydrodynamic perturbations we should study their dynamics on top of the boost-invariant plasma. However, the adiabaticity of the late-time evolution suggests that the answer may be captured by the quasinormal modes of the static black-brane solutions. We have found no evidence of instabilities in these modes (see SM [50]). In summary, our findings indicate that no fine-tuning is necessary in the sector of initial conditions that preserve boost invariance asymptotically. However, this requirement is itself a form of fine-tuning which, in addition, necessitates that the total energy in the system be infinite.

Discussion—We have presented asymptotically AdS solutions in which arbitrarily large curvatures are dynamically generated in macroscopic regions despite the existence of a positive-energy theorem. When condition (5) is violated, the evolution in the bulk can no longer be reliably continued using classical GR. We see from (6) and (12) that this happens at a time $\Lambda\tau_p \sim N^{1/3\gamma^2}$. This is longer than (if $\gamma < \gamma_c$) or comparable to (if $\gamma = \gamma_c$) the time at which the boundary energy density becomes of order N^0 , which scales as $\Lambda\tau_0 \sim N^{3/(2-3\gamma^2)}$ [see Eq. (30) in SM [50]]. Therefore, finite- N effects must be included in order to assess if the high-curvature regime is reached. If it is, then the hydrodynamic evolution at the boundary can be continued, provided the equation of state can be determined at arbitrarily low energy densities [74].

In some cases the high-curvature regime is avoided by a phase transition to a regular geometry. This is illustrated by the 11-dimensional M-theory model of [28], whose black brane solutions were constructed in [29]. This system can be consistently reduced and truncated to a four-dimensional model with six scalars. In this effective description, the horizon values of the scalar fields move along a specific direction in the six-scalar manifold as the energy density decreases. Along this direction, the potential behaves as in (4) with $\gamma = \gamma_c$. In the $\mathcal{E} \rightarrow 0$ limit, the black-brane solutions approach a supersymmetric geometry that is singular even in 11 dimensions. However, in this limit there exists another supersymmetric solution that is regular and describes a confining phase [76]. The reduction of this solution to four dimensions is singular. Moreover, heating it up does not lead to the appearance of a classical horizon, since its free energy scales as N^0 . For these reasons, the confining 11-dimensional solution is not easily discovered in the four dimensional effective theory. Nevertheless, at times $\tau > \tau_0$ we expect a phase transition, akin to hadronization in Quantum Chromodynamics, from the black-brane solutions to the confining branch. Since the confining

solution is regular, this transition cuts off the curvature growth in 11 dimensions (but not in the effective four-dimensional description).

Our results may be viewed as a rather mild violation of WCCC, because the time it takes for the curvature to become Planckian, τ_p , diverges in the classical limit, $N \rightarrow \infty$. It would be interesting to investigate whether instabilities can accelerate the approach to the singularity, along the lines of [77–79].

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Data availability—The data that support the findings of this Letter are openly available [80].

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