

Exact renormalization group evolution of the inflation dynamics: Reconciling α attractors with ACT

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We present a nonperturbative framework for the dynamics of slow-roll inflation that consistently incorporates quantum corrections, based on an alternative functional renormalization group (RG) approach. We derive the coupled Friedmann-RG flow equations governing the joint evolution of spacetime, the inflaton field, and its effective potential. Applying this formalism to α -attractor E-models, we find that the RG flow induces a dynamical destabilization of the inflationary trajectory, leading to a premature termination of slow roll. Remarkably, the resulting predictions bring α attractors into full agreement with the latest Atacama Cosmology Telescope data without introducing new physics beyond a consistent quantum-corrected treatment of the inflaton dynamics.

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I. INTRODUCTION

Since it was suggested in the 1980s [1–3], the cosmic inflation paradigm, in which a scalar field dubbed the *inflaton* rolled down a relatively flat potential in the early Universe, has passed successfully many of the observational tests (see, e.g., [4] for a review). In particular, single-field slow-roll inflation models explain remarkably well why the primordial spectrum of scalar perturbation observed in the cosmic microwave background (CMB) by *WMAP* [5] and *Planck* [6] is nearly flat. Moreover, searches for tensor modes in the CMB from BICEP/Keck [7] tightly constrained inflationary scenarios, ruling out vanilla models of chaotic inflation, and favoring concave *plateaulike* inflation potentials such as in Starobinsky [2] or, more generally, α -attractor [8–10] scenarios. Recently, high-resolution CMB observations from the Atacama Cosmology Telescope (ACT) [11] and the South Pole Telescope [12], provided the community with a new window to constrain primordial cosmology [13–46]. In particular, data from ACT [47] favor a spectral index $n_s = 0.974 \pm 0.003$, which is on the edge of excluding α -attractor scenarios at 95% of confidence level.

Reconciling α -attractor scenarios with ACT so far required invoking additional new physics leading to a distortion of the α -attractor potential, effectively pushing the spectral index to slightly higher values around the pivot scale [14,28,48–51]. One interesting such possibility is to consider the running of the inflation potential induced by the coupling of the inflaton field to light fermions [52] with a coupling tuned to accommodate the ACT data. In this context, like in numerous works incorporating the effect of quantum corrections on the inflationary trajectory [53–62], renormalization is treated perturbatively, and in a quasistatic manner in which the renormalization scale of the theory evolves very slowly with the slow-rolling field, deep inside the quasi-de Sitter era.

In this paper, instead, we consider the inflation self-interactions only and explore how quantum corrections may affect the inflationary dynamics beyond the leading-order slow-roll approximation. To do so, we use nonperturbative *functional renormalization* techniques based on the Wetterich average effective action [63–65], which allows us to track the evolution of the inflaton potential while the Hubble scale decreases and quantum perturbations exit the horizon and do not contribute to the renormalization flow. In the context of nonperturbative functional renormalization, several works have implemented spatial coarse graining on de Sitter or Friedmann-Lemaître-Robertson-Walker (FLRW) backgrounds, finding scale-dependent effective potentials, interesting nonperturbative infrared phenomena, or explicit dependence on the underlying quantum state [66–70] (see also Ref. [71] for backreaction effects on fixed de Sitter). Complementarily, open-system, in-in, and stochastic approaches have captured how dissipative and noise effects emerge in the long-wavelength inflationary

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dynamics by integrating out short-wavelength fluctuations [72–74].

Here, we take a qualitatively different approach: we use the flow equation derived in Ref. [75], coupling the Hubble-scale-dependent flow of the inflaton potential to the Friedmann dynamics to track the dynamics of the classical background field together with cosmic expansion. As we will see, this evolution can drastically affect the inflationary dynamics toward the end of inflation, considerably changing the field value at which inflation ends, and, thus, the location of the potential where the inflation observables at horizon exit need to be computed.

II. TIME-DEPENDENT COARSE GRAINING IN SPACE

We start by presenting the basic steps leading to our key flow equation for the inflation potential with the Hubble scale. The derivation presented in this section is based on the calculations detailed in Ref. [75].

We consider the FLRW metric $ds^2 = g_{00}dt^2 + a^2(t)(d\vec{x})^2$, and for each time slice, we perform a coarse graining in space by integrating out ultraviolet (UV) Fourier modes of the inflaton. The physical energy scale providing a natural separation between UV and infrared (IR) modes is the Hubble rate H , whose covariant expression with respect to time redefinition is

$$H = \frac{1}{\sqrt{-g_{00}}} \frac{\dot{a}}{a}, \quad (1)$$

which is proportional to the extrinsic curvature for FLRW spacetime. The resulting effective description is provided by the potential U_H , for which we derive a Wilsonian flow equation. The latter must be solved simultaneously with Friedmann equations, which are modified by the H dependence of the potential.

A. Wilsonian flow equation

The coarse graining in space is characterized by a time-dependent scale k , which is eventually related to the Hubble rate. Details of the derivation can be found in Ref. [75], but we present here the main steps that led to our results. The inflaton is decomposed as

$$\Phi = \phi(t) + \varphi(t, \vec{x}), \quad (2)$$

where φ is the quantum field, with three-dimensional Fourier transform defined with the comoving momentum $\vec{p}$:

$$\varphi(t, \vec{p}) = \int d^3x \varphi(t, \vec{x}) e^{i\vec{p}\cdot\vec{x}}. \quad (3)$$

Following Wetterich's approach to Wilsonian renormalization [63–65], but with a coarse graining in space, we

integrate over Fourier modes $\varphi(t, \vec{p})$ weighted by a cutoff function $C_k(p)$ in the path integral, with $p = |\vec{p}|$. The scalar field bare action involved in the path integral is then [75]

$$S_k[\Phi] = \int dt \sqrt{-g} \int d^3x \left(-\frac{g^{\mu\nu}}{2} \partial_\mu \Phi \partial_\nu \Phi - U_{\text{bare}}(\Phi) \right) - \frac{1}{2} \int dt \sqrt{-g} \int \frac{d^3p}{(2\pi)^3} \varphi(t, \vec{p}) \varphi(t, -\vec{p}) C_k(p), \quad (4)$$

where the cutoff function plays the role of a mass squared of the order of k^2/a^2 for IR modes (with $p^2 < k^2$) and is suppressed for UV modes (with $p^2 > k^2$). When defining the one-particle-irreducible (1PI) effective action Γ_k , UV modes are then integrated out without any constraint, while IR modes get frozen, leading to the *average effective action* [63]. By construction, one imposes $\lim_{k \rightarrow 0} C_k(p) = 0$, which implies that, in the deep infrared limit $k \rightarrow 0$, Γ_k becomes the usual 1PI effective action.

Following Ref. [65], we compute this effective action using Eq. (4), but for an arbitrary time-dependent cutoff function $k(t)$. We then perform a functional derivative of this effective action with respect to $k(t)$, and we obtain

$$\frac{\delta \Gamma_k}{\delta k(t)} = i \frac{\sqrt{-g(t)}}{2} \int \frac{d^3p}{(2\pi)^3} \int \frac{d^3q}{(2\pi)^3} \delta(\vec{p} + \vec{q}) \times \int dt' \delta(t - t') \mathcal{O}(t, t', \vec{p}, \vec{q}) \partial_k C_k(p), \quad (5)$$

where

$$\mathcal{O}(t, t', \vec{p}, \vec{q}) = \left[\sqrt{|g|} (C_k(p) - i\varepsilon) \delta(\vec{p} + \vec{q}) \delta(t - t') - \frac{\delta^2 \Gamma_k}{\delta \varphi_b(t, \vec{p}) \delta \varphi_b(t', \vec{q})} \right]^{-1} \quad (6)$$

and φ_b stands for the vacuum expectation value of the quantum fluctuation φ .

We then consider the local potential approximation (LPA), where for every scale k we assume that the effective action for the background ϕ has the form¹

$$\Gamma_k[\phi] = \int dt \sqrt{-g} \int d^3x \left(-\frac{g^{\mu\nu}}{2} \partial_\mu \phi \partial_\nu \phi - U_k(\phi) \right). \quad (7)$$

We use the Litim cutoff function [76], adapted to our situation, which reads

$$C_k(p) = a^{-2}(k^2 - p^2) \Theta(k^2 - p^2). \quad (8)$$

¹We estimate the wave-function renormalization in the Appendix, based on an adiabatic approximation, and we find that it is indeed negligible.

In terms of the physical scale $\kappa \equiv k/a$, we find the Euclidean flow equation

$$\partial_\kappa U_\kappa(\phi) = \frac{aT^{-1}}{6\pi^2} \mathcal{D}_E^{-1}(a^{-1}\kappa^4), \quad (9)$$

where

$$\mathcal{D}_E = -\frac{d^2}{dt^2} - 3H\frac{d}{dt} + \kappa^2 + \partial_\phi^2 U_\kappa(\phi) \quad (10)$$

and T^{-1} is a frequency cutoff, which is unavoidable since the cutoff function C_k acts in three-dimensional Fourier space only.² In Ref. [75], we showed that this cutoff can be fixed by requiring the flat-spacetime limit of our flow equation to reproduce the standard four-dimensional functional-RG flow, leading to $T^{-1} = 3\kappa/16$.

We then divide by κ and multiply by the operator $\mathcal{D}_E$ both sides of Eq. (9), go back to real time, and choose $\kappa = H$, to finally obtain the flow equation

$$\ddot{X} + H\dot{X} + (U_{\phi\phi} - H^2 - \dot{H})X = \frac{H^4}{32\pi^2}, \quad (11)$$

where $X \equiv H^{-1}\partial_H U$ and a dot denotes a total time derivative.

B. Modified Friedmann equations

The action describing gravity and matter is

$$\begin{aligned} \Sigma &= \int d^4x \sqrt{-g} \left(\frac{M_p^2}{2} R - \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - U_H(\phi) \right) \\ &= \int \left\{ \frac{3}{2} M_p^2 \left[-\frac{2}{g_{00}} \left(\frac{\ddot{a}}{a} + \left(\frac{\dot{a}}{a} \right)^2 \right) + \frac{\dot{g}_{00}}{g_{00}^2} \frac{\dot{a}}{a} \right] \right. \\ &\quad \left. - \frac{1}{2g_{00}} (\dot{\phi})^2 - U_H(\phi) \right\} \sqrt{-g_{00}} a^3 d^4x, \end{aligned} \quad (12)$$

and the field equations are obtained from the variation with respect to the dynamical variables g_{00} , a , and ϕ . When taking the variation of the action with respect to g_{00} before setting it to -1 , we should not forget the covariant definition (1) of the Hubble rate which appears in the potential. With the above definition of X , we obtain

(i) $\delta\Sigma/\delta\phi = 0$ for the scalar field evolution

$$\ddot{\phi} + 3H\dot{\phi} + \partial_\phi U = 0; \quad (13)$$

(ii) $\delta\Sigma/\delta g_{00} = 0$ for the Hamiltonian constraint

$$3M_p^2 H^2 = \rho = \frac{1}{2} \dot{\phi}^2 + U - H^2 X, \quad (14)$$

where ρ is the energy density; and

(iii) $\delta\Sigma/\delta a = 0$ for the scale factor dynamical equation

$$\begin{aligned} p &= -M_p^2 (2\dot{H} + 3H^2) \\ &= \frac{1}{2} \dot{\phi}^2 - U + H^2 X + \frac{1}{3} \frac{d(HX)}{dt}, \end{aligned} \quad (15)$$

where p is the pressure.

The last two equations lead to the dynamical equation for H :

$$2M_p^2 \dot{H} = -\dot{\phi}^2 - \frac{1}{3} \frac{d(HX)}{dt}. \quad (16)$$

If one takes into account $\dot{U} = U_\phi \dot{\phi} + HX\dot{H}$, one can check that the two Friedmann equations (14) and (16) imply the scalar field equation (13), as expected, which is consistent with the continuity equation $\dot{\rho} = -3H(\rho + p)$. Trading the cosmic time variable for the number of e -folds $N \equiv \ln(a)$, with primes denoting now derivatives with respect to N , the master set of equations we need to solve is

$$\begin{aligned} X'' &= \frac{H^2}{32\pi^2} - \left(1 + \frac{H'}{H} \right) X' - \left(\frac{\partial_\phi^2 U}{H^2} - \frac{H'}{H} - 1 \right) X, \\ \phi'' &= - \left(3 + \frac{H'}{H} \right) \phi' - \frac{\partial_\phi U}{H^2}, \\ H' &= - \frac{H}{2M_p^2} \left(\frac{(\phi')^2 + X/3}{1 + X/(6M_p^2)} \right), \end{aligned} \quad (17)$$

in which it is assumed that $X > -6M_p^2$, as will be the case throughout this paper.

C. de Sitter solution

In Ref. [75], we showed that this system of equations naturally has a pure de Sitter solution, in which $X' = H' = \phi' = 0$, as long as the field is located at an extremum of the potential (such that $\partial_\phi U = 0$) and

$$X = - \frac{H^2}{32\pi^2 (1 - \partial_\phi^2 U/H^2)}. \quad (18)$$

This case is trivial but representative of what happens physically: to maintain Hubble constant and all the time derivatives in the master equation to remain zero, one needs to adjust the flow of the cosmological constant with the Hubble scale ($X \neq 0$) using Eq. (18), together with the Hamiltonian constraint.

²For usual four-dimensional Wilsonian flows, the cutoff function C_k acts as a full regulator and no extra hard cutoff is needed.

D. From de Sitter to slow-roll inflation

Although this solution is a fixed point of the set of equations describing the dynamics, this equilibrium does not need to be stable. Indeed, we showed in Ref. [75] that such a pure de Sitter solution is a saddle of the dynamics, in the case where the curvature of the potential is smaller than the squared Hubble scale, which typically is the case for slow-roll inflation potentials during inflation. This means that attempting to fine-tune the RG flow at initial time to protect the potential from running with cosmic expansion cannot protect the potential from running throughout the whole duration of inflation.

In what follows, instead, we will consider initial conditions in which the initial condition for the flow is exactly zero, such that the potential does not flow initially, observe that this initial apparent adiabaticity cannot be stable over time, and track where the running of the potential over time is bringing the dynamics.

In this paper, we will restrict ourselves to the case of the E-model class of α attractors [77], although the following analysis could, in principle, be generalized to any kind of single-field slow-roll inflation. The potential we consider is defined, at the Hubble scale H_0 , as

$$U(\phi, H_0) = V(\phi) = V_0 \left(1 - e^{-\sqrt{\frac{2}{3\alpha}} \frac{\phi}{M_P}} \right)^2. \quad (19)$$

This class of potentials is particularly attractive, as it encapsulates many different inflation scenarios, such as Starobinsky inflation ($\alpha = 1$) [2], scenarios from $\mathcal{N} = 4$ supergravity (SUGRA) [78–80] ($\alpha = 1/3$), or so-called Goncharov–Linde (GL) models [81–84] ($\alpha = 1/9$).

III. NUMERICAL ANALYSIS

The set of equations (17) is, in principle, infinite dimensional. Indeed, calculating partial derivatives of the potential $U_H(\phi)$ with respect to ϕ would require solving this system of nonlinear differential equations for multiple trajectories $\phi(N)$ and computing on a lattice partial derivatives at every point. Instead, one can use a decomposition of the potential on a basis that helps tracking the evolution of the potential along the trajectory of interest. Before we present our results, we detail here our full numerical procedure.

A. Exponential basis decomposition

One possibility of such decomposition, given the form of the potential (19), is to use an exponential basis

$$U_H(\phi) = \sum_{n=0}^{\infty} U_n(H) e^{-n\phi/M}, \quad (20)$$

where we defined $M = \sqrt{3\alpha/2} M_P$. The corresponding variable X is then written

$$X = \sum_{n=0}^{\infty} X_n e^{-n\phi/M}, \quad X_n = H^{-1} \frac{dU_n}{dH}, \quad (21)$$

and the total derivative for X becomes

$$X' = \sum_n \left(X'_n - \frac{n\phi'}{M} X_n \right) e^{-n\phi/M}. \quad (22)$$

The exponential expansions above are plugged into the flow equation [first line of Eq. (17)], and the identification of each power of the exponential leads to

$$\begin{aligned} X''_n = & \delta_{n0} \frac{H^2}{32\pi^2} + \left(\frac{2n\phi'}{M} - 1 - \frac{H'}{H} \right) X'_n \\ & + \left[n \frac{\phi''}{M} - \left(n \frac{\phi'}{M} \right)^2 + \left(1 + \frac{H'}{H} \right) \left(1 + \frac{n\phi'}{M} \right) \right] X_n \\ & - \frac{1}{H^2 M^2} \sum_{\ell=0}^n \ell^2 U_\ell X_{n-\ell}. \end{aligned} \quad (23)$$

After truncating this ladder of equations to a sufficiently high order to ensure numerical stability, it is to be solved simultaneously with the Friedmann equations for $\phi(N)$ and $H(N)$ in Eq. (17).

B. Initial conditions

For the potential of Eq. (19) and a given set of potential parameters (V_0, α) , the initial value of the U_n is determined at all orders:

$$U_0 = V_0, \quad U_1 = -2V_0, \quad U_2 = V_0, \quad U_{n>2} = 0. \quad (24)$$

For every retained coefficient, we impose

$$X_n(0) = X'_n(0) = 0.$$

We specify $\phi(0) = \phi_0$ and $\phi'(0) = 0$. The Hamiltonian fixes

$$H_0^2 = \frac{V(\phi_0)}{3M_P^2}, \quad H'(0) = 0.$$

The initial surface is not a stationary solution, since

$$X''(0) = \frac{H_0^2}{32\pi^2}, \quad \phi''(0) = -\frac{V_\phi(\phi_0)}{H_0^2}.$$

Together, these equations guarantee that the Universe starts in an instant de Sitter state defined by $\phi' = H' = 0$ and deviates from its de Sitter state only because of the field quickly reaching its slow-roll attractor and the flow deviating from zero.

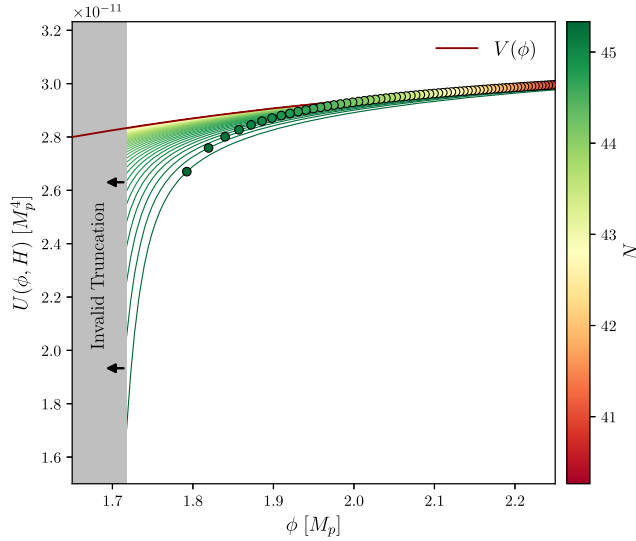


FIG. 1. Full exact renormalization group (ERG) evolution of the potential $U(\phi, H(N))$ (colored lines), for $\alpha = 0.2$, $\phi_0 = 3.175M_p$, and for various values of the number of e -folds since horizon crossing [assumed to be at $\phi(0) = \phi_0$]. For each curve $U(\phi, H(N))$, colored circles depict the corresponding values of the background field $\phi(N)$ at time N .

C. Effect on the inflationary dynamics

A representative example of the inflationary dynamics and its exact renormalization group (ERG) flow evolution is represented in Fig. 1, in the case where $\alpha = 0.2$. In this figure, both the Hubble scale and field value evolve with time (or the number of e -folds) while the scalar potential is driven by the flow equation, according to the set of equations (17). The solution $\phi(N)$ is depicted with colored dots, while colored curves indicate the overall shape of the potential when evaluated at the corresponding values of $H(N)$ for arbitrary ϕ . In that plot, the gray-shaded area stands for the region where the truncation of the sums in the exponential basis decomposition cannot be trusted (which we evaluate numerically, by varying the truncation order and demanding that its relative impact on the potential shape is smaller than 0.1%). As one can see from the figure, while the field rolls down, the potentials drop at larger and larger field values over time, thus destabilizing the potential and prematurely accelerating the field, enforcing slow-roll exit earlier than in the absence of running. Remarkably, this has the important effect of pushing the field value at which the CMB modes exit the horizon further away on the plateau of the potential, modifying drastically the inflation observables, as we will see in the next subsection.

D. Inflation observables

To explore the effect of the inflationary ERG flow on CMB observables, we scan over the parameter space for various values of $\alpha \in [10^{-5}, 1]$ and varied values of ϕ_0 . For each parameter choice, we calculate the slow-roll parameters

$$\epsilon_V \equiv \frac{M_p^2}{2} \left(\frac{\partial_\phi U}{U} \right)^2, \quad \eta_V \equiv M_p^2 \frac{\partial_\phi^2 U}{U}, \quad (25)$$

and normalize the potential at $\phi = \phi_0$ such that the primordial power spectrum of scalar perturbations has the correct amplitude at the *Planck* pivot scale [6]

$$A_s \equiv \frac{U(\phi_0, H_0)}{24\pi^2 M_p^4 \epsilon_V(\phi_0)} \simeq 2.1 \times 10^{-9}. \quad (26)$$

Then we track the dynamics of the inflaton, Hubble scale, flow functions X_n , and potential functions U_n dynamically until the system exits slow roll when $\epsilon_H \equiv -H'/H = 1$. We then calculate $N_\star$, the number of e -folds between horizon exit and the end of inflation, and select parameters for which $50 < N_\star < 60$. The tensor-to-scalar ratio and spectral index of the primordial scalar perturbation power spectrum are then evaluated using their slow-roll expression at horizon exit, $r \simeq 16\epsilon_V(\phi_0)$ and $n_s \simeq 1 - 6\epsilon_V(\phi_0) + 2\eta_V(\phi_0)$.

E. Results

Our results are presented in Fig. 2, in which the left panel stands for the usual case without running and the right panel includes the full ERG flow in the dynamics. In the latter case, we assess truncation convergence by increasing the retained order and comparing the effect of order variation on the potential shape throughout the inflation trajectory. In the example depicted in Fig. 1, this means that the end of inflation takes place before the colored dots enter the gray region. Results are compared to the most recent constraints from *Planck*, BICEP, and Keck [7] and ACT [47]. As expected from the previous qualitative discussion, the effect of the running is to provoke the ending of inflation much earlier compared to the case of a rigid potential. This effectively pushes the value of the field at horizon crossing to larger values where the potential is flatter, effectively increasing the value of the spectral index.

IV. SUMMARY AND DISCUSSION

In this paper, we have explored the evolution of the inflationary dynamics in the presence of an evolving effective potential. Because of the quantum modes exiting the horizon during inflation, and, hence, stopping contributing to the effective potential felt by the classical inflaton background field, this evolution ultimately destabilizes the slow-roll regime, enforcing an early ending of inflation, and effectively pushing the inflaton to be further up its potential when the CMB modes exit the horizon. Using the example of α -attractor E-models, we saw that the effect of this ERG running is to reconcile the model with recent CMB data from ACT that were initially thought to exclude α -attractor models with 95% confidence level. The quantitative predictions, in principle, depend on the cutoff function and the choice of coarse-graining scale. As discussed in [75],

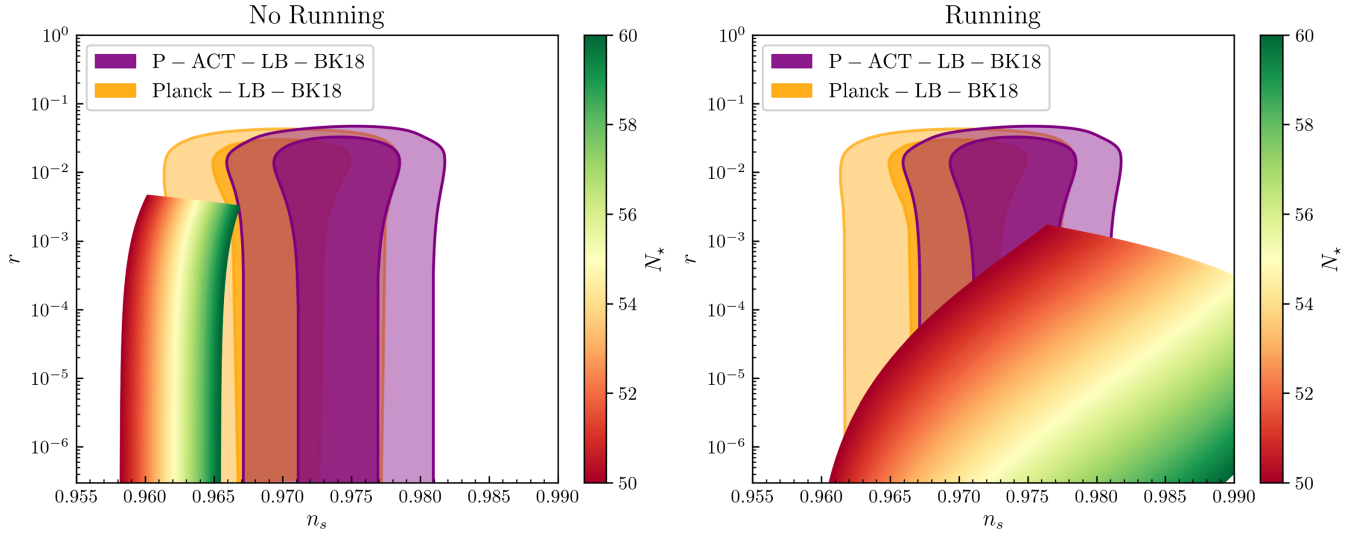


FIG. 2. Tensor-to-scalar ratio r as a function of the spectral index n_s calculated with (right panel) or without (left panel) including the ERG flow of the potential.

though, we have the freedom to choose the initial potential for the Wilsonian flow, and this renormalization-scheme dependence can be compensated by a different initial potential. We expect in future work to find a combination (initial potential, cutoff function, and coarse-graining scale) which leaves physical predictions invariant.

While this work proposes a new framework for the study of inflation, it also opens several new avenues that ought to be explored in the future. First of all, in this work, we neglected metric fluctuations and treated gravity at the level of a fixed FLRW background; i.e., we ignored both scalar-metric perturbation mixing and quantum graviton loops. Including scalar-metric mixing is known to induce curvature-scale corrections to the inflaton dynamics that can become quantitatively important beyond slow roll but are not expected to qualitatively modify the structure of the effective potential (see, e.g., [85]). Quantum graviton loops are further suppressed by H^2/M_p^2 in the semiclassical regime and are not expected to dominate the RG flow unless gravity becomes strong (see, e.g., [86,87]). While secular graviton effects have been discussed, they typically require an incredibly large number of e -folds and their interpretation for local observables is subtle (see, e.g., [88,89]). We, therefore, expect metric fluctuations to mainly renormalize numerical coefficients and shift the precise location of the dynamical destabilization identified here, without qualitatively altering the mechanism or its impact on inflationary observables.

For these reasons, it is apparent that this framework cannot be used to properly describe the behavior of the field after inflation ends, and it is possible that including metric fluctuations in the calculation—or using alternative regularization methods that do not feature explicit spinodal instabilities [90] or adopting an alternative Wilsonian renormalization scheme relevant to cosmology [91]—may help

pushing our understanding of renormalization beyond the sole study of slow-roll inflation. Finally, our treatment does not account for the stochastic effect exerted by quantum modes that exit the horizon on the classical background field characteristic of open-EFT systems. Upgrading this framework using the closed-time-path formalism [92–95] is, thus, something we are planning to explore in following studies.

By dynamically incorporating the effect of quantum corrections into the description of the Universe’s evolution, our work sets the path for a new approach to scalar theories of the early Universe. This may lead the community to reconsider the impact of cosmic inflation theories on cosmological observables, as well as broader classes of theories, such as quintessence scenarios, early dark energy models, or even axionlike particle phenomenology.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX: JUSTIFICATION OF THE LPA

We evaluate here the order of magnitude of the wave function renormalisation $Z - 1$ of the inflaton, in order to justify the LPA (Z is defined by the dressed term $Z(\dot{\phi})^2/2$). For this, we consider the one-loop contribution in the adiabatic approximation. Hence we assume that quantum fluctuations dress the inflaton quickly compared to the expansion of the universe, and quantum corrections are calculated as if the scale factor a and the background field ϕ were constant.

Starting from the potential

$$U(\phi + \varphi) = V_0(1 - e^{-(\phi+\varphi)/M})^2, \quad (\text{A1})$$

an expansion in powers of φ leads to all power-law self-interactions on the background ϕ . As we show below, the cubic interaction leads to a finite wave function renormalisation, which is small and independent of the scale factor, and therefore can be neglected. The quartic interaction does not renormalise the wave function at one loop, since the resulting self energy does not depend on the external momentum. Higher-order interactions appear in higher-loop-order diagrams only.

The mass and the cubic coupling constant (both defined at the minimum $\varphi = -\phi$) are given by

$$m^2 = \frac{2V_0}{M^2}, \quad \lambda_3 = -\frac{6V_0}{M^3}, \quad (\text{A2})$$

with $[\lambda_3] = 1$. The imaginary-time one-loop self energy Σ arising from the cubic interaction, and for an external momentum $k_\mu = (\omega, 0)$, is given by

$$\begin{aligned} \sqrt{g}\Sigma(\omega) &= \frac{(\sqrt{g}\lambda_3)^2}{2} \int \frac{dp_0}{2\pi} \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{g}[p_0^2 + (p/a)^2 + m^2]} \\ &\times \frac{1}{\sqrt{g}[(p_0 + \omega)^2 + (p/a)^2 + m^2]}, \end{aligned} \quad (\text{A3})$$

where the integral over $\vec{p}$ is regularised by a cutoff Λ , such that

$$\begin{aligned} \Sigma(\omega) &= \frac{\lambda_3^2}{8\pi^3} \int_{-\infty}^{\infty} dp_0 \int_0^{\Lambda/a} p^2 dp \\ &\times \frac{p^2 dp}{[p_0^2 + p^2 + m^2][(p_0 + \omega)^2 + p^2 + m^2]}. \end{aligned} \quad (\text{A4})$$

The wave-function renormalisation in real time ($\omega^2 \rightarrow -\omega^2$) is defined by the expansion

$$\Sigma(\omega) = \Sigma(0) - \frac{\omega^2}{2}(Z - 1) + \mathcal{O}(\omega^4), \quad (\text{A5})$$

where $\Sigma(0)$ is a correction to the mass term, which is logarithmically divergent. The quadratic order in ω is convergent, though

$$\begin{aligned} Z - 1 &= -\frac{d^2\Sigma}{d\omega^2} = \frac{\lambda_3^2}{4\pi^3} \int_{-\infty}^{\infty} dp_0 \int_0^{\Lambda/a} p^2 dp \\ &\times \left(\frac{1}{(p_0^2 + p^2 + m^2)^3} - \frac{4p_0^2}{(p_0^2 + p^2 + m^2)^4} \right), \end{aligned} \quad (\text{A6})$$

such that one can take the limit $\Lambda \rightarrow \infty$, where the scale factor does not contribute anymore. We have then

$$Z - 1 = \frac{1}{96\pi^2} \frac{\lambda_3^2}{m^2} = \frac{3V_0}{16\pi^2 M^4}, \quad (\text{A7})$$

and indeed $|Z - 1| \ll 1$.

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