

Phase Transitions in an Expanding Medium: Hot Remnants

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We analyze the dynamics of a first order confinement-deconfinement phase transition in an expanding medium using an effective boundary description fitted to the holographic Witten model. We observe and analyze hot plasma remnants, which do not cool down or nucleate bubbles despite the expansion of the system. The appearance of the hot remnants, the dynamics of their shrinking, and subsequent dissolution and further heating up is very robust and persists in such diverse scenarios as boost-invariant expansion with a flat Minkowski metric and cosmological expansion in a Friedmann-Robertson-Walker spacetime.

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Introduction—In this Letter we study the behavior of a physical system that expands and cools, and consequently undergoes a first order phase transition. Studying an expanding medium is interesting in view of various important applications, in particular, an expanding and cooling fireball of quark-gluon plasma in high energy collisions, as well as the expanding of the early universe. In the present Letter we primarily consider boost-invariant expansion providing a simple kinematical implementation of the physics of expanding plasma, but we also discuss uniform cosmological expansion.

We employ a relatively simple model combining hydrodynamics and a dynamical order parameter or “mixing fraction” introduced in [1]. The model was fitted to describe the main features of the holographic Witten model [2] (without flavors), that were worked out in [3], with regard to the physics associated to the phase transitions. However, the structure of the model is very general and we expect that it can be applied in many other settings. Note that the confinement-deconfinement phase transition [4–9] has been studied intensively using holography.

Suppose that a physical system has a first order phase transition at $T = T_c$ and starts off in the stable high energy phase (which is the deconfined phase in the Witten model)

at $T > T_c$. When the system expands, it cools and crosses the phase transition temperature. Then bubbles of the low energy (“confined”) phase start to nucleate. When the temperature falls further below T_c , the bubbles expand, coalesce and the even more supercooled deconfined plasma tends to nucleate even more, which is expected to realize the phase transition.

The main result of this Letter in part challenges this simple picture: we observe the very generic appearance of hot remnants of the high energy phase, whose temperature remains in the vicinity of T_c despite the expansion of the medium (see Fig. 1). For this reason, as the plasma is not supercooled, these remnants do not dissipate through bubble nucleation of the low energy phase, but rather shrink. Then the system starts to “dissolve” with the remaining decreasing fraction of the deconfined phase being counterintuitively even heated up. These findings seem very generic and robust and appear in such diverse scenarios as boost-invariant expansion in flat Minkowski space or cosmological expansion in a FRW universe. The physics behind this behavior is an interplay between the hydrodynamics of the high energy phase, latent heat, and the dynamics of the order parameter, and is thus largely independent from the details of the model. After introducing our model, we will describe the various stages in the evolution of the remnants, namely, the initial “heating up,” a shrinking stage, and subsequent dissolution.

Simplified model of the holographic system—As the bulk numerical relativity description of time-dependent dynamics in the holographic Witten model is extremely challenging [10], in [1], we introduced a boundary model which describes very well the domain wall in the dual field theory, obtained through numerical relativity in holography [3]. The model combines together the hydrodynamic description of the

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deconfined phase with a simple description of the confined phase, by adding an additional dynamical degree of freedom γ describing the mixing of the two phases and thus also the shape of the domain wall in holography (see also [11,12]). Indeed, the derivative terms in the action for γ are necessary for reproducing the profile and finite width of the holographic domain wall of [3], as seen in the components of the energy-momentum tensor extracted from holography. Covariantizing the derivative terms in a way consistent with Lorentz symmetry leads to a kinetic term for γ , making it a dynamical field.

Specifically, the total energy-momentum tensor of the theory is a mixture of the hydrodynamic tensor of the deconfined phase and the trivial $T^{\mu\nu} = \eta^{\mu\nu}$ of the confined phase, together with a contribution describing the domain wall ($T_{\Sigma}^{\mu\nu}$):

$$T^{\mu\nu} = (1 - \Gamma)((\varepsilon + p)u^{\mu}u^{\nu} + p\eta^{\mu\nu}) + \Gamma\eta^{\mu\nu} + T_{\Sigma}^{\mu\nu}, \quad (1)$$

where $\Gamma(\gamma) = \gamma^2(3 - 2\gamma)$ is the mixing fraction (see Supplemental Material [13] for full formulas and rationale for this choice of $\Gamma(\gamma)$ [13]). Here we use units where $p(T) = T^4$, $\varepsilon(T) = 3 T^4$ and the phase transition occurs at $T_c = 1$. The dynamics of γ follows from the Lagrangian

$$\mathcal{L} = \frac{c}{2}(\partial\gamma)^2 + V_{\text{TOT}}(\gamma, T), \quad (2)$$

where

$$V_{\text{TOT}}(\gamma, T) = -(1 - \Gamma)p(T) - \Gamma + \frac{c}{2}q_*^2 T^{\beta}(\gamma)\gamma^2(1 - \gamma)^2 + 1. \quad (3)$$

Here c is related to the surface tension of the domain wall at $T = T_c$ and q_* sets its width. The potential for $T < T_c$, for $T = T_c$ and for $T > T_c$ is shown in Fig. 2. In the case of the Witten model with three-dimensional boundary theory $\beta(\gamma) = 2(1 + \Gamma)$, while for the four-dimensional (“conformal” version) $\beta(\gamma) = 4$ [1]. Note that the above form of the double well potential was fitted only in the neighborhood of $T_c = 1$ up to its first T derivative. For definiteness, we use the above form also away from T_c . The potential $V_{\text{TOT}}(\gamma, T)$ has the interpretation of the Landau free energy of a system with an order parameter γ .

One can view the model (1)–(3) as an extension of the (perfect fluid) hydrodynamics of the deconfined phase (at $\Gamma = \gamma = 0$) by a dynamical order parameter field γ describing the passage to the confined phase (at $\Gamma = \gamma = 1$). The explicit T dependence in the Lagrangian (2) induces a very specific coupling between γ and hydrodynamics through the formalism of [14]. This coupling is necessary in order to well describe [1] the domain wall in the Witten model [3]. Note that we do not have any dissipation terms which are usually added (see, e.g., [15,16]) in order to generate friction to determine bubble wall velocity, as our model

by itself reproduces correct bubble wall velocities in holographic theories [17].

Boost invariant expansion—Boost invariant expansion, first introduced by Bjorken [18], can be seen as a way of modeling extremely energetic heavy-ion collisions, where the system undergoes expansion along the longitudinal (collision) axis. Consequently, this leads to decreasing energy density and hence also decreasing temperature. The inherent time dependence and simplicity led boost-invariant expansion to be a fruitful arena for investigating various dynamical phenomena in holography (see, e.g., [19–24]). The expansion automatically drives the cooling system through any phase transitions or crossover to the low energy confined phase. Technically, one writes the standard Minkowski metric in adapted coordinates

$$ds^2 = -d\tau^2 + \tau^2 dy^2 + dx^2, \quad (4)$$

with y being the (pseudo-)rapidity along the longitudinal axis and τ is the proper time in the longitudinal plane. The assumption of boost-invariance amounts to the independence of the dynamics on y .

Stages in the evolution of the remnant—As initial conditions for our evolution we take constant temperature $T = 0.95$ and we perturb the $\Gamma = \gamma = 0$ plasma phase by two Gaussians. These mimic two seeds which might arise from thermal or quantum fluctuations which are not encompassed by our dynamical model. In Fig. 1, we show the subsequent evolution, with increasing domains of the stable low energy phase appearing ($\Gamma = \gamma = 1$). In between these domains, we have a trapped region of supercooled plasma. The boundaries of this region are domain walls which move towards each other. As demonstrated in [17], the moving domain walls are preceded by hydrodynamic waves at $T \sim T_c$. The decrease of the temperature to the supercooled value in the middle (which also decreases with τ) occurs on the wave fronts. The energy that results in heating up of the plasma comes from the latent heat

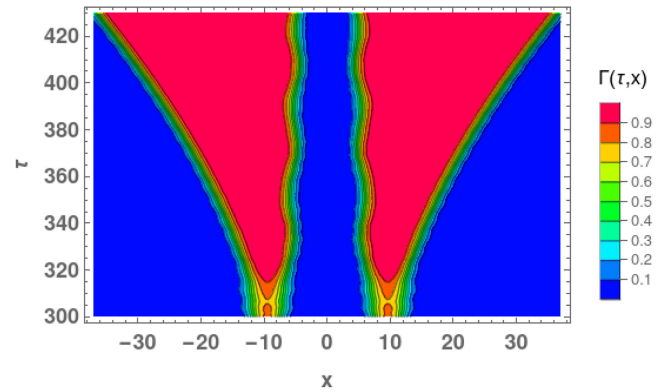
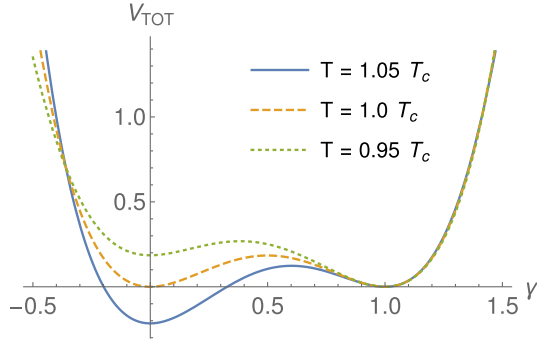
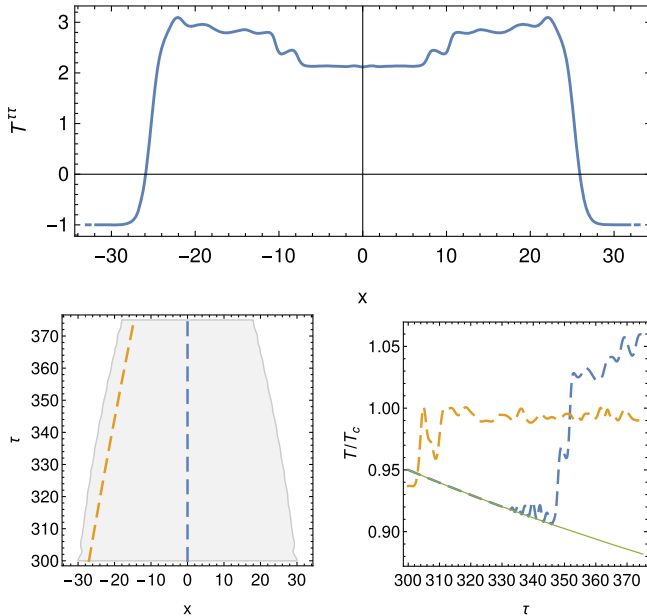


FIG. 1. Evolution of the confined ($\Gamma = 1$) and deconfined phase ($\Gamma = 0$) from two separated Gaussian perturbations of a boost-invariant expanding deconfined plasma system.


 FIG. 2. The potential at $T < T_c$, at $T = T_c$ and $T > T_c$.

released by the phase transition. The profiles of the energy density and the temporal dependence of the temperature in the middle (blue dashed curves) and on the waves (orange dashed curves) are shown in Fig. 3. Once the waves reach the center, the whole plasma region stops being supercooled. The duration of this process can be estimated as $\tau_1 - \tau_0 \sim L_0/(2c_s)$, where c_s is the speed of sound and L_0 is the initial width. The approximate size of the region after this time is $L_1 \sim L_0 - 2\bar{v}_d(\tau_1 - \tau_0)$, where $\bar{v}_d$ is the average domain wall velocity which can be estimated as in [17] (see also [25,26]).

After some oscillations, the subsequent evolution of the plasma remnant becomes quite simple. We observe that the temperature is approximately T_c and stays *constant*, while the fluid velocity is quite small. We should understand how to reconcile this constant temperature with the boost-invariant expansion [along the longitudinal coordinate y


 FIG. 3. The central plasma region of Fig. 1. Hydrodynamic waves at $T \sim T_c$ at the edges of the central region at $\tau = 335$ as seen through the energy density $T^{\tau\tau}$ (top). Evolution of temperature on the wave and at the center (bottom).

(not shown)] which normally leads to $T \sim 1/\tau^{(1/3)}$. To this end, we can make a rough estimate by considering the conservation of energy with the shrinking of the transverse size of the blob releasing latent heat, which can keep the temperature constant.

In view of a similar computation in the FRW spacetime, it is convenient to use an integrated version of the conservation equation in a general time-dependent geometry [27]

$$\partial_t \left(\int_{\Sigma} d^{D-1}x \sqrt{\tilde{\gamma}} T^{00} \right) + \frac{1}{2} \int_{\Sigma} d^{D-1}x \sqrt{\tilde{\gamma}} T^{ij} \partial_i \tilde{\gamma}_{ij} = 0, \quad (5)$$

where $\tilde{\gamma}_{ij}$ is the induced metric on the spatial hypersurface Σ . Neglecting at 0th order any x dependence within the blob, we obtain

$$\partial_{\tau}(\tau E) + P = 0 \quad (6)$$

with

$$E = 3p(\tau)L(\tau) - [R - L(\tau)], \quad (7)$$

$$P = p(\tau)L(\tau) + [R - L(\tau)] \quad (8)$$

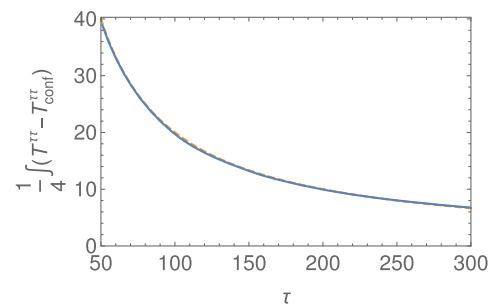
coming from the components $T^{\tau\tau}$ and T^{yy} integrated over x and denoting by R the overall transverse size of the system and by $L(\tau)$ the size of the remnant.

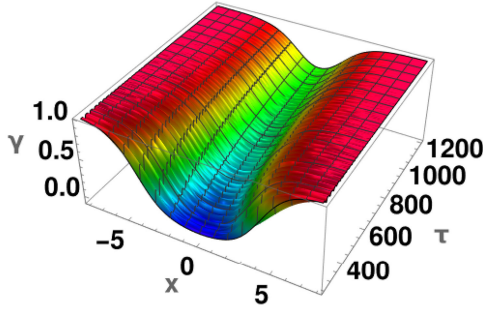
If the remnant is in equilibrium with the confined (i.e., low-energy) phase, then $p(\tau) = 1$ and we see that the size of the remnant scales like

$$L(\tau) \sim \frac{1}{\tau} \quad (9)$$

in agreement with the numerical simulation (see Fig. 4), where we estimate the size by integrating $T^{\tau\tau} - T_{\text{conf}}^{\tau\tau}$ along the spatial direction. The subtraction ensures that only the hot remnant contributes. Incidentally, this behavior at a constant temperature gives an alternative way of satisfying the constancy of entropy per unit rapidity giving

$$S \sim \tau L(\tau) T^3 \propto \text{const} \quad \text{as} \quad L(\tau) \sim \frac{1}{\tau}. \quad (10)$$


 FIG. 4. The evolution of the size of the shrinking remnant estimated through the integral of $T^{\tau\tau}$. The blue curve is numerical data and the dashed orange curve shows the expected $1/\tau$ decay.


 FIG. 5. Evolution of γ when the remnant starts dissolving.

An analogous estimate for Bjorken flow [18] [i.e., standard expansion of plasma $p(\tau) \sim 1/\tau^{(4/3)}$, obtained by setting $L(\tau) = R \equiv L_0$ in Eqs. (7) and (8)] realizes the constant entropy differently

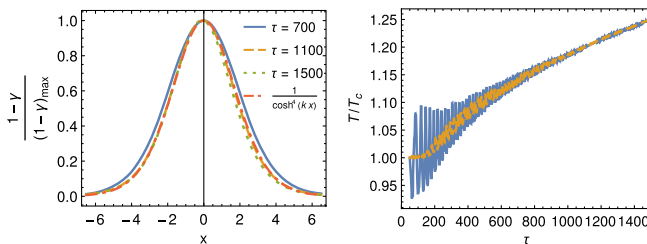
$$S \sim \tau L_0 T(\tau)^3 \propto \text{const} \quad \text{as} \quad T(\tau) \sim 1/\tau^{1/3}. \quad (11)$$

The above rough estimate of the size of the blob assumed a thin wall approximation which means that all sizes are much larger than the width of the domain wall (set in our model by q_*). Therefore, we can estimate the duration of the shrinking stage as $\tau_2/\tau_1 = L_1/L_{\text{dw}}$, where $L_{\text{dw}} \sim 1/q_*$ is the typical width of the domain wall. Once the width of the remnant shrinks to this size, its evolution enters a qualitatively new stage which we now describe.

The value of γ now starts to rise towards $\gamma = 1$, meaning that the fraction of the deconfined high energy phase is decreasing (see Fig. 5). Interestingly enough, the profile of $1 - \gamma$ seems to preserve the same approximate shape [see Fig. 6 (left)]

$$1 - \gamma(\tau) \sim \frac{b(\tau)}{\cosh(d(\tau)q_*\frac{x}{4})}, \quad (12)$$

with $b(\tau)$ behaving roughly like $1/\sqrt{\tau}$ and $d(\tau) \sim 1$ in the range of the numerical simulation. Simultaneously, we observe that the temperature of the residual deconfined plasma slowly *increases* [see Fig. 6 (right)]. We can understand both of these features in the following way.


 FIG. 6. The profile of $1 - \gamma$ normalized to the same height at different times during the dissolving process, together with the approximate profile (12) (left). A slow rise of the temperature (blue) of the dissolving remnant high temperature phase compared with virial theorem prediction (orange) (right).

As the expansion is slow enough, the temporal derivatives in the γ equation of motion are much smaller than the spatial derivatives. Therefore the profile of γ at each instant of time is essentially determined by the virial approximation

$$c\gamma'^2 = 2V_{\text{TOT}}(\gamma, T). \quad (13)$$

Consequently, the maximum of $1 - \gamma$ has to occur at a zero of the potential $V_{\text{TOT}}(\gamma_*, T)$. As for the dissolving remnant $\gamma \rightarrow 1$, this necessarily implies that the temperature must move away from $T = T_c \equiv 1$ and *increase* towards $T > T_c$ (see e.g. the blue curve in Fig. 2 which has a zero at $\gamma_* \sim 0.35$). Indeed, we observe an increase in the temperature in the center of the dissolving blob in accordance to the equation $V_{\text{TOT}}[\gamma(t), T(t)] = 0$ [see Fig. 6 (right)]. It is important to note that T in our model denotes the (hydrodynamic) temperature of the deconfined phase, as the confined phase in the Witten model is temperature independent.

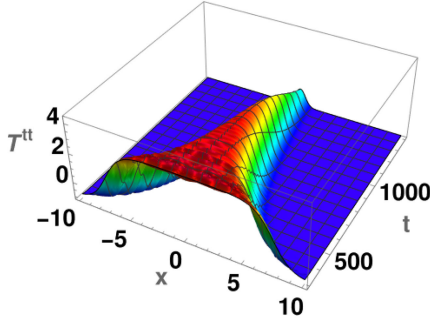
For the description of the profile [Fig. 6 (left)], the relevant part of the potential is the range between $\gamma = 1$ and the first zero $\gamma = \gamma_*$. The approximate shape of the profile (12) follows from the virial equation (13) with a simple approximation for the potential in that range

$$V_{\text{TOT}}(\gamma = 1 - g) \sim d^2 q_*^2 g^2 (1 - \sqrt{g/b}). \quad (14)$$

In the above way we can understand the structure of the temporal snapshots of the dissolving remnant. In order to get a rough estimate of the temporal dependence, we turn to the entropy current $j^\mu = -\partial_\tau V_{\text{TOT}}(\gamma, T)u^\mu$, which is conserved in our model as we do not include dissipation (see Supplemental Material). As in the dissolving remnant configuration the flow velocity is $u^\mu = (1, 0, 0)$, and the entropy conservation reduces to $\partial_\tau j^\tau + (1/\tau)j^\tau = 0$, which gives $s = j^\tau = [f(x)/\tau]$. Using the formula for the entropy current and expanding $s = -\partial_\tau V_{\text{TOT}}$ around $\gamma = 1$ gives $1 - \Gamma \sim 1/\tau$ or, equivalently, $1 - \gamma \sim 1/\sqrt{\tau}$.

Let us comment here on the role played by the dynamical nature of the mixing fraction γ following from (2). The shrinking phase, as estimated by (6)–(8), is essentially independent of the dynamical nature of γ and would be very similar even for infinitely thin domain walls. The kinetic term for γ is important, however, for the dissolving phase, which is happening on the scale of the domain wall width.

Expanding FRW universe—As we saw in the preceding discussion, the appearance and subsequent evolution of the hot remnants could be understood on very general grounds as a consequence of dynamics driven by longitudinal (boost-invariant) expansion. Therefore, it is particularly interesting to consider the system in the context of cosmological expansion, in order to check if similar phenomena also appear despite the differences in the character of the expansion. Moreover, a phase transition in the early universe may be the most physically important application of the phenomena discussed in this Letter.


 FIG. 7. Evolution of T'' in the FRW scenario.

Concretely, we analyze our model in an expanding FRW universe

$$ds^2 = -dt^2 + a^2(t)d\vec{x}^2. \quad (15)$$

For definiteness, we consider here the radiation dominated phase with $a(t) \sim t^{1/2}$, although we checked that similar results hold also, e.g., for $a(t) \sim e^{Ht}$ with small H . Here we work with the physical four-dimensional spacetime, while for numerical reasons all our domains have planar symmetry. Note that the backreaction of the matter system on the cosmological expansion is not taken into account (cf. [28]).

The increasing scale factor $a(t)$ drives the cooling of the system, so a uniform “high energy” phase [with $p(T) = T^4$] in the spacetime (15) expands and cools as $T(t) \sim 1/a(t)$, which keeps the entropy conserved. The plasma becomes supercooled, and nucleates bubbles of confined phase, which we mimic by adding Gaussian seeds and, just as before, leads to the appearance of hot remnants surrounded by the confined phase.

As shown in Fig. 7, we observe the same qualitative stages of the subsequent evolution, first shrinking with the temperature being constant at T_c and then subsequent dissolution. In the FRW spacetime, one can estimate the shrinking of a planar blob with

$$\partial_t(a^3 E) + 3a^2 \dot{a} P = 0 \quad (16)$$

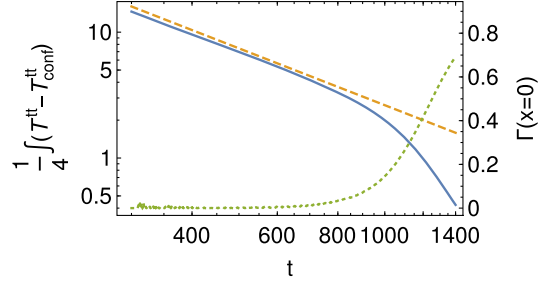
instead of (6). This gives

$$L(t) \sim \frac{1}{a(t)^3}. \quad (17)$$

In case of general, nonplanar blobs we would get $\text{Volume}(t) \sim [1/a(t)^3]$, which again conserves entropy (note that here the volume is measured in the $\vec{x}$ coordinates). In Fig. 8, we see a very good agreement with (17) in the shrinking phase (i.e., when $\Gamma \sim 0$ in the blob).

In order to analyze the dissolution of the remnant, we have to examine the γ equation of motion

$$-\partial_t^2 \gamma - \frac{3\dot{a}}{a} \partial_t \gamma + \frac{1}{a^2} \Delta_x \gamma = \frac{1}{c} \frac{\partial V_{\text{TOT}}(\gamma, T)}{\partial \gamma}. \quad (18)$$


 FIG. 8. Decay of the averaged T'' in the FRW case (blue curve) compared to power-law decay $t^{-(3/2)}$ (dashed orange curve). We also show the value of Γ in the middle of the blob (dotted green curve).

Neglecting the explicit temporal derivatives, the spatial part can be translated into a virial equation, but modulated by the scale factor $a(t)$. This indicates that the width of the dissolving blob shrinks with $a(t)$ in contrast to staying constant in the boost-invariant case. The virial theorem also leads to the conclusion of a rise of the temperature of the dissolving fluid. This is in very good agreement with the numerical simulations as shown in Fig. 9. Note that the observed behavior is driven by the expansion of the universe, and qualitatively different from heated droplets in earlier studies (see Refs. [29,30] and references therein) in particular due to the rise of the temperature above T_c .

Discussion—In this Letter, we explore some generic features of the evolution of holographic Witten model with a confinement-deconfinement phase transition in an expanding setup using an effective QFT boundary description. We find the appearance of hot remnants of the deconfined phase, which get reheated to the phase transition temperature and stay at that temperature while shrinking in size. Subsequently, we observe a distinct stage of dissolution, where the fraction of the deconfined phase decreases, while starting to heat up. These stages and their properties are quite robust, as they appear both in a boost-invariant expansion in flat space and in cosmological expansion in a FRW universe. The observed phenomena do not depend on the details of the employed models despite the

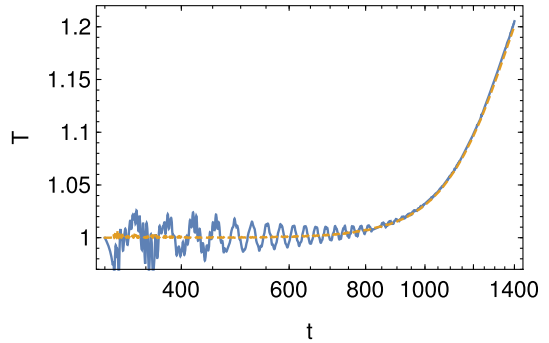


FIG. 9. Evolution of the temperature in the middle of the dissolving blob in the FRW scenario (blue curve) compared to the prediction from the virial theorem (dashed orange curve).

dimensionality and the details of the potential specific for the Witten models being different.

We note that despite the rather complicated and numerically difficult simulations, the appearance and the qualitative features of the hot remnants can be understood quite explicitly using simple physical arguments, which is also a strong indication for their robustness and independence on the details of the models.

We hope that this will facilitate the investigation of their potential role and impact in various applications, in particular, phase transitions in the early universe (see Ref. [31] and references therein), which seem to be the most promising context for the physics described in the present Letter. Indeed, an experimental verification of these phenomena would be of utmost interest.

In this respect, there are a number of physically relevant possible extensions of the present Letter, like going to spherical bubble topology in the FRW case, as well as the case of very fast expansion.

Apart from direct physical applications, this Letter also opens up very intriguing questions on the holographic side. The simplified boundary description for the holographic Witten model used in this Letter, can be understood as a proxy for the (slowly-varying) regime of Einstein's equations in the bulk. From this perspective, the final dissolving phase of the remnant is especially intriguing from the bulk gravitational side, as the dual description entails the disappearance of a black hole and the very final asymptotic stage could be interpreted as a system of glueballs as being a small perturbation on top of the confining geometry. The precise elucidation of what happens on the bulk gravity side remains as a very interesting open problem.

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Data availability—The data that support the findings of this Letter are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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