

Magnetic anti-de Sitter wormholes as seeds for Higgs inflation

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We show how certain types of magnetic asymptotically anti-de Sitter Euclidean wormholes can catalyze the onset of inflation. These wormholes can be embedded as saddle point solutions of general relativity coupled to the Standard Model, the inflaton being identified with the Higgs particle. Our scenario is based on the assumption that the quantum effective potential for the Higgs turns negative at a certain high energy window, in line with current measured values for the Higgs and top quark masses. Within our proposal, we can estimate various parameters and physical quantities of interest, in consistency with current observational bounds.

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I. INTRODUCTION

The tremendous advances of theoretical and experimental physics during the previous century led to an extremely detailed understanding of the history and evolution of our Universe, from the present era back to its primordial origin, when its size was billion times smaller than that of a proton. Cosmic inflation [1–3] is our leading candidate theory describing the physics of the rapidly expanding early universe at this primordial epoch, leading to predictions that conform very well with our current cosmological observational data—a surprising feat considering the vast hierarchy of scales between the primordial and current epoch.

Nevertheless the onset of inflation and the epochs preceding it were never really understood in detail, despite numerous bold efforts of theoretical physicists such as [4,5]. Effective field theory itself offers an operationally concrete setting to pose the relevant questions: Why and how did the universe start with appropriate conditions high up on the inflaton potential? While these are initial/boundary value questions within the context of purely inflationary theory, we do expect that a more complete theory could explain them in a meaningful and natural way, and perhaps in a more dynamical fashion.

In this paper we propose and show how certain types of semiclassical solutions of general relativity coupled to the standard model can catalyze the onset of inflation. These are Z_2 ($\tau \leftrightarrow -\tau$) symmetric “wineglass anti-de Sitter wormholes” [6] that describe the nonperturbative nucleation process of the inflationary universe, whose initial state is determined from the $\tau = 0$ slice of the Euclidean wormhole geometry.

Our paradigm is based on two assumptions: that the inflaton is to be identified with the experimentally observed Higgs particle (the so called “Higgs inflation” scenario [7]), and that the effective Higgs potential turns negative at a certain high energy window [8,9]. While the first assumption is not entirely crucial for the viability of our scenario, it naturally resolves the problem of explaining what happens to the inflaton in our current energy scales and why we have not yet observed any other scalar particle besides the Higgs. The second assumption while being structurally more important for our proposal, it aligns with current experimental data for the Higgs and top quark masses [10] that lead to a metastability of the electroweak vacuum [9].

Our proposed model has various phenomenologically interesting outputs. As other models of Higgs inflation (that essentially generalise Starobinsky’s R^2 model [11]), it can explain very well current cosmological data [12,13]. At the same time it fares better regarding some troublesome points of former Higgs inflation models such as the unitarity issue [14,15]. It also predicts the presence of primordial magnetic fields with a very large coherence length. While these magnetic fields get diluted to an enormous degree due to the accelerated inflationary expansion [16], they could perhaps play an important role even at later stages of the

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evolution of the universe if further physical mechanisms that can amplify and sustain them become active [16–20]. Finally, it has the potential of being embeddable in a UV complete model of quantum gravity such as string theory, due to the Euclidean AdS asymptotics of the relevant wormhole backgrounds [6,21–25].

We use the abbreviations: SM for the standard model, EW for the electroweak scale, EFT for effective field theory, RG for the renormalization group, AdS for anti-de Sitter and natural units $\hbar = c = k_B = \epsilon_0 = 1$ throughout.

II. THE MODEL

Our starting point is the EFT composed of the Einstein-Hilbert action together with a nonminimally coupled scalar field ϕ (that will be identified with the Higgs field) and the rest of the radiation and matter fields of the SM. In Euclidean signature the action is

$$\mathcal{S}_E = \int d^4x \sqrt{g_E} \left[-\frac{1}{2\kappa} R - \frac{\xi\phi^2}{2} R + \frac{1}{2} \partial_\mu \phi \partial^\mu \phi + V(\phi) + \mathcal{L}_{\text{rad}}^E + \mathcal{L}_{\text{matter}}^E \right], \quad (1)$$

where $\kappa \equiv 1/M_{\text{Pl}}^2$ and $V(\phi)$ is the scalar field potential. The nonminimal coupling term typically arises when considering the one-loop effective action of the SM coupled to gravity via quantum loop corrections and is required to renormalize the stress energy tensor of scalar fields in curved backgrounds [26]. The scalar potential should then also be understood to be the loop-corrected effective potential. In general, the EFT contains higher derivative and more complicated nonminimal coupling terms, such as $W(\phi)R^m$. In this work we focus in the *minimal Higgs inflation* scenario where the $\xi\phi^2 R$ term is the most relevant term in such an expansion [7,15]. This is further corroborated by the fact that the characteristic size of the throat of our wormhole backgrounds, is at least $10^5 M_{\text{Pl}}^{-1}$, rendering higher curvature and quantum corrections subdominant [27].

The action (1) can be expressed in the Einstein frame through a Weyl rescaling of the form

$$g_{\mu\nu} \rightarrow \Omega^{-2}(\phi) g_{\mu\nu}, \quad \text{with} \quad \Omega^2(\phi) = 1 + \kappa \xi \phi^2. \quad (2)$$

After performing a field redefinition of the form

$$d\chi = \sqrt{(1 + (1 + 6\xi)\kappa\xi\phi^2)/(1 + \kappa\xi\phi^2)} d\phi, \quad (3)$$

the Einstein frame action reads,

$$\mathcal{S}_E = \int d^4x \sqrt{g_E} \left(-\frac{R}{2\kappa} + \frac{\partial_\mu \chi \partial^\mu \chi}{2} + U(\chi) + \mathcal{L}_{\text{rad,matter}}^E \right) \quad (4)$$

where the potential of the canonical scalar field is given by $U(\chi) = V(\phi(\chi))/\Omega^4(\phi(\chi))$.

Since we want to understand the dynamics of the early universe (having a small size), the leading contribution to the Einstein equations stems from the stress energy tensor of possible axionic fields scaling as a^{-6} , followed by that of the radiation and (nonrelativistic) matter fields, scaling as a^{-4} and a^{-3} , respectively [a is the scale factor in the Friedmann Lemaitre Robertson Walker (FLRW) Ansätze (7)]. The axionic contribution was analyzed in [6], but in what follows, we shall only consider the radiation (density and pressure) components to the stress energy tensor ρ_{rad} , P_{rad} , arising from the experimentally observed SM gauge fields [28]. Another advantage is that their actions are scale-invariant and so their contribution is also invariant under the rescaling (2).

Our approach is phenomenological, as we do not solve the resulting equations of motion (EOMs) for the radiation fields directly, but treat them as a fluid (see [22,29–33] for wormhole examples with radiation fields where the complete field equations are solved). However, this approach suffices to demonstrate the existence of the backgrounds we are interested in and has the merit that it applies to radiation fields in a universal way—all the SM radiation fields contribute to our background solutions. While in their Lorentzian equation of state

$$P_{\text{rad}}^L = \frac{1}{3} \rho_{\text{rad}}^L, \quad \rho_{\text{rad}}^L \geq 0, \quad (5)$$

their energy density is positive definite, the “Euclidean energy” density of the radiation, ρ_{rad}^E , can be negative if the field has a dominant magnetic component. For example, in the $U(1)$ E/M case the “energy” density from the Euclidean stress energy tensor is [34]

$$T_{00}^E = \frac{1}{2} (E^2 - B^2). \quad (6)$$

We therefore find that the magnetic component of the gauge fields gives rise to a negative contribution to the Euclidean “energy,” while the electric component to a positive one (a similar result holds for the axion written as a two-form field). As we demonstrate, *a negative Euclidean energy density is essential to our scenario, implying the necessity of a magnetic radiation component* that was prominent in the early universe.

Arguments regarding the dominance of $SO(4)$ symmetric Euclidean configurations in the path integral [35], consistent with the observed isotropy and homogeneity of our universe, lead us to use the following Euclidean FLRW type of Ansätze for describing its nucleation process

$$ds_{\text{EFLRW}}^2 = d\tau^2 + a^2(\tau) d\Omega_3^2, \quad \chi = \chi(\tau). \quad (7)$$

In this Ansätze, the Euclidean EOMs take the form [37]

$$\begin{aligned}\tilde{\chi}'' + 3\frac{a'\tilde{\chi}'}{a} - \frac{d\tilde{U}(\chi)}{d\tilde{\chi}} &= 0, \\ \frac{2a''}{a} + \frac{a'^2}{a^2} - \frac{1}{a^2} + 3\left(\tilde{U}(\chi) + \frac{\tilde{\chi}'^2}{2}\right) + \frac{\tilde{\rho}_{\text{rad}}^E}{a^4} &= 0, \\ \frac{a'^2}{a^2} - \frac{1}{a^2} + \left(\tilde{U}(\chi) - \frac{\tilde{\chi}'^2}{2}\right) - \frac{\tilde{\rho}_{\text{rad}}^E}{a^4} &= 0,\end{aligned}\quad (8)$$

where the prime denotes a derivative with respect to τ . The conditions that describe a “wineglass” shaped wormhole with a local maximum $a(0) = a_{\text{max}}$ for the scale factor are [6]

$$a''(0) < 0, \quad a'(0) = 0, \quad \chi'(0) = 0. \quad (9)$$

These are favorable initial conditions for a subsequent Lorentzian inflationary evolution after the universe nucleates at $\tau=0$ [6], since upon $\tau=it, d^2a/dt^2(0) > 0$. Additionally, there is another point in the Euclidean evolution where the scale factor must reach a local minimum, a_{min} , i.e., $a'(\tau_{\text{min}}) = 0$ and $a''(\tau_{\text{min}}) > 0$. Since the Euclidean evolution of the scalar happens in the potential $-\tilde{U}(\chi)$ [6], in order to find the wormhole solutions of our interest, we assume an effective potential $\tilde{U}(\chi)$ with a positive global maximum and a negative global minimum, for which (see Fig. 1 for an example)

$$\tilde{U}_{\text{global max.}} > \tilde{U}_0 > \tilde{U}_{\tau_{\text{min}}} > \tilde{U}_{\text{global min.}} \quad (10)$$

From the third equation in (8), we deduce that when $a' = 0$, $x \equiv a^2$ obeys a quadratic equation with solutions

$$x_{\pm} = \frac{-1 \mp \sqrt{1 - 4\mathcal{H}_E\tilde{\rho}_{\text{rad}}^E}}{2\mathcal{H}_E}, \quad \mathcal{H}_E = \frac{\tilde{\chi}'^2}{2} - \tilde{U}(\chi), \quad (11)$$

where $\mathcal{H}_E \geq -\tilde{U}_0$ plays the role of a “Euclidean energy,” that eventually gets dissipated during the motion in the lower Fig. 1. Combining the second and third equations in (8), we determine the acceleration when $a' = 0$

$$\frac{a''}{a} = -3\tilde{U} - \mathcal{H}_E + \frac{1}{x_{\pm}}. \quad (12)$$

A short analysis shows that it is impossible to satisfy the wormhole conditions when $\tilde{\rho}_{\text{rad}}^E > 0$, so a negative energy density is needed. This requirement arises from a Euclidean analog of gravitational focusing where gravity shrinks the space/scale factor, while negative Euclidean energy (magnetic flux in our case) enables a bounce and re-expansion. The competition is between the curvature term $1/a^2$ and the flux term $\sim 1/a^4$, which dominates at small a . Other fields, like axions ($\sim 1/a^6$) or magnetized matter ($\sim 1/a^3$), can play a similar role.

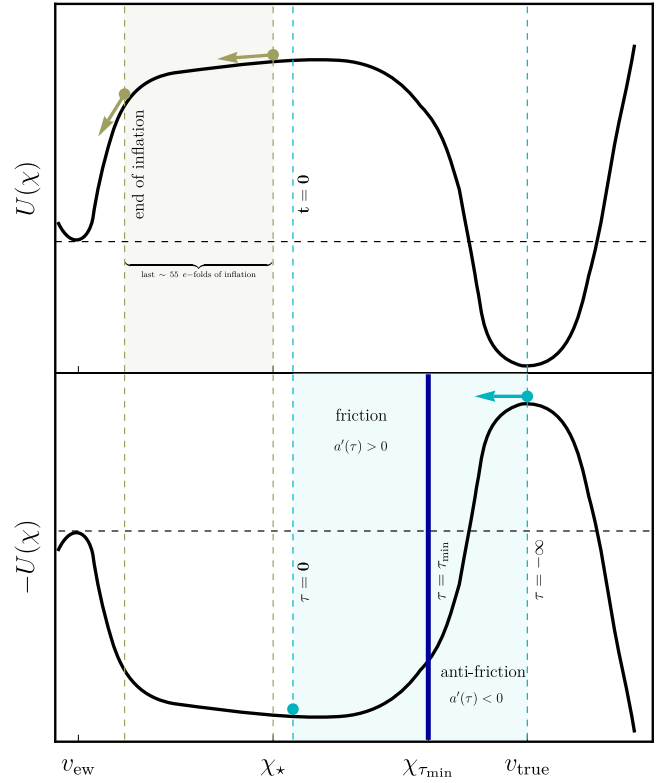


FIG. 1. *Upper:* The Lorentzian evolution. In this case the scalar field naturally reaches the EW vacuum. The condition of having a sufficiently large number of e-folds (wide hilltop region) also helps in the stability of the EW vacuum. *Lower:* The Euclidean evolution corresponding to a half wormhole configuration. It starts with antifriction $a'(\tau) < 0$, since the space is asymptotically Euclidean anti de Sitter. At $\tau = 0 = t$ the universe nucleates and then follows the slow-roll trajectory.

Let us now assume that $\tilde{\rho}_{\text{rad}}^E < 0$, with

$$0 > \tilde{U}_0\tilde{\rho}_{\text{rad}}^E \geq -\frac{1}{4}, \quad (13)$$

so that the scale factor is real. We now find both branches of solutions $x_{\pm}$, $x_{+} \geq x_{-}$, becoming the same at the boundary of the inequality above, see Fig. 2. We therefore choose $a_{\tau_{\text{min}}}$ in the x_{-} and a_{max} in the x_{+} branch. For the a_{max} solution ($\mathcal{H}_E = -\tilde{U}_0$) we find that $a''(0) < 0$ always. For the x_{-} solution, we find that $a''(\tau_{\text{min}}) > 0$ as long as $\mathcal{H}_E^{\text{max}} > \mathcal{H}_E(\tau_{\text{min}}) > -\tilde{U}_0$. The existence of the upper bound is tied to the fact that if $\mathcal{H}_E$ is too high at τ_{min} (for example if one sends $\tilde{U}_{\text{global min.}} \rightarrow -\infty$ in Fig. 1), the friction will not be enough to stop the particle, resulting in overshooting. In the rest we shall assume that the effective potential does not exhibit such pathological characteristics [38].

The result of this analysis is that we can generically obtain the wanted wormhole solutions for a spatially compact early universe, using relativistic fields with a

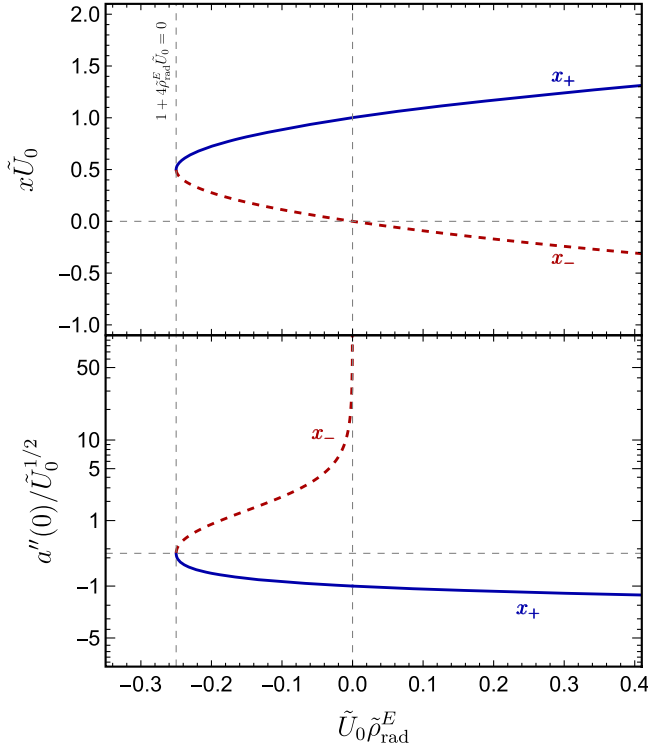


FIG. 2. Normalized scale factor and its second derivative at $\tau = 0$ as functions of $\tilde{U}_0\tilde{\rho}_{\text{rad}}^E$. The conditions $a_{\text{max}} > 0$, and $a''(0) < 0$ are satisfied for $0 > \tilde{U}_0\tilde{\rho}_{\text{rad}}^E > -\frac{1}{4}$ when using the x_+ branch.

slightly dominant magnetic component, as long as the inequality (13) is respected. In this article we do not analyse further possible mechanisms by which such primordial magnetic fields can be produced, see [18,20,39–41] for reviews that focus on their generation during and after inflation.

III. HIGGS INFLATION

As discussed earlier, our model can be embedded in Higgs inflation [7], if we consider a tree level potential $V(\phi) = \lambda(\phi^2 - v_{\text{ew}}^2)^2/4$, where quantum loop effects induce a running quartic coupling, $\lambda(\mu)$ (μ is the RG scale), that remains positive up to the inflationary scale μ_* . The nonminimal coupling, $\xi(\mu)$, has the effect of flattening the effective potential at large field values in the Einstein frame, which can lead to a sufficient number of e -folds for inflation. Assuming values for the Higgs and top quark mass around the central values provided by [10], we seem to be in the metastable region for the EW vacuum. This introduces a negative minimum (v_{true}) in the effective potential, and the quartic coupling turns negative in a certain energy window [42]. We are interested in the parametric regime which places this window above the inflationary scale μ_* , corresponding to earlier times in the evolution of the universe prior to the onset of

inflation [44]. In this scenario, we obtain an inflationary hilltop plateau for $v_{\text{true}} > \chi > v_{\text{ew}}$, and the shape of the quantum corrected effective potential is shown schematically in the upper panel of Fig. 1. While we obtain this scenario for appropriate values of SM parameters, well within experimental bounds (see [43,46,47] for more detailed discussions and a list of references), we remain agnostic regarding the actual value of v_{true} and the detailed shape of the effective potential at such high energy scales, since it is sensitive to the precise UV completion of our model and the scale at which beyond the SM physics, such as supersymmetry, could kick in. Nevertheless in low energy effective actions arising from string theory, one generically finds effective potentials with a global negative minimum, typically preserving some amount of supersymmetry, so our scenario is very natural in potential string motivated UV completions of our model [48].

If we adopt our scenario of Higgs inflation, the potential in the plateau region at horizon crossing is approximately $U(\chi_*) \simeq \lambda_*/(4\kappa^2\xi_*^2)$ and this ratio of the couplings remains approximately constant for most of the inflationary period due to the slow-roll condition. The amplitude of the primordial curvature perturbation on super-Hubble scales generated by single-field slow-roll inflation, $A_s^* = \frac{N_*^2\lambda_*}{72\pi^2\xi_*^2}$, is constrained by recent observations [12] to approximately $\sim 2.1 \times 10^{-9}$ at the pivot scale $k_* = 0.05 \text{ Mpc}^{-1}$. This constraint, in turn, determines the ratio $\lambda_*/\xi_*^2 \sim (4-5) \times 10^{-10}$ and $U(\chi_*) \simeq (1.0-1.2) \times 10^{-10} M_{\text{Pl}}^4$, for $N_* = 60 - 55$ e -folds of inflation.

We should emphasize at this point that these conditions are at the inflationary scale, at lower energies from the end of inflation to the current EW scale $\lambda(\mu)$ has an appreciable nontrivial running, so that [9] $\lambda_{\text{ew}} \simeq 0.126 \gg \lambda_*$ (above the inflationary scale it turns negative). Hence our model accommodates a nonminimal coupling whose value ξ_* at the inflationary scale is much lower than that of the standard tree level discussions of Higgs inflation, where $\lambda_* \simeq \lambda_{\text{ew}}$ [7,50]. This result helps with the issue of violating perturbative unitarity at the scale $\Lambda_u = M_{\text{Pl}}/\xi$, that Higgs inflation models exhibit [14,15,51,52], since for small ξ_* it is pushed well above the inflationary scale, and toward the Planck scale, where we inevitably need to assume a certain UV completion of our EFT model.

Within our model, we can also make an estimate on the value of primordial magnetic fields at the end of inflation and compare it with current experimental bounds. At the time $\tau = t = 0$, where the Lorentzian evolution begins, the field is already positioned in the plateau region, allowing us to approximate

$$U_0 \simeq U(\chi_*) \simeq (1.0-1.2) \times 10^{-10} M_{\text{Pl}}^4. \quad (14)$$

Using Eqs. (11) and (13) we find that the maximum value of the scale factor is approximately $a_{\text{max}}^2 \simeq 3/(kU_0)$.

Substituting this in the definition $\kappa\rho_{\text{rad}}^E/3 = \tilde{\rho}_{\text{rad}}^E/a^4$ and using Eq. (14) we obtain

$$\rho_{\text{rad}\star}^E \simeq \frac{\kappa U_0^2 \tilde{\rho}_{\text{rad}}^E}{3}, \quad 0 > \rho_{\text{rad}\star}^E > -(2.5-3.0) \times 10^{-11} M_{\text{Pl}}^4. \quad (15)$$

We then separate from this density its magnetic component $\rho_{\text{rad}\star}^E = \rho_{\text{rest}\star}^E - B_\star^2/2$. Assuming that no fine-tuning takes place, but rather a small asymmetry between the magnitudes of the magnetic and electric components, the scale of the magnetic field density at the beginning of inflation is $B_\star^2 \simeq 10^{-11} M_{\text{Pl}}^4$. During inflation, the universe was a very poor conductor, leading to an adiabatic dilution of this primordial density by a factor of $\simeq 10^{-96}-10^{-104}$ for $N_\star = 55-60$ e -folds, resulting into $B_{\text{reh}} \simeq 1000-0.1$ G at the beginning of reheating. The rapid expansion also leads to fields with very large coherence length after horizon re-entry. During reheating and up until horizon re-entry the universe turns into a very good conductor. By the time we enter the radiation era, the electric fields get eliminated by the currents of the highly conducting plasma, and the universe is left with classical large-scale magnetic fields [16]. One can then define a ratio $r_B = B^2/8\pi\rho_\gamma$ (ρ_γ is the energy density of relativistic particle species), that remains approximately constant during the history of the universe from reheating until galaxy formation [16]. Possible galactic dynamo mechanisms [40,53–56] that can sustain the currently observed magnetic fields require a minimum “seed” field of the order of $B_{\text{seed}} \simeq 10^{-22}$ G (that is model dependent) and a coherence scale of at least 10 Kpc. Since $\rho_\gamma \simeq 10^{-51}$ GeV⁴ today, we find a sufficient seed ratio $r_B \simeq 10^{-34}$. Evolving back the current radiation fields to the reheating era, we find $\rho_{\gamma,\text{reh}} \simeq 10^{62}$ GeV⁴ assuming instantaneous thermalization. In consistency with [16] we find that the magnetic fields in our model after reheating, while being coherent to extremely large scales, get diluted to such a high degree during the expansion of the universe, making it impossible for them to seed the galactic dynamo mechanism unless additional physical mechanisms are invoked [16–18,20].

IV. DISCUSSION

In this article, we proposed that certain magnetic Euclidean (half) wormhole saddles can catalyze the onset of inflation, while at the same time being embeddable in a model of Higgs inflation containing only experimentally observed SM fields. This constitutes an important development compared to [6], since the models we consider in the present work lead to predictions that are in line with current observational data.

The most interesting future direction is to analyze in great detail the quantum effective potential for the Higgs,

assuming various scenarios for particle physics at the inflationary scale and above (for example the presence or not of supersymmetry and grand unification). Our wormhole backgrounds exist only if the electroweak vacuum is metastable, the energy window where the effective potential turns negative being higher than the inflationary scale. RG corrections could also potentially leave intermediate scale imprints affecting the detailed numbers of observational quantities, such as the spectral index. This allows us to place stringent bounds on particle masses and perhaps other SM parameters, by solving the RG equations for each such scenario and therefore to solidify or rule out our proposal. The presence of primordial magnetic fields in our model also offers a potential window into understanding the origin of the large scale magnetization of our Universe, albeit it suffers from the same inflationary dilution issues that plague similar proposals in the literature.

Finally, due to their anti-de Sitter asymptotics, such Euclidean wormhole saddles [6,21–25] could descend from a fundamental UV complete construction (string theory or the AdS/CFT correspondence). At a more conservative level, these Euclidean saddles are an effective proxy for the expected complicated chaotic dynamics of the early universe, allowing us to perform consistent semiclassical computations, in analogy to studies of quantum mechanical black holes [57].

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DATA AVAILABILITY

No data were created or analyzed in this study.

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