

R^2 -Higgs inflation: R^3 contribution and preheating after ACT and SPT dataTanmoy Modak *Department of Physical Sciences, Indian Institute of Science Education and Research Berhampur, Berhampur 760003, Odisha, India*

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The R^2 -Higgs inflation is one of the simplest yet best-fit models consistent with Planck data. The higher spectral index n_s , recently reported by the combined cosmic microwave background (CMB) data from the Atacama Cosmology Telescope (ACT), South Pole Telescope (SPT), and Planck, along with baryonic acoustic oscillation data from the Dark Energy Spectroscopic Instrument (DESI), disfavors the single-field-like regime of R^2 -Higgs inflation at approximately 2σ . Following a doubly covariant formalism, we show that the R^2 -Higgs inflation, when modified by a dimension-six R^3 term, can account for the high n_s reported by CMB + BAO. In this regard, we find that preheating may play a pivotal role. We also show that if the nonminimal coupling between the Ricci scalar R and the Higgs field is $\mathcal{O}(10)$, then preheating via the production of Higgs quanta may help explain the reported observations.

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The cosmic inflation [1–3] offers an elegant solution to the so-called flatness, horizon, and exotic relic problems. The quantum fluctuations generated during inflation later transformed into density perturbations, which subsequently formed the large-scale structures such as the CMB anisotropies. The features predicted by the inflationary paradigm, such as the acoustic peaks in the CMB and the nearly scale-invariant adiabatic primordial power spectrum, are in excellent agreement with the WMAP [4] and Planck [5] data. While several predictions of the inflationary paradigm are in good agreement, the “last missing piece” of cosmic inflation, i.e., the primordial gravitational waves generated during inflation, is yet to be discovered. Ongoing experiments [6,7] are actively searching for this last missing piece which are believed to be engraved in the CMB B -mode polarization, with future missions proposed [8].

The spectral index n_s also provides a powerful test for inflationary dynamics. Indeed, several inflationary models have been ruled out based on their predicted values of n_s when compared to the observed 1σ and 2σ contours in the n_s vs r plane (r being the tensor-to-scalar ratio) of the Planck 2018 data [5]. Among other inflationary models, the Starobinsky or R^2 inflation [1,9–12] was one of the best-fit models and lay within the $\sim 1\sigma$ range in the n_s vs r plane of the Planck data [5]. The marginalized n_s from the

ACT [13], SPT [14], and the combined Planck + ACT + SPT (CMB-SPA) data [14] are still consistent with each other and found to be still on the lower side in the Λ CDM model.¹ However, the n_s is shifted toward a higher value in the Λ CDM model if one combines CMB with the BAO data from the DESI-DR2 collaboration [15]. This higher n_s is separately reported by both the ACT [13] and SPT [14] collaborations; the ACT collaboration found $n_s = 0.9752 \pm 0.0030$ by combining the CMB data from ACT, Planck (with lensing), and BAO data from DESI-DR2 (denoted as P-ACT-LB2) [13], whereas the SPT collaboration reported $n_s = 0.9728 \pm 0.0027$ by combining data from SPT, ACT, Planck, and DESI-DR2 (denoted as CMB-SPA + DESI) [14]. This high n_s from CMB + BAO excludes pure R^2 inflation at $>2\sigma$ [13,14].

The R^2 inflation model, while successful in many aspects, is not fully realistic. Following the discovery of the Higgs boson at the Large Hadron Collider, incorporating the Higgs field into inflationary dynamics has become essential. This model, i.e., the baseline Starobinsky model extended with the Higgs field and a

¹The marginalized n_s values within the Λ CDM model from different CMB experiments and their combinations are as follows: Planck: 0.9657 ± 0.0040 ; SPT-3G D1: 0.951 ± 0.011 ; ACT DR6: 0.9682 ± 0.0069 ; SPT + ACT: 0.9671 ± 0.0058 ; SPT + Planck: 0.9636 ± 0.0035 ; and combination of Planck + SPT + ACT i.e. CMB-SPA: 0.9684 ± 0.0030 , respectively [14]. Note that all of these experiments considered both T , E polarization data and lensing reconstruction and, the prior of τ_{reio} is taken to be the same as in Planck PR4. The n_s from the joint SPT-3G D1 + DESI is 0.949 ± 0.012 , which is actually much lower than both P-ACT-LB2 $n_s = 0.9752 \pm 0.0030$ [13] and from CMB-SPA+DESI $n_s = 0.9728 \pm 0.0027$ [14].

coupling between R and the Higgs field, dubbed the R^2 -Higgs inflation [16–27] (see also Refs. [28–31]), has similar inflationary predictions as R^2 inflation and is also a best-fit model to Planck data [5]. Like R^2 inflation, the R^2 -Higgs inflation is also a best-fit model to Planck data [5], but is now in tension with the measured high n_s from CMB + BAO [13,14], at least for the majority of the single-fieldlike regime. Several attempts have already been made to reconcile both R^2 inflation, Higgs inflation, and R^2 -Higgs inflation with the observed high n_s [32–49].

The inflationary observables, such as n_s and r , can receive significant modifications if one adds a dimension-six R^3 term to the action of R^2 -Higgs inflation [50,51]. Such curvature modifications are motivated not only by $f(R)$ theories of gravity but also from a purely phenomenological point of view [52–64] (see also Refs. [65–67]). In this paper, we study the implications of the R^3 term in R^2 -Higgs inflation to account for the observed high n_s (see Refs. [38,40] for a similar discussion). We show that the R^3 term can indeed account for the observed n_s , as well as the amplitude of the scalar power spectrum, as reported by the ACT [13] and SPT [14] collaborations. We adopt a doubly covariant formalism for our analysis, which is well suited for the two-field model under consideration. We derive all equations of motion (EoMs) for the background and perturbations without any assumptions other than the linearized approximation for perturbations from the metric and scalar sectors of the model. Our analysis stands in stark contrast to most existing studies in the literature, which predominantly follow the slow-roll approximation. We provide a comparison between the slow-roll approximated values of the amplitude of the power spectrum of curvature perturbations and n_s , and their corresponding values estimated by directly solving the EoMs. This is of particular importance, given that both observables are measured with high precision.

A precise understanding of the postinflationary reheating epoch is essential to match between the CMB scale, where the Λ CDM parameters such as n_s are measured and, the scale where inflation took place. In this regard, it has been shown that the reheating temperature is by far the most important parameter to account for the observed high n_s [33,39]. Practically all studies conducted so far that aim to account for the high n_s in the R^2 -Higgs inflation (or R^2 inflation for that matter) treated the reheating temperature as a free parameter. However, a thorough analysis is still lacking to determine whether the parameter space yielding a high n_s leads to perturbative reheating or whether thermalization proceeds via preheating. E.g., it has been shown that the presence of $\xi_H \sim 10$, where ξ_H is the nonminimal coupling between Higgs to Ricci scalar, may preheat the Universe via production of Higgs particles [68]. In our work, we study the impact of preheating in the R^3 modified R^2 -Higgs inflation and its connection in matching the inflationary and the CMB scale. We construct

gauge-invariant scalar perturbations considering all linear perturbations from scalar sector as well as those from the metric and compute the EoMs utilizing covariant formalism. The EoMs are then solved to determine the corresponding preheating temperature to match the scales. We show that the presence of the R^3 term may lead to faster preheating via Higgs production than that of R^2 -Higgs inflation without it. Our study illustrates that the R^3 term not only helps account for high n_s , it also impacts on the scale matching by inducing faster preheating.

The paper is organized as follows. In Sec. II, we summarize the details of the model and derive the relevant EoMs, followed by details of the inflationary dynamics in Sec. III. The preheating dynamics, along with the estimation of the preheating temperature and its impact on matching the CMB reference scale, is discussed in Sec. IV. We summarize our findings with an outlook in Sec. V.

II. THE ACTION

The action of the R^2 -Higgs inflation with the dimension six R^3 term in the Jordan frame is

$$S_J = \int d^4x \sqrt{-g_J} \times \left[\frac{M_P^2}{2} f(R_J, \Phi) - g_J^{\mu\nu} (\partial_\mu \Phi)^\dagger \partial_\nu \Phi - V(\Phi, \Phi^\dagger) \right], \quad (2.1)$$

where

$$f(R_J, \Phi) = R_J + \frac{\xi_R}{2M_P^2} R_J^2 + \frac{1}{3M_{P\xi_c}^4} R_J^3 + \frac{2\xi_H}{M_P^2} |\Phi|^2 R_J, \quad (2.2)$$

$$V(\Phi, \Phi^\dagger) = \lambda |\Phi|^4, \quad (2.3)$$

and, $M_P = \sqrt{1/(8\pi G)} \approx 2.4 \times 10^{18}$ GeV is the reduced Planck mass with G being Newton's gravitational constant.² The chosen metric convention is $(-1, +1, +1, +1)$ and $\sqrt{-g_J}$ denotes the determinant of the metric and R_J is the space-time Ricci scalar. The Φ is the Higgs field with hypercharge +1. Note that here we ignored the mass term in the Higgs potential since it does not have any significant impact on the inflationary dynamics under consideration.

It is in general more convenient to study the dynamics of R^2 -Higgs inflation in the Einstein frame. To transform the generic $f(R_J, \Phi)$ theory action in Eq. (2.1) to the scalar-tensor theory we first introduce an auxiliary field Ψ

²Note that, in references such as [38,50,60] a different parametrization of $f(R)$ is followed $f(R) = (R + \frac{1}{6M^2} R^2 + \frac{\delta_3}{36M^4} R^3)$. Here parameters M and δ_3 are related to ξ_R and ξ_c as $M^2 = \frac{M_P^2}{3\xi_R}$ and $\delta_3 = \frac{4}{3\xi_R^2 \xi_c}$.

and perform a Legendre transformation. The action is rewritten as

$$S_J = \int d^4x \sqrt{-g_J} \left[\frac{M_P^2}{2} \left(f(\Psi, \Phi) + \frac{\partial f(\Psi, \Phi)}{\partial \Psi} (R_J - \Psi) \right) - g_J^{\mu\nu} (\partial_\mu \Phi)^\dagger \partial_\nu \Phi - V(\Phi, \Phi^\dagger) \right]. \quad (2.4)$$

The Legendre transformation is well defined as long as $f(R, \Phi)$ is convex [60]. This leads to constrained relationship $\Psi > -\frac{1}{2} M_P^2 \xi_c \xi_R$. It is customary to introduce a physical degree of freedom

$$\Theta = \frac{\partial f(\Psi, \Phi)}{\partial \Psi}, \quad (2.5)$$

and rewrite the action in Eq. (2.4) as

$$S_J = \int d^4x \sqrt{-g_J} \times \left[\frac{M_P^2}{2} \Theta R_J - U(\Theta, \Phi) - g_J^{\mu\nu} (\partial_\mu \Phi)^\dagger \partial_\nu \Phi - V(\Phi, \Phi^\dagger) \right]. \quad (2.6)$$

Note that Eq. (2.5) has two solutions and we have chosen

$$\Psi = -\frac{1}{2} M_P^2 \xi_c (\xi_R - \zeta(\Theta, \Phi)) \quad (2.7)$$

where,

$$\zeta(\Theta, \Phi) = \sqrt{\xi_R^2 + \frac{4}{\xi_c} \left(\Theta - 1 - \frac{2\xi_H(\Phi^\dagger \Phi)}{M_P^2} \right)}. \quad (2.8)$$

The convexity condition here is satisfied during inflation if $\xi_c > 0$ for $\xi_R \gg \xi_H$. The potential $U(\Theta, \Phi)$ with two degrees of freedom Θ and Φ takes the form

$$\begin{aligned} U(\Theta, \Phi) &= \frac{M_P^2}{2} [\Psi \Theta - f(\Psi, \Phi)] \\ &= \frac{M_P^4 \xi_c^2}{48} \left[(\xi_R - \zeta(\Theta, \Phi)) \right. \\ &\quad \times (\xi_R^2 + \xi_R \zeta(\Theta, \Phi) - 2\zeta(\Theta, \Phi)^2) \left. \right]. \end{aligned} \quad (2.9)$$

For the other solution of Eq. (2.5), we remark that convexity is possible for $\xi_c < 0$; however, negative ξ_c may lead to unbounded $U(\Theta, \Phi)$ from below. After performing the Weyl rescaling of the metric $g_J^{\mu\nu} = \Theta g_E^{\mu\nu}$, we get the Einstein frame action

$$S_E = \int d^4x \sqrt{-g_E} \times \left[\frac{M_P^2}{2} R_E - \frac{3M_P^2}{4} g_E^{\mu\nu} \partial_\mu (\ln \Theta) \partial_\nu (\ln \Theta) - \frac{1}{2\Theta} g_E^{\mu\nu} (\partial_\mu \Phi)^\dagger \partial_\nu \Phi - V_E \right] \quad (2.10)$$

with

$$V_E = \frac{1}{\Theta^2} [V(\Phi, \Phi^\dagger) + U(\Theta, \Phi)], \quad (2.11)$$

$$R_J = \Theta \left[R_E + 3\Box_E \Theta - \frac{3}{2} g_E^{\mu\nu} \partial_\mu (\ln \Theta) \partial_\nu (\ln \Theta) \right]. \quad (2.12)$$

We have ignored the surface term $\Box_E = g_E^{\mu\nu} \partial_\mu \partial_\nu$ in the action S_E . With a field redefinition

$$\phi = M_P \sqrt{\frac{3}{2}} \ln \Theta, \quad (2.13)$$

and decomposing the Higgs field in the unitary gauge

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ h \end{pmatrix}, \quad (2.14)$$

the action Eq. (2.10) takes the form

$$S_E = \int d^4x \sqrt{-g_E} \left[\frac{M_P^2}{2} R_E - \frac{1}{2} G_{IJ} g_E^{\mu\nu} D_\mu \phi^I D_\nu \phi^J - V_E(\phi^I) \right], \quad (2.15)$$

where

$$\begin{aligned} V_E(\phi^I) &= e^{-2\sqrt{\frac{2}{3}} \frac{\phi}{M_P}} \left[\frac{\lambda}{4} h^4 + \left\{ \frac{M_P^4 \xi_c^2}{48} [(\xi_R - \tilde{\zeta}(\phi, h))(\xi_R^2 + \xi_R \tilde{\zeta}(\phi, h) - 2\tilde{\zeta}(\phi, h)^2)] \right\} \right], \\ \text{with } \tilde{\zeta}(\phi, h) &= \sqrt{\xi_R^2 + \frac{4}{\xi_c} \left(\Theta - 1 - \frac{\xi_H h^2}{M_P^2} \right)}. \end{aligned} \quad (2.16)$$

The G_{IJ} is the 2×2 metric for the field space manifold $\phi^I \in \{\phi, h\}$ which has only two nonzero elements $G_{\phi\phi} = 1$ and $G_{hh} = e^{-\sqrt{\frac{2}{3}} \frac{\phi}{M_P}}$.

III. INFLATIONARY DYNAMICS

A. Background and perturbation

In order to study the inflationary dynamics, we first need to find the equation of motion (EoM) of the scalar fields $\phi^I \in \{\phi, h\}$, which is obtained from varying the action Eq. (2.15),

$$\square\phi^K + \Gamma_{IJ}^K g_E^{\alpha\mu} D_\alpha \phi^I D_\mu \phi^J - G^{KM} V_{E,M} = 0, \quad (3.1)$$

where Γ_{IJ}^K are the Christoffel symbols for the field-space metric. Due to the multifield nature of the action and presence of noncanonical kinetic terms, we closely follow the covariant formalism as detailed in Refs. [69–71] at the linear order in perturbation. The fields $\phi^I(x^\mu)$ are decomposed into homogeneous background part (φ^I) and perturbation ($\delta\phi^I$) as

$$\phi^I(x^\mu) = \varphi^I(t) + \delta\phi^I(x^\mu), \quad (3.2)$$

where t is the cosmic time. We denote background fields $\varphi^I(t) = \{\varphi(t), h_0(t)\}$, i.e., the $\{\varphi(t)$ and $h_0(t)$ are the background fields for $\phi(x^\mu)$ and $\phi^I(x^\mu)$ respectively. The linear order perturbed Friedmann-Robertson-Walker (FRW) metric is given as [72–74]

$$ds^2 = -(1 + 2\mathcal{A})dt^2 + 2a(t)(\partial_i \mathcal{B})dx^i dt + a(t)^2[(1 - 2\psi)\delta_{ij} + 2\partial_i \partial_j \mathcal{E}]dx^i dx^j, \quad (3.3)$$

where $\mathcal{A}$, $\mathcal{B}$, ψ , and $\mathcal{E}$ are the metric scalar perturbations and $a(t)$ is scale factor. Throughout this work we consider the longitudinal gauge where the scalar perturbations $\mathcal{E}$ and $\mathcal{B}$ vanish.

At the linear order of perturbation, one can simply find the EoMs of the background fields from Eq. (3.1)

$$\mathcal{D}_I \dot{\varphi}^I + 3H\dot{\varphi}^I + G^{\phi J} V_{E,J}(\varphi^I) = 0, \quad (3.4)$$

where $\mathcal{D}_I$ and $\mathcal{D}_J$ field space covariant derivatives [69–71]

$$\mathcal{D}_I A^I = \dot{A}^I + \Gamma_{JK}^I \dot{\varphi}^J A^K, \quad (3.5)$$

$$\mathcal{D}_J A^I = \partial_J A^I + \Gamma_{JK}^I A^K. \quad (3.6)$$

The Hubble parameter is defined as

$$H^2 = \left(\frac{\dot{a}}{a}\right)^2 = \frac{1}{3M_{\text{Pl}}^2} \left(\frac{1}{2} G_{IJ} \dot{\varphi}^I \dot{\varphi}^J + V_0(\varphi^I)\right), \quad (3.7)$$

$$\dot{H} = -\frac{1}{2M_{\text{Pl}}^2} (G_{IJ} \dot{\varphi}^I \dot{\varphi}^J). \quad (3.8)$$

The slow-roll parameter ϵ and the number of e -foldings $\mathcal{N}$ are expressed as

$$\epsilon = -\frac{\dot{H}}{H^2}, \quad \mathcal{N}(t) \equiv \ln \frac{a(t)}{a(t_{\text{end}})}, \quad (3.9)$$

where t_{end} is the corresponding cosmic time when the inflation ends and $a(t_{\text{end}})$ is the scale factor at the end of inflation. We denote the end of inflation when $\epsilon(t = t_{\text{end}}) = 1$. In what follows we shall use the t and $\mathcal{N}$ interchangeably. The energy density associated with background fields is

$$\rho_{\text{inf}} = \frac{1}{2} G_{IJ} \dot{\varphi}^I \dot{\varphi}^J + V_0(\varphi^I), \quad (3.10)$$

where G_{IJ} is evaluated at the background field order.

Before moving to the perturbation dynamics let us first understand the impact of the ξ_c . Our primary aim is to investigate whether the single fieldlike regime of R^2 -Higgs inflation could still be a viable parameter space to account high n_s after adding the R^3 term. The single fieldlike regime can be obtained by the so-called valley approximation by solving $\frac{\partial V_{\text{eff}}}{\partial h} = 0$ for h in terms of ϕ and, inserting the solution back to V_E [51]. We denote this effective single field potential as $W_{\text{eff}}(\phi)$. In the valley approximation, for suitable values of ξ_R , ξ_H , and ξ_c , the $W_{\text{eff}}(\phi)$ acts as effective inflaton potential and ϕ plays the role of inflaton. In Fig. 1, we plot $W_{\text{eff}}(\phi)$ for different ξ_H , ξ_R , and ξ_c . The parameter λ is kept at a fixed value 10^{-2} throughout this paper for simplicity. It is clear that if ξ_c becomes large the potential becomes more flat i.e. akin to the case of pure

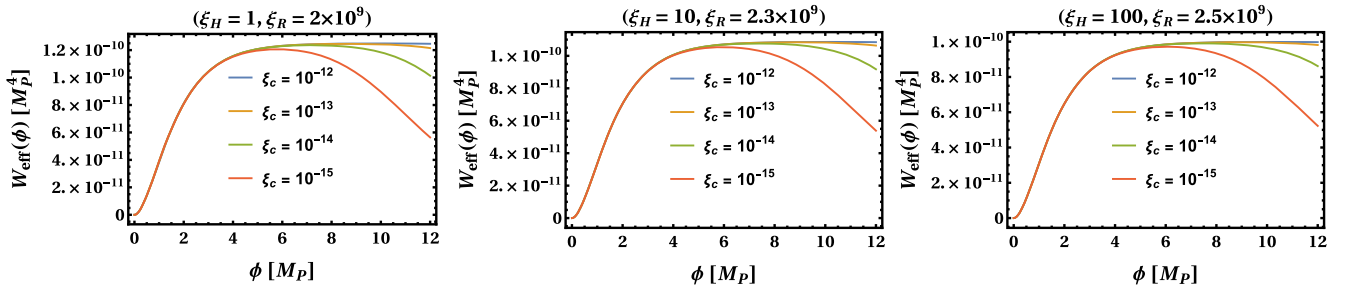


FIG. 1. The effective potential $W_{\text{eff}}(\phi)$ vs ϕ plots for $(\xi_H = 1, \xi_R = 2 \times 10^9)$, $(\xi_H = 10, \xi_R = 2.3 \times 10^9)$, and $(\xi_H = 100, \xi_R = 2.5 \times 10^9)$ for different values of ξ_c .

TABLE I. Two benchmark points for our analysis. Scales are given in units of the Planck mass M_P . See text for details.

BP	ξ_R	ξ_H	ξ_c	$\varphi(t_{\text{in}}) [M_P]$	$h_0(t_{\text{in}}) [M_P]$
<i>a</i>	2.12×10^9	1.5	7×10^{-14}	5.32	9.8×10^{-5}
<i>b</i>	2.3×10^9	10	8×10^{-15}	5.35	3×10^{-7}

R^2 -Higgs inflation. For smaller ξ_c , the $W_{\text{eff}}(\phi)$ does not remain asymptotically flat for large ϕ . This is expected, since in the R^3 term ξ_c sits in the denominator [see Eq. (2.3)]; hence a smaller ξ_c impacts more in changing the shape of $W_{\text{eff}}(\phi)$. In the following we shall see that $\xi_c \sim 10^{-13} - 10^{-14}$ is sufficient to induce large n_s to match the reported value of CMB + BAO.

We utilize the valley approximation as a guideline to identify a couple of single-fieldlike benchmark points (BPs) as given in Table I. In what follows, we take these BPs as reference and solve the coupled EoMs in Eq. (3.4) directly with the full potential $V_E(\phi^I)$. The initial time for the numerical analysis is set to $t_{\text{in}} = 0$. Note that Eq. (3.7) is solved simultaneously taking $\ln(a)$ as a variable with initial condition $\ln(a(t_{\text{in}})) = 0$. For consistency, we have checked the $\dot{H}$ estimated from Eq. (3.8) matches with the one taking time derivative of Eq. (3.7) within second decimal place. We remark here that the exact solution of the EoMs is required, rather than the slow-roll approximation, due to the high precision of current measurements of parameters such as n_s . In Fig. 2, we plot the field evolution and Hubble parameter for BP*a* for illustration. Note that we solved the Eq. (3.4) numerically with initial conditions $\varphi(t_{\text{in}})$ as summarized in Table I and $\dot{\varphi}(t_{\text{in}}) = 0$.

B. Dynamics of the perturbations

The cosmological perturbations generated from the field fluctuations $\delta\phi^I(x^\mu)$ can account for the observed nearly scale invariant adiabatic curvature perturbation. In general $\delta\phi^I(x^\mu)$ are gauge-dependent quantities but one can conveniently construct the covariant field fluctuations $\mathcal{Q}^I(x^\mu)$ which relate the $\phi^I(x^\mu)$ to corresponding $\varphi^I(t)$ via a unique geodesic in the field-space where $\delta\phi^I$ is expressed as [69,75]

$$\delta\phi^I = \mathcal{Q}^I - \frac{1}{2}\Gamma_{JK}^I \mathcal{Q}^K \mathcal{Q}^J + \frac{1}{3!}(\Gamma_{MN}^I \Gamma_{JK}^N - \Gamma_{JK,M}^I) \mathcal{Q}^K \mathcal{Q}^J \mathcal{Q}^M + \dots \quad (3.11)$$

At the linear order $\mathcal{Q}^I = \delta\phi^I$ and the gauge-independent Mukhanov-Sasaki variables are [73,76,77]

$$\mathcal{Q}^I = \mathcal{Q}^I + \frac{\dot{\varphi}^I}{H} \psi = \delta\phi^I + \frac{\dot{\varphi}^I}{H} \psi. \quad (3.12)$$

Note that is $\mathcal{Q}^I$ s are doubly covariant with respect both the space-time and field-space transformations. In the field-space manifold $\dot{\varphi}^I$ and $\mathcal{Q}^I$ transform like vectors. We may now insert Eqs. (3.12) and (3.3) into Eq. (3.1) to find EoMs for $\mathcal{Q}^I$ s at linear order as

$$\mathcal{D}_i^2 \mathcal{Q}^I + 3H\mathcal{D}_i \mathcal{Q}^I - \frac{\nabla^2}{a^2} \mathcal{Q}^I + \mathcal{M}_J^I \mathcal{Q}^J = 0, \quad (3.13)$$

with

$$\mathcal{M}_L^I = G^{IJ}(\mathcal{D}_L \mathcal{D}_J V_E) - \mathcal{R}_{JKL}^I \dot{\varphi}^J \dot{\varphi}^K - \frac{1}{M_P^2 a^3} \mathcal{D}_I \left(\frac{a^3}{H} \dot{\varphi}^I \dot{\varphi}_L \right), \quad (3.14)$$

where the $\mathcal{R}_{JKL}^I$ is the field-space Riemann tensor which is evaluated, along with $\mathcal{M}_L^I$, at the background order.

The two independent perturbations $\mathcal{Q}^\phi$ and $\mathcal{Q}^h$ can be decomposed into adiabatic and isocurvature perturbations; however, we first need two unit vectors $\hat{\sigma}^I$ and $\hat{\omega}^I$. The former is defined as

$$\hat{\sigma}^I = \frac{\dot{\varphi}^I}{\dot{\sigma}}, \quad (3.15)$$

with $\dot{\sigma} = \sqrt{G_{IJ} \dot{\varphi}^I \dot{\varphi}^J}$, while the latter is defined as

$$\hat{\omega}^I = \frac{\omega^I}{\omega}, \quad (3.16)$$

where ω^I is called “turning vector” defined as $\omega^I = \mathcal{D}_I \hat{\sigma}^I$ and the magnitude $\omega = |\omega^I| = \sqrt{G_{IJ} \omega^I \omega^J}$. It should be

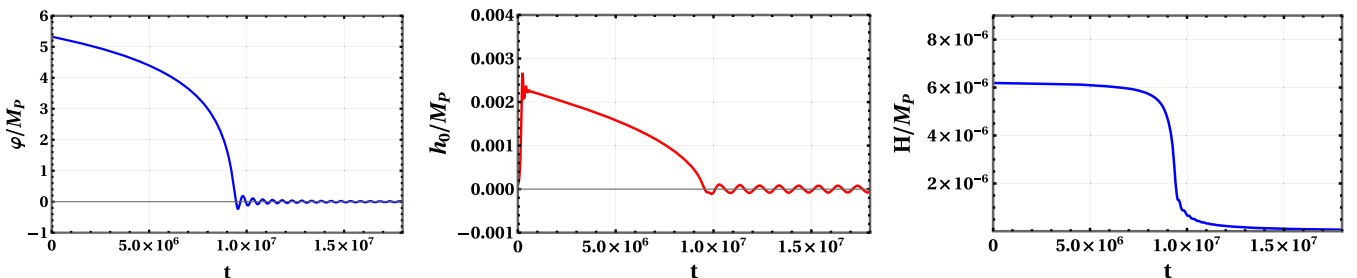


FIG. 2. The evolution of φ , h_0 , and H with respect to cosmological time t for BP*a*.

noted that $\omega_I \delta^I = 0$. We can now decompose the curvature perturbations and isocurvature perturbations as

$$Q_\sigma = \hat{\sigma}_I Q^I, \quad Q_s = \hat{\omega}_I Q^I. \quad (3.17)$$

The gauge invariant curvature (adiabatic) and isocurvature perturbations are

$$\mathcal{R} = \frac{H}{\dot{\sigma}} Q_\sigma, \quad \mathcal{S} = \frac{H}{\dot{\omega}} Q_s. \quad (3.18)$$

The dimensionless power spectra for the adiabatic and entropy perturbations are [73,74,78]

$$\mathcal{P}_{\mathcal{R}}(t; k) = \frac{k^3}{2\pi^2} |\mathcal{R}|^2, \quad (3.19)$$

$$\mathcal{P}_{\mathcal{S}}(t; k) = \frac{k^3}{2\pi^2} |\mathcal{S}|^2. \quad (3.20)$$

One can now readily evaluate the power spectrum from Eqs. (3.19) and (3.20) for a given Fourier mode k . The power spectrum for curvature perturbation freeze after it exits horizon. Therefore we simply evaluate $\mathcal{P}_{\mathcal{R}}$ for different modes at the end of inflation. The power spectrum of the isocurvature mode $\mathcal{P}_{\mathcal{S}}(t; k)$ on the other hand may change in the superhorizon scales. This is primarily due to the small but nonvanishing off-diagonal elements of M_J^I which induces mild power transfer from adiabatic to isocurvature mode. As we expect $\mathcal{P}_{\mathcal{R}}$ to be orders of magnitude larger during inflation compared to $\mathcal{P}_{\mathcal{S}}(t; k)$, a small power transfer from $\mathcal{P}_{\mathcal{R}}$ shall induce large change in $\mathcal{P}_{\mathcal{S}}(t; k)$. Finally, the spectral index n_s is given as

$$n_s = 1 + \frac{d \ln \mathcal{P}_{\mathcal{R}}(k)}{d \ln k}. \quad (3.21)$$

The measured values of inflationary parameters are [6,14]

$$\log(10^{10} \mathcal{A}_{s*}) = 3.0574 \pm 0.0094, \quad (3.22)$$

$$n_s^* = 0.9728 \pm 0.0027, \quad (3.23)$$

$$r_* \lesssim 0.036 \quad \text{at 95\% CL}, \quad (3.24)$$

where $*$ denotes the respective parameters are measured at a reference scale $k_{\text{ref}} = 0.05 \text{ Mpc}^{-1}$ and $\mathcal{A}_{s*}$ is amplitude of the power spectrum of the curvature perturbation evaluated at k_{ref} . Here we have considered CMB-SPA+DESI from Ref. [14] which includes all CMB measurements as well as BAO data from DESI.

In Fig. 3 we show the $\mathcal{P}_{\mathcal{R}}$ and n_s as a function of Fourier modes k evaluated from Eqs. (3.19) and (3.21). The reference mode $k_{\text{ref}} = 0.05 \text{ Mpc}^{-1}$ today corresponds to a mode $k_* = a(t_*)H(t_*)$ which exits the horizon at time t_* before the end of inflation such that the corresponding power spectrum $\mathcal{P}_{\mathcal{R}}(t_*; k_*) = \mathcal{A}_{s*}$ and n_s evaluated at $k = k_*$ is equal to n_s^* . From Fig. 3, we find that the $k_* = 3.4624 \times 10^{-5} (1.3843 \times 10^{-4}) M_P$, for BP_a (BP_b), which corresponds to $t_* = 2.784 \times 10^5 (5.3037 \times 10^5)$ and $\mathcal{P}_{\mathcal{R}}(k_*) = 2.102 \times 10^{-9} (2.104 \times 10^{-9})$ and $n_s^* = 0.9714 (0.9733)$. The tensor-to-scalar ratio is estimated simply by utilizing single-field approximation $r \approx 16\epsilon(t_*)$. It is justified since the observable is not measured and only the upper limit exists. We find $r_* \approx 3.75 \times 10^{-3} (3.45 \times 10^{-3})$ for the respective BPs.

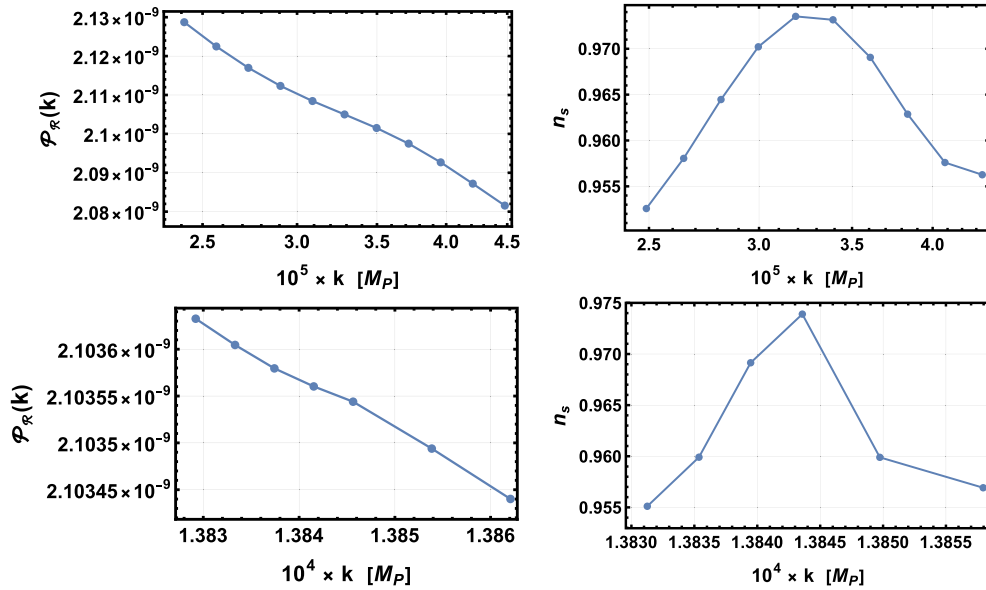


FIG. 3. The $\mathcal{P}_{\mathcal{R}}(k)$ and n_s vs k for both the BPs. Here the n_s is evaluated directly utilizing Eq. (3.21) and $\mathcal{P}_{\mathcal{R}}(k)$ via Eq. (3.19).

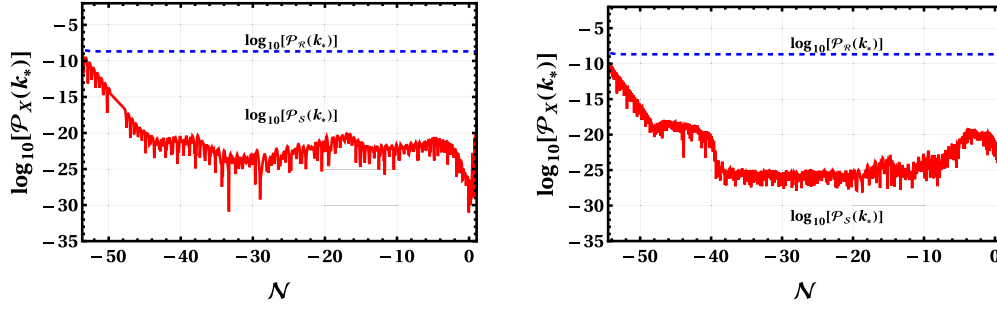


FIG. 4. The evolution of $\mathcal{P}_R(k_*)$ and $\mathcal{P}_S(k_*)$ after horizon exit for both the BPs.

To generate Fig. 3, we note that the EoMs in Eq. (3.13) are solved for each Fourier modes k using initial condition [79]

$$Q^I(t_{\text{in}}) \simeq \frac{H}{\sqrt{2}k^3} \left(i + \frac{k}{aH} \right) e^{i\frac{k}{aH}}, \quad (3.25)$$

with all relevant modes initialized at $t = t_{\text{in}}$ to ensure they are well within the horizon. In Fig. 4, we plot the evolution of $\mathcal{P}_R(k_*)$ and $\mathcal{P}_S(k_*)$ as a function of cosmic time. For both BPs, as expected, $\mathcal{P}_S$ is orders of magnitude lower than that of $\mathcal{P}_R$ during inflation for the reference mode k_* . Moreover, as explained above, $\mathcal{P}_R(k_*)$ remains frozen but, $\mathcal{P}_S(k_*)$ changes during inflation, albeit its magnitude is still much smaller than the former. The number of e -folding when the reference mode exits horizon before the end of inflation is $\Delta N_* = \mathcal{N}|_{t=t_*} - \mathcal{N}|_{t=t_{\text{end}}}$, which is 53.66 and 54.51 for BPb and BPc, respectively. At this point it is useful to compare the $\mathcal{A}_{s*}$ and n_{s*} values obtained above to the ones from slow-roll approximation. In slow-roll approximation

$$\mathcal{A}_s(t) = \frac{V_E(t)}{24\pi^2 M_{\text{Pl}}^4 \epsilon(t)}, \quad (3.26)$$

where T_{pre} and $T_0 = 2.7K$ are the temperature at the end of preheating and today. The number of relativistic degrees of freedom at the end of preheating and today are $g_{\text{pre}} = 106.75$ and $g_0 = 43/11$ respectively and $\mathcal{N}_{\text{pre}} = \ln(a(t_{\text{pre}})/a(t_{\text{end}}))$ is the number of e -folding elapsed from the end of inflation to the end of preheating. Note here that while finding Eq. (3.29) we have assumed that the thermalization after is completed within one Hubble time as in Ref. [25]. We may now express all dimensionful quantities

$$n_s(t) \approx 1 - 6\epsilon(t) + 2\eta(t). \quad (3.27)$$

Once evaluated at $t = t_*$ these expressions translate to $n_s^* = 0.9652$ (0.9746) and $\mathcal{A}_{s*} = 2.068 \times 10^{-9}$ (2.0465×10^{-9}) for BPa (BPb). This emphasizes that the slow-roll approximation is perhaps not ideal given the high precision recent cosmological measurements.

The task now is to match the reference mode k_* in units of M_{Pl} to CMB reference mode $k_{\text{ref}}/a_0 = 0.05 \text{ Mpc}^{-1}$. The k_{ref} is connected to k_* via the relationship

$$k_{\text{ref}} = k_* = a(t_*)H(t_*) = \frac{a(t_*)}{a(t_{\text{end}})} \frac{a(t_{\text{end}})}{a(t_{\text{pre}})} \frac{a(t_{\text{pre}})}{a_0} a_0 H(t_*), \quad (3.28)$$

where $a(t_{\text{pre}})$ is the corresponding scale factor when preheating is completed. Here we simply assumed that the radiation dominated era has immediately started after the end of preheating. The explicit details of the preheating epoch for the BPa and BPb is discussed shortly. One can reexpress Eq. (3.28) as [25]

$$\begin{aligned} \Delta N_* &= \ln\left(\frac{a(t_{\text{end}})}{a(t_*)}\right) = \ln\left(\frac{H(t_*)}{(k_{\text{ref}}/a_0)}\right) + \ln\left(\frac{a(t_{\text{end}})}{a(t_{\text{pre}})}\right) + \ln\left(\frac{a(t_{\text{pre}})}{a_0}\right) \\ &= \ln\left(\frac{H(t_*)}{(k_{\text{ref}}/a_0)}\right) - \mathcal{N}_{\text{pre}} + \ln\left\{\frac{T_0}{T_{\text{pre}}} \left(\frac{g_0}{g_{\text{pre}}}\right)^{1/3}\right\}, \end{aligned} \quad (3.29)$$

in M_{Pl} units and match the right-hand side of Eq. (3.29) i.e. ΔN_* to the left-hand side. While ΔN_* was already found above, parameters $\mathcal{N}_{\text{pre}}$ and T_{pre} are yet to be estimated. A detailed discussion on them is deferred to Sec. IV.

IV. PREHATING

In order to find preheating we write the second order action for the fluctuations Q^I (with $I = \{1, 2\}$) as [69,70,80,81]

$$S_{(Q)}^{(2)} = \int d^3x dt a^3 \left[-\frac{1}{2} \bar{g}_E^{\mu\nu} G_{IJ} \mathcal{D}_\mu Q^I \mathcal{D}_\nu Q^J - \frac{1}{2} \mathcal{M}_{IJ} Q^I Q^J \right], \quad (4.1)$$

where $\bar{g}_E^{\mu\nu} \equiv (-1, a^2(t), a^2(t), a^2(t))$ is the spatially flat unperturbed FLRW metric and $\mathcal{M}_{IJ}$ and G_{IJ} are computed at the background order. However, it is more convenient to quantize the fields for preheating in conformal time τ with invariant $ds^2 = a^2(\tau) \eta_{\mu\nu} dx^\mu dx^\nu$. The second order action in conformal time with transformation $\partial_0 \rightarrow \partial_\tau/a$ and field rescaling $X^I(x^\mu) \equiv a(t)Q^I(x^\mu)$ is given as

$$S_{(X)}^{(2)} = \int d^3x d\tau \left[-\frac{1}{2} \eta^{\mu\nu} G_{IJ} (\mathcal{D}_\mu X^I) (\mathcal{D}_\nu X^J) - \frac{1}{2} \mathcal{M}_{IJ} X^I X^J \right], \quad (4.2)$$

with $\eta^{\mu\nu} = (-1, 1, 1, 1)$.

$$\mathcal{M}_{IJ} = a^2 \left(\mathcal{M}_{IJ} - \frac{1}{6} G_{IJ} R_E \right), \quad \text{with } R_E = \frac{6a''}{a^3}, \quad (4.3)$$

where we have used the shorthand notation $(')$ conformal time derivative. The EoMs of the scalar field fluctuations $X^I(x^\mu)$ can now readily be derived from either Eq. (4.2) or from Eq. (3.13),

$$\mathcal{D}_\tau^2 X^I - \left[\nabla^2 - a^2 \left(\mathcal{M}_I^I - \frac{1}{6} R_E G_I^I \right) \right] X^I = 0, \quad (4.4)$$

where we have used the diagonality of M_J^I . The energy momentum tensor for doubly covariant conformally rescaled field fluctuations X^I is

$$T_{\mu\nu}^{(X)} = G_{IJ} (\mathcal{D}_\mu X^I) (\mathcal{D}_\nu X^J) + \eta_{\mu\nu} \left[-\frac{1}{2} \eta^{\alpha\beta} G_{IJ} (\mathcal{D}_\alpha X^I) (\mathcal{D}_\beta X^J) - \frac{1}{2} \mathcal{M}_{IJ} X^I X^J \right]. \quad (4.5)$$

We transform the second order action in Eq. (4.2) to momentum space [68,70,80–82]

$$S_{(X)} = \int d\tau \mathcal{L}_{(X)} = \int d\tau \frac{d^3k}{(2\pi)^3} \times \left[\frac{1}{2} |\partial_\tau \tilde{X}^I|^2 - \frac{1}{2} \omega_{(I)}^2(\tau, k) |\tilde{X}^I|^2 \right] \quad (4.6)$$

with

$$\omega_{(I)}^2(\tau, k) = (k^2 + a^2 m_{\text{eff},(I)}^2(\tau))$$

$$\text{and } m_{\text{eff},(I)}^2(\tau) = \mathcal{M}_I^I - \frac{1}{6} R_E = \frac{1}{a^2} \mathcal{M}_I^I = m_{\text{eff},(I)}^2 = \sum_i m_{i,(I)}^2 \quad (4.7)$$

where we have denoted

$$m_{1,(I)}^2 = G^{(I)J} (\mathcal{D}_{(I)} \mathcal{D}_J V_E), \quad m_{2,(I)}^2 = -\mathcal{R}_{JK(I)} \dot{\phi}^J \dot{\phi}^K, \\ m_{3,(I)}^2 = -\frac{1}{M_{\text{Pl}}^2 a^3} \mathcal{D}_I \left(\frac{a^3}{H} \dot{\phi}^{(I)} \dot{\phi}_{(I)} \right), \quad m_{4,(I)}^2 = -\frac{R_E}{6}, \quad (4.8)$$

and the field-space indices (I) are not summed over. In Fig. 5 the $m_{\text{eff},(I)}^2(\tau)$ as a function of $\mathcal{N}$ is plotted for illustration. We see that in both the BPs $m_{\text{eff},(\phi)}^2(\tau)$ is tachyonic but becomes positive after the end of inflation. In contrast $m_{\text{eff},(h)}^2(\tau)$ is positive before the end of inflation while it oscillates around zero after the end of inflation. We postpone a discussion regarding consequences of them for a later part of this section.

The canonical momentum is defined as

$$\hat{\pi}^I(\tau, \mathbf{k}) = \partial_\tau \hat{X}^I(\tau, \mathbf{k}),$$

$$\text{with } [\hat{X}^I(\tau, \mathbf{k}), \hat{\pi}^J(\tau, \mathbf{q})] = i(2\pi)^3 \delta^{IJ} \delta^{(3)}(\mathbf{k} + \mathbf{q}), \quad (4.9)$$

where we have elevated the classical field fluctuations $\tilde{X}^I$ to their respective quantized versions $\hat{X}^I$. We now decompose the quantized fluctuations $\hat{X}^\phi$ and $\hat{X}^h$ in momentum space as [70,80]

$$\hat{X}^\phi = [(v_{1k}(\tau) e_1^\phi(\tau) \hat{a}_1(\mathbf{k}) + v_{2k}(\tau) e_2^\phi(\tau) \hat{a}_2(\mathbf{k})) + (v_{1k}^*(\tau) e_1^\phi(\tau) \hat{a}_1^\dagger(-\mathbf{k}) + v_{2k}^*(\tau) e_2^\phi(\tau) \hat{a}_2^\dagger(-\mathbf{k}))], \quad (4.10)$$

$$\hat{X}^h = [(y_{1k}(\tau) e_1^h(\tau) \hat{a}_1(\mathbf{k}) + y_{2k}(\tau) e_2^h(\tau) \hat{a}_2(\mathbf{k})) + (y_{1k}^*(\tau) e_1^h(\tau) \hat{a}_1^\dagger(-\mathbf{k}) + y_{2k}^*(\tau) e_2^h(\tau) \hat{a}_2^\dagger(-\mathbf{k}))], \quad (4.11)$$

$\hat{a}_m(\mathbf{k})$ and $\hat{a}_m^\dagger(-\mathbf{k})$ (with $m \in 1, 2$) are annihilation and creation operators which follow the commutator relationships

$$[\hat{a}_m(\mathbf{k}), \hat{a}_n(\mathbf{q})] = [\hat{a}_m^\dagger(\mathbf{k}), \hat{a}_n^\dagger(\mathbf{q})] = 0, \\ [\hat{a}_m(\mathbf{k}), \hat{a}_n^\dagger(\mathbf{q})] = (2\pi)^3 \delta_{mn} \delta^{(3)}(\mathbf{k} - \mathbf{q}), \quad (4.12)$$

with

$$\hat{a}_m(\mathbf{k})|0\rangle = 0, \quad \langle 0|\hat{a}_m^\dagger(\mathbf{k}) = 0. \quad (4.13)$$

The field space vielbeins follow the relationship

$$\delta^{mn} e_m^I(\tau) e_n^J(\tau) = G^{IJ}(\tau), \quad \text{and } \mathcal{D}_\tau e_J^m = 0, \quad (4.14)$$

for all m and J .

As our focus of interest is the single fieldlike regime, the off-diagonal elements are $\mathcal{M}_h^\phi \sim 0$ and $\mathcal{M}_\phi^h \sim 0$ and, hence $e_2^\phi \sim 0$, $e_1^h \sim 0$ [68,81]. Therefore, Eq. (4.5) becomes two decoupled source-free EoMs as

$$v_{1k}'' + \omega_{(\phi)}^2 v_{1k} \simeq 0, \quad (4.15)$$

$$y_{2k}'' + \omega_{(h)}^2 y_{2k} \simeq 0, \quad (4.16)$$

with $\omega_{(I)}^2$ given by Eq. (4.7). We solve the EoMs Eqs. (4.15) and (4.16) utilizing the Bunch-Davis (BD) initial condition

$$\begin{aligned} \lim_{t \rightarrow -\infty} v_{1k}(k, t) &= \lim_{t \rightarrow -\infty} y_{2k}(k, t) = \frac{e^{-\frac{ikt}{a}}}{\sqrt{2k}}, \\ \lim_{t \rightarrow -\infty} \dot{v}_{1k}(k, t) &= \lim_{t \rightarrow -\infty} \dot{y}_{2k}(k, t) = -\frac{i}{a} \sqrt{\frac{k}{2}} e^{-\frac{ikt}{a}}, \end{aligned} \quad (4.17)$$

where the relevant modes under consideration are initialized about $\mathcal{N} \sim -3$ i.e. $3e$ -foldings before the end of inflation to ensure they are well within the horizon.

The combined vacuum averaged comoving energy densities for the inflaton and Higgs fluctuations is [68,82]

$$\rho_{\phi h} = \int \frac{d^3 k}{(2\pi)^3} (\rho_k^{(\phi)} + \rho_k^{(h)}), \quad (4.18)$$

where $\rho_k^{(\phi)}$ and $\rho_k^{(h)}$ are the respective fluctuations per mode defined as

$$\begin{aligned} \rho_k^{(\phi)} &= \frac{1}{2} G_{\phi\phi} (|v_{1k}'|^2 + \omega_{(\phi)}^2 |v_{1k}|^2) e_1^\phi e_1^\phi \\ &= \frac{1}{2} (|v_{1k}'|^2 + \omega_{(\phi)}^2 |v_{1k}|^2), \end{aligned} \quad (4.19a)$$

$$\begin{aligned} \rho_k^{(h)} &= \frac{1}{2} G_{hh} (|y_{2k}'|^2 + \omega_{(h)}^2 |y_{2k}|^2) e_2^h e_2^h \\ &= \frac{1}{2} (|y_{2k}'|^2 + \omega_{(h)}^2 |y_{2k}|^2). \end{aligned} \quad (4.19b)$$

The physical energy densities of the fluctuation of ϕ is

$$\begin{aligned} \rho_{(\phi)} &= \frac{1}{a^4} \int \frac{d^3 k}{(2\pi)^3} \rho_k^{(\phi)} \\ &= \frac{1}{a^2} \int \frac{k^2}{4\pi^2} dk \left[|\dot{v}_{1k}|^2 + \left(\frac{k^2}{a^2} + |m_{\text{eff},(\phi)}^2(t)| \right) |v_{1k}|^2 \right], \end{aligned} \quad (4.20)$$

and a similar expression can be found for h . The vacuum subtracted quantum energy densities are given as [68]

$$\rho_{(I)}^q = \rho_{(I)} - \rho_{(I)}^{\text{BD}}, \quad \text{with} \quad \rho_{(I)}^{\text{BD}} = \frac{1}{a^4} \int dk \frac{k^3}{4\pi^2}, \quad (4.21)$$

where $\rho_{(I)}^{\text{BD}}$ is the associated BD vacuum energy densities for respective field.

The momentum upper limit in Eqs. (4.20) and (4.21) are obtained by finding the mode for which the relative error between ρ_k and ρ_k^{BD} is about 10% with $\rho_k > \rho_k^{\text{BD}}$ in each time step. For illustration, we have plotted the spectrum in Fig. 6 for different $\mathcal{N}$. The momentum upper limit is identified as the k value where the blue or red curves stop respectively slightly above the orange dotted BD spectrum (within 10% of each other with $\rho_k > \rho_k^{\text{BD}}$)

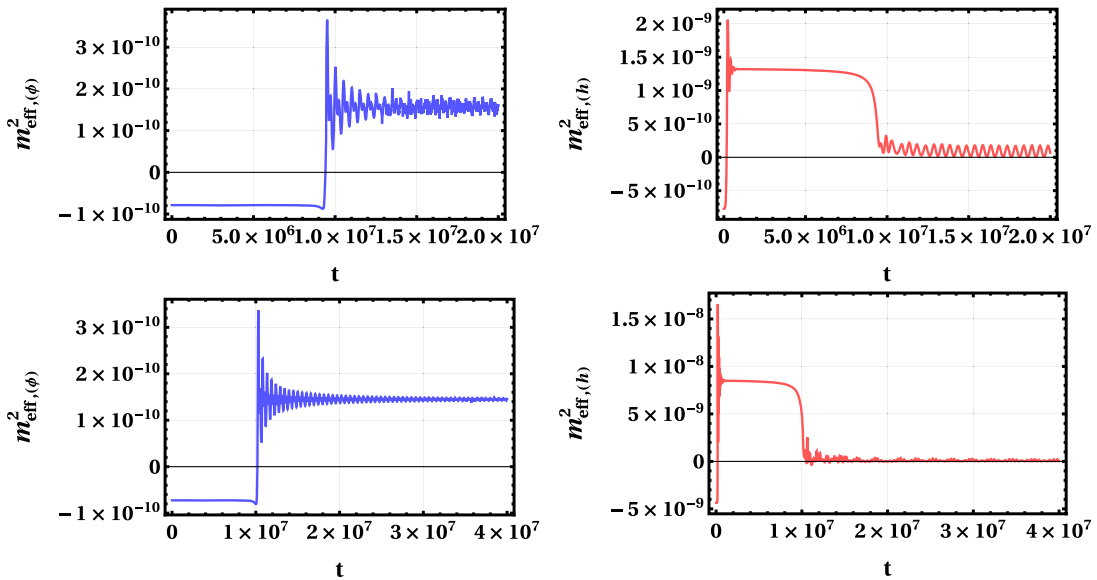


FIG. 5. The evolution of $m_{\text{eff},(I)}^2(\tau)$ as a function of $\mathcal{N}$. See text for details.

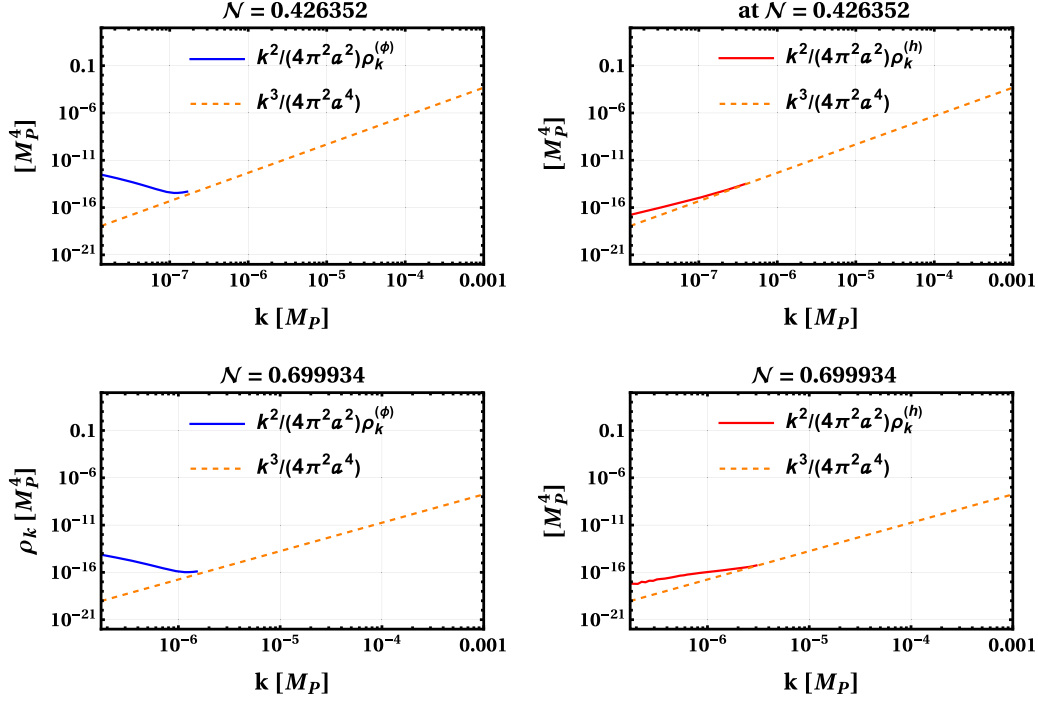


FIG. 6. The spectrum for inflaton fluctuations ($k^2/(4\pi^2 a^2)\rho_k^{(\phi)}$) and the BD vacuum ($k^3/(4\pi^2 a^4)$) are plotted in the left panels for both the BPs in solid blue and dotted orange respectively. In the right panels the corresponding spectrum of Higgs ($k^2/(4\pi^2 a^2)\rho_k^{(h)}$) fluctuations are plotted in solid against the BD vacuum in dotted orange. Note that these spectra are essentially the integrand of Eqs. (4.20) and (4.21) for the respective fluctuations and BD vacuum.

for specific $\mathcal{N}$ [68]. Modes larger than this k do not leave the vacuum in the respective time steps. The lower limit on the other hand is found by considering all subhorizon modes at a particular time step, since only these subhorizon modes take part in thermalization process in our single-fieldlike scenario. In practice, we follow an adaptive numerical framework where we consider all subhorizon modes at a particular time while the upper limit is obtained from the above-mentioned method. We redirect readers to Ref. [68] for details of this numerical approach.

We plot $\rho_{(\phi)}^q$ (blue), $\rho_{(h)}^q$ (red), and ρ_{inf} (black) in Fig. 7 for illustration. It is clear that $\rho_{(\phi)}^q$ is orders of magnitude lower than ρ_{inf} for both BP a and BP b . This is primarily because $m_{\text{eff},(\phi)}^2 < 0$; as a result, $\rho_{(\phi)}^q$ receives a tachyonic (exponential) amplification before the end of inflation for both BPs. However, after the end of inflation, $m_{\text{eff},(\phi)}^2 > 0$, and therefore $\rho_{(\phi)}^q$ experiences neither tachyonic growth nor parametric resonance. In contrast, $m_{\text{eff},(h)}^2 > 0$ before the end of inflation, but oscillates around the minimum after the end of inflation; as a result, $\rho_{(h)}^q$ experiences parametric resonance. We now denote the completion of preheating as the time when $\rho_{\text{inf}} = \rho_{\text{pre}} = \rho_{(X)}^q$ (with X being either ϕ or h), at a time $\mathcal{N}_{\text{pre}}$ [70]. It is clear from Fig. 7 that $\rho_{(\phi)}^q$ cannot provide successful preheating for

any of the BPs, but preheating is possible for BP b via $\rho_{(h)}^q$. This is primarily due to the smaller ξ_H in BP a . We find that for BP b , $\rho_{(h)}^q \approx \rho_{\text{inf}}$ at $\mathcal{N} \approx 1.8$. As discussed above, we assume that thermalization is complete immediately after preheating is completed [68]. It is clear that preheating via the production of Higgs quanta is not possible for BP a ; however, there is a subtlety, which we shall return to shortly.

We are now equipped with all quantities needed to evaluate the right side of Eq. (3.29). The preheating temperature T_{pre} is evaluated equating ρ_{pre} to the energy density of the thermal bath

$$\rho_{\text{inf}}|_{\mathcal{N}=\mathcal{N}_{\text{pre}}} \equiv \rho_{\text{pre}} = \frac{g_{\text{pre}}\pi^2}{30} T_{\text{pre}}^4, \quad (4.22)$$

with all relevant quantities summarized in Table II. We find that $T_{\text{pre}} \approx 4.9 \times 10^{14}$ GeV for BP b . This leads to an e -folding value of 55.39 for the right-hand side of Eq. (3.29). In contrast, as discussed above, the left-hand side yields $54.51e$ -foldings. This simply means that the matching is very close but not exact. Indeed, one can slightly readjust the model parameters ξ_R , ξ_H , and ξ_c , along with $\varphi(t_{\text{in}})$ and $h_0(t_{\text{in}})$, to achieve exact matching; however, we refrain from such adjustments here. This is primarily because we have not considered the decay of the inflaton and Higgs condensates, as well as the production

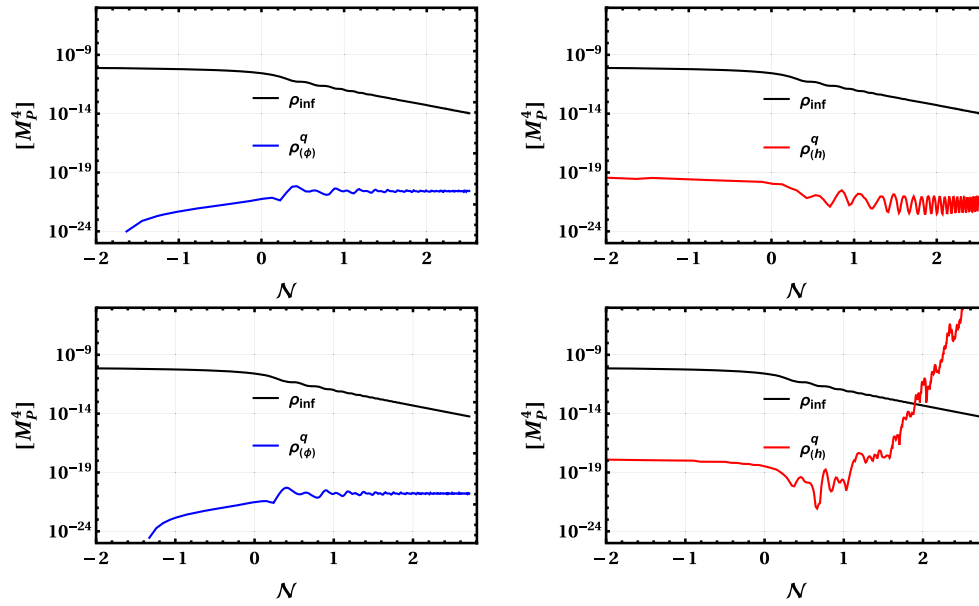


FIG. 7. The vacuum subtracted quantum energy densities $\rho_{(\phi)}^q$ and $\rho_{(h)}^q$ for BPa (upper panel) and BPb (lower panel) respectively along with ρ_{inf} for comparison.

of inflaton and Higgs quanta. Such decays will reduce ρ_{pre} , leading to a lower T_{pre} . We leave a detailed study for future work.

Although the preheating for BPa via ϕ and h particle production is incomplete, the associated longitudinal gauge bosons (or, conversely, Goldstone bosons) may complete preheating even for BPa [68]. It has been found that for pure R^2 -Higgs inflation (i.e., without the R^3 term), with similar parameter values and $\xi_H \sim 1$, the Goldstone bosons can preheat the Universe within $\mathcal{N} \simeq 3$ with a $T_{\text{pre}} \approx 5 \times 10^{14}$ GeV [68]. Inserting these numbers, we find that the left- and right-hand sides of Eq. (3.29) give 53.66 and 54.16, respectively. As discussed before, one can readjust the model parameters to achieve exact matching. A proper estimation of preheating for both BPa and BPb, including contributions from Goldstone and gauge bosons, would require decomposing the Higgs field Φ in the Coulomb gauge instead of the unitary gauge adopted here, as the latter becomes ill-defined during zero crossings of the Higgs condensate after the end of inflation [68,70]. This is deferred to a more detailed future publication.

TABLE II. The details of preheating for both the benchmark points chosen for our analysis.

BP	Preheating field(s)	$\mathcal{N}_{\text{pre}}$	$\rho_{\text{pre}}[M_P^4]$	$T_{\text{pre}} [\text{GeV}]$
a	...	...	...	...
b	h	1.9	5.9×10^{-14}	4.9×10^{14}

V. SUMMARY AND OUTLOOK

In this paper, we discussed how the addition of an R^3 term to R^2 -Higgs inflation may account for the observed high n_s in the latest CMB + BAO data within the Λ CDM model [13,14]. Considering two representative benchmark parameter sets, we found that $\xi_c \sim 10^{-13} - 10^{-14}$ is sufficient to explain the observed n_s , $\mathcal{A}_s$, and r in the single-fieldlike regime. We show that Higgs preheating plays a pivotal role in matching the reference CMB scale to the inflationary scale with $\xi_H \gtrsim 10$. Smaller values, $\xi_H \sim 1$, can indeed provide successful preheating, however, via the production of Goldstone bosons [68], which is not considered here. In addition, we have not considered the decay of the inflaton and Higgs condensate, nor the decay of the produced quanta and their backreaction on the background dynamics. This induces some uncertainties in our result, but the model parameters can be readjusted to achieve exact matching of scales easily. We also note that, in general, a negative value of ξ_c can account for the observed value of n_s , as discussed in Ref. [38]. In the case of R^2 -Higgs inflation, the presence of the Higgs boson introduces an additional field direction in the potential V_E . Consequently, a positive ξ_c can also reproduce the observed n_s if one departs from the exact single-field regime, as shown above. In this first attempt, we have simply assumed $\xi_c > 0$ and leave a more detailed analysis for future work.

We also identify further limitations here. It has been found that gauge boson production may play a vital role in R^2 -Higgs inflation for $\xi_H \gtrsim 1$ [68,70]. However, to consider both Goldstone and gauge preheating, one must

abandon the unitary gauge adopted here. This is primarily because the unitary gauge becomes ill-defined at each zero crossing of the Higgs condensate h_0 [68,70]. Moreover, the produced gauge bosons may decay leptonically and inhibit the completion of gauge preheating [68,70]. These factors introduce further uncertainties into our results. To the best of our knowledge, such effects have not been considered in the literature for R^2 -Higgs inflation following the ACT and SPT data.

At this point it is also useful to understand the role of BAO data within the Λ CDM model. Note that the BAO data does not directly constraint the n_s but, once they are combined with the CMB data correlated uncertainties shift the n_s to a higher value. Indeed, the CMB data from ACT, SPT, and Planck are consistent with each other and the combination of all CMB data i.e. the CMB-SPA measurement found $n_s = 0.9684 \pm 0.0030$ in the Λ CDM model. This is consistent with Planck 2018 within 1σ and pure Starobinsky, Higgs, and R^2 -Higgs models are still best fit models to the data. It has been shown that the higher n_s is due to the discrepancy between the CMB and BAO data and the discrepancy between BAO parameters and n_s in the CMB data within the Λ CDM model. While this apparent

tension between CMB and BAO may well be artifact of unknown systematics, it could well be an indication of new physics beyond Λ CDM [47]. Nonetheless, additional independent measurements of the CMB and BAO are required to resolve this tension. On the CMB side, the Simons Observatory is expected to achieve a precision of $\sigma(n_s) \sim 0.002$ [83]. BAO data from the Dark Energy Survey [84], and the improved sensitivity of Euclid in measuring BAO [85], may also contribute significantly. While we await more precise measurements, within the Λ CDM model, this high n_s may indicate the presence of an R^3 term in R^2 -Higgs inflation. If confirmed in the future data, along with detection of r , this may offer exquisite information about quantum gravity.

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DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

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- [1] A. A. Starobinsky, A new type of isotropic cosmological models without singularity, *Phys. Lett. B* **91**, 99 (1980).
 - [2] K. Sato, First order phase transition of a vacuum and expansion of the universe, *Mon. Not. R. Astron. Soc.* **195**, 467 (1981).
 - [3] A. H. Guth, The inflationary universe: A possible solution to the horizon and flatness problems, *Phys. Rev. D* **23**, 347 (1981).
 - [4] G. Hinshaw *et al.* (WMAP Collaboration), Nine-Year Wilkinson Microwave Anisotropy Probe (WMAP) Observations: Cosmological parameter results, *Astrophys. J. Suppl. Ser.* **208**, 19 (2013).
 - [5] Y. Akrami *et al.* (Planck Collaboration), Planck 2018 results. X. Constraints on inflation, *Astron. Astrophys.* **641**, A10 (2020).
 - [6] P. A. R. Ade *et al.* (BICEP, Keck Collaborations), Improved constraints on primordial gravitational waves using Planck, WMAP, and BICEP/Keck observations through the 2018 observing season, *Phys. Rev. Lett.* **127**, 151301 (2021).
 - [7] P. Ade *et al.* (Simons Observatory Collaboration), The Simons Observatory: Science goals and forecasts, *J. Cosmol. Astropart. Phys.* **02** (2018) 056.
 - [8] E. Allys *et al.* (LiteBIRD Collaboration), Probing cosmic inflation with the LiteBIRD cosmic microwave background polarization survey, *Prog. Theor. Exp. Phys.* **2023**, 042F01 (2023).
 - [9] A. A. Starobinsky, The perturbation spectrum evolving from a nonsingular initially de-Sitter cosmology and the microwave background anisotropy, *Sov. Astron. Lett.* **9**, 302 (1983).
 - [10] A. Vilenkin, Classical and quantum cosmology of the Starobinsky inflationary model, *Phys. Rev. D* **32**, 2511 (1985).
 - [11] M. B. Mijic, M. S. Morris, and W.-M. Suen, The R^2 cosmology: Inflation without a phase transition, *Phys. Rev. D* **34**, 2934 (1986).
 - [12] K.-i. Maeda, Inflation as a transient attractor in R^2 cosmology, *Phys. Rev. D* **37**, 858 (1988).
 - [13] T. Louis *et al.* (ACT Collaboration), The Atacama Cosmology Telescope: DR6 power spectra, likelihoods and Λ CDM parameters, [arXiv:2503.14452](https://arxiv.org/abs/2503.14452).
 - [14] E. Camphuis *et al.* (SPT-3G Collaboration), SPT-3G D1: CMB temperature and polarization power spectra and cosmology from 2019 and 2020 observations of the SPT-3G main field, [arXiv:2506.20707](https://arxiv.org/abs/2506.20707).
 - [15] M. Abdul Karim *et al.* (DESI Collaboration), DESI DR2 results II: Measurements of baryon acoustic oscillations and cosmological constraints, *Phys. Rev. D* **112**, 083515 (2025).
 - [16] A. Salvio and A. Mazumdar, Classical and quantum initial conditions for Higgs inflation, *Phys. Lett. B* **750**, 194 (2015).
 - [17] Y. Ena, Higgs scalaron mixed inflation, *Phys. Lett. B* **770**, 403 (2017).
 - [18] S. Pi, Y.-I. Zhang, Q.-G. Huang, and M. Sasaki, Scalaron from R^2 -gravity as a heavy field, *J. Cosmol. Astropart. Phys.* **05** (2017) 042.

- [19] D. Gorbunov and A. Tokareva, Scalaron the healer: Removing the strong-coupling in the Higgs- and Higgs-dilaton inflations, *Phys. Lett. B* **788**, 37 (2019).
- [20] A. Gundhi and C. F. Steinwachs, Scalaron-Higgs inflation, *Nucl. Phys. B* **954**, 114989 (2020).
- [21] M. He, R. Jinno, K. Kamada, S. C. Park, A. A. Starobinsky, and J. Yokoyama, On the violent preheating in the mixed Higgs- R^2 inflationary model, *Phys. Lett. B* **791**, 36 (2019).
- [22] M. He, A. A. Starobinsky, and J. Yokoyama, Inflation in the mixed Higgs- R^2 model, *J. Cosmol. Astropart. Phys.* **05** (2018) 064.
- [23] D. Y. Cheong, S. M. Lee, and S. C. Park, Primordial black holes in Higgs- R^2 inflation as the whole of dark matter, *J. Cosmol. Astropart. Phys.* **01** (2019) 032.
- [24] D. D. Canko, I. D. Gialamas, and G. P. Kodaxis, A simple $F(\mathcal{R}, \phi)$ deformation of Starobinsky inflationary model, *Eur. Phys. J. C* **80**, 458 (2020).
- [25] M. He, R. Jinno, K. Kamada, A. A. Starobinsky, and J. Yokoyama, Occurrence of tachyonic preheating in the mixed Higgs- R^2 model, *J. Cosmol. Astropart. Phys.* **01** (2020) 066.
- [26] M. He, Perturbative reheating in the mixed Higgs- R^2 model, *J. Cosmol. Astropart. Phys.* **05** (2020) 021.
- [27] F. Pineda and L. O. Pimentel, Non-perturbative approach for scalar particle production in Higgs- R^2 inflation, [arXiv:2409.09999](#).
- [28] E. J. Copeland, C. Rahmede, and I. D. Saltas, Asymptotically safe Starobinsky inflation, *Phys. Rev. D* **91**, 103530 (2015).
- [29] I. D. Saltas, Higgs inflation and quantum gravity: An exact renormalisation group approach, *J. Cosmol. Astropart. Phys.* **02** (2015) 048.
- [30] A. de la Cruz-Dombriz, E. Elizalde, S. D. Odintsov, and D. Sáez-Gómez, Spotting deviations from R^2 inflation, *J. Cosmol. Astropart. Phys.* **05** (2016) 060.
- [31] S. D. Odintsov, V. K. Oikonomou, and F. P. Fronimos, Canonical scalar field inflation with string and R^2 -corrections, *Ann. Phys. (Amsterdam)* **424**, 168359 (2021).
- [32] I. D. Gialamas, A. Karam, A. Racioppi, and M. Raidal, Has ACT measured radiative corrections to the tree-level Higgs-like inflation?, [arXiv:2504.06002](#).
- [33] M. Drees and Y. Xu, Refined predictions for Starobinsky inflation and post-inflationary constraints in light of ACT, *Phys. Lett. B* **867**, 139612 (2025).
- [34] D. S. Zharov, O. O. Sobol, and S. I. Vilchinskii, Reheating acts on Starobinsky and Higgs inflation, [arXiv:2505.01129](#).
- [35] L. Liu, Z. Yi, and Y. Gong, Reconciling Higgs inflation with ACT observations through reheating, [arXiv:2505.02407](#).
- [36] M. R. Haque, S. Pal, and D. Paul, Improved predictions on Higgs-Starobinsky inflation and reheating with ACT DR6 and primordial gravitational waves, *Phys. Lett. B* **869**, 139852 (2025).
- [37] Yogesh, A. Mohammadi, Q. Wu, and T. Zhu, Starobinsky like inflation and EGB gravity in the light of ACT, *J. Cosmol. Astropart. Phys.* **10** (2025) 010.
- [38] A. Addazi, Y. Aldabergenov, and S. V. Ketov, Curvature corrections to Starobinsky inflation can explain the ACT results, *Phys. Lett. B* **869**, 139883 (2025).
- [39] R. Kallosh and A. Linde, On the present status of inflationary cosmology, *Gen. Relativ. Gravit.* **57**, 135 (2025).
- [40] I. D. Gialamas, T. Katsoulas, and K. Tamvakis, Keeping the relation between the Starobinsky model and no-scale supergravity active, *J. Cosmol. Astropart. Phys.* **09** (2025) 060.
- [41] S. Saini and A. Nautiyal, Power law α -Starobinsky inflation, [arXiv:2505.16853](#).
- [42] W. J. Wolf, Inflationary attractors and radiative corrections in light of ACT data, [arXiv:2506.12436](#).
- [43] X. Wang, K. Kohri, and T. T. Yanagida, Primordial black holes save R^2 inflation, [arXiv:2506.06797](#).
- [44] M. Piva, Classification of $f(R)$ theories of inflation and the uniqueness of Starobinsky model, [arXiv:2507.02637](#).
- [45] N. Sidik Risdianto, R. H. S. Budhi, N. Shobcha, and A. Salim Adam, The preheating stage on the Starobinsky inflation after ACT, [arXiv:2507.12868](#).
- [46] D. S. Zharov, O. O. Sobol, and S. I. Vilchinskii, ACT observations, reheating, and Starobinsky and Higgs inflation, *Phys. Rev. D* **112**, 023544 (2025).
- [47] E. G. M. Ferreira, E. McDonough, L. Balkenhol, R. Kallosh, L. Knox, and A. Linde, The BAO-CMB tension and implications for inflation, [arXiv:2507.12459](#).
- [48] J. Ellis, M. A. G. Garcia, N. Nagata, D. V. Nanopoulos, and K. A. Olive, Deformations of Starobinsky inflation in no-scale SU(5) and SO(10) GUTs, [arXiv:2508.13279](#).
- [49] V. K. Oikonomou, Strong gravity effects on $\mathcal{R}^2$ -corrected single scalar field inflation and compatibility with the ACT data, *Phys. Lett. B* **871**, 139972 (2025).
- [50] T. Modak, L. Röver, B. M. Schäfer, B. Schosser, and T. Plehn, Cornering extended Starobinsky inflation with CMB and SKA, *SciPost Phys.* **15**, 047 (2023).
- [51] S. M. Lee, T. Modak, K.-y. Oda, and T. Takahashi, Ultra-violet sensitivity in Higgs-Starobinsky inflation, *J. Cosmol. Astropart. Phys.* **08** (2023) 045.
- [52] T. Saidov and A. Zhuk, Bouncing inflation in nonlinear $R^2 + R^4$ gravitational model, *Phys. Rev. D* **81**, 124002 (2010).
- [53] Q.-G. Huang, A polynomial $f(R)$ inflation model, *J. Cosmol. Astropart. Phys.* **02** (2013) 035.
- [54] H. Motohashi, Consistency relation for R^p inflation, *Phys. Rev. D* **91**, 064016 (2015).
- [55] T. Asaka, S. Iso, H. Kawai, K. Kohri, T. Noumi, and T. Terada, Reinterpretation of the Starobinsky model, *Prog. Theor. Exp. Phys.* **2016**, 123E01 (2016).
- [56] K. Bamba and S. D. Odintsov, Inflationary cosmology in modified gravity theories, *Symmetry* **7**, 220 (2015).
- [57] T. Miranda, J. C. Fabris, and O. F. Piattella, Reconstructing a $f(R)$ theory from the α -attractors, *J. Cosmol. Astropart. Phys.* **09** (2017) 041.
- [58] D. Y. Cheong, H. M. Lee, and S. C. Park, Beyond the Starobinsky model for inflation, *Phys. Lett. B* **805**, 135453 (2020).
- [59] G. Rodrigues-da Silva, J. Bezerra-Sobrinho, and L. G. Medeiros, Higher-order extension of Starobinsky inflation: Initial conditions, slow-roll regime, and reheating phase, *Phys. Rev. D* **105**, 063504 (2022).
- [60] V. R. Ivanov, S. V. Ketov, E. O. Pozdeeva, and S. Y. Vernov, Analytic extensions of Starobinsky model of inflation, *J. Cosmol. Astropart. Phys.* **03** (2021) 058.
- [61] A. S. Koshelev, K. S. Kumar, and A. A. Starobinsky, Generalized non-local R^2 -like inflation, *J. High Energy Phys.* **07** (2022) 146.

- [62] Y. Shtanov, V. Sahni, and S. S. Mishra, Tabletop potentials for inflation from $f(R)$ gravity, *J. Cosmol. Astropart. Phys.* **03** (2022) 023.
- [63] Q.-Y. Wang, Y. Tang, and Y.-L. Wu, Inflation in Weyl scaling invariant gravity with R^3 extensions, *Phys. Rev. D* **107**, 083511 (2023).
- [64] J. Kim, X. Wang, Y.-I. Zhang, and Z. Ren, Enhancement of primordial curvature perturbations in R^3 -corrected Starobinsky-Higgs inflation, *J. Cosmol. Astropart. Phys.* **09** (2025) 011.
- [65] L. Sebastiani, G. Cognola, R. Myrzakulov, S. D. Odintsov, and S. Zerbini, Nearly Starobinsky inflation from modified gravity, *Phys. Rev. D* **89**, 023518 (2014).
- [66] S. D. Odintsov, V. K. Oikonomou, and L. Sebastiani, Unification of constant-roll inflation and dark energy with logarithmic R^2 -corrected and exponential $F(R)$ gravity, *Nucl. Phys. B* **923**, 608 (2017).
- [67] S. Nojiri, S. D. Odintsov, and V. K. Oikonomou, Modified gravity theories on a nutshell: Inflation, bounce and late-time evolution, *Phys. Rep.* **692**, 1 (2017).
- [68] Y. Cado, C. Englert, T. Modak, and M. Quirós, Implication of preheating on gravity assisted baryogenesis in R^2 -Higgs inflation, *Phys. Rev. D* **111**, 023042 (2025).
- [69] J.-O. Gong and T. Tanaka, A covariant approach to general field space metric in multi-field inflation, *J. Cosmol. Astropart. Phys.* **03** (2011) 015; *J. Cosmol. Astropart. Phys.* **02** (2012) E01.
- [70] E. I. Sfakianakis and J. van de Vis, Preheating after Higgs inflation: Self-resonance and gauge boson production, *Phys. Rev. D* **99**, 083519 (2019).
- [71] D. I. Kaiser, E. A. Mazenc, and E. I. Sfakianakis, Primordial bispectrum from multifield inflation with nonminimal couplings, *Phys. Rev. D* **87**, 064004 (2013).
- [72] H. Kodama and M. Sasaki, Cosmological perturbation theory, *Prog. Theor. Phys. Suppl.* **78**, 1 (1984).
- [73] V. F. Mukhanov, H. A. Feldman, and R. H. Brandenberger, Theory of cosmological perturbations. Part 1. Classical perturbations. Part 2. Quantum theory of perturbations. Part 3. Extensions, *Phys. Rep.* **215**, 203 (1992).
- [74] K. A. Malik and D. Wands, Cosmological perturbations, *Phys. Rep.* **475**, 1 (2009).
- [75] J. Elliston, D. Seery, and R. Tavakol, The inflationary bispectrum with curved field-space, *J. Cosmol. Astropart. Phys.* **11** (2012) 060.
- [76] M. Sasaki, Large scale quantum fluctuations in the inflationary universe, *Prog. Theor. Phys.* **76**, 1036 (1986).
- [77] V. F. Mukhanov, Theory of cosmological perturbations. Part 1. Classical perturbations. Part 2. Quantum theory of perturbations. Part 3. Extensions, *Phys. Rept.* **215**, 203 (1992).
- [78] B. A. Bassett, S. Tsujikawa, and D. Wands, Inflation dynamics and reheating, *Rev. Mod. Phys.* **78**, 537 (2006).
- [79] S. Antusch, D. Nolde, and S. Orani, Hill crossing during preheating after Hilltop inflation, *J. Cosmol. Astropart. Phys.* **06** (2015) 009.
- [80] M. P. DeCross, D. I. Kaiser, A. Prabhu, C. Prescod-Weinstein, and E. I. Sfakianakis, Preheating after multifield inflation with nonminimal couplings, I: Covariant formalism and attractor behavior, *Phys. Rev. D* **97**, 023526 (2018).
- [81] M. P. DeCross, D. I. Kaiser, A. Prabhu, C. Prescod-Weinstein, and E. I. Sfakianakis, Preheating after multifield inflation with nonminimal couplings, III: Dynamical space-time results, *Phys. Rev. D* **97**, 023528 (2018).
- [82] Y. Cado, C. Englert, T. Modak, and M. Quirós, Baryogenesis in R^2 -Higgs inflation: The gravitational connection, *Phys. Rev. D* **109**, 043026 (2024).
- [83] M. Abitbol *et al.* (Simons Observatory Collaboration), The Simons Observatory: Science goals and forecasts for the enhanced large aperture telescope, *J. Cosmol. Astropart. Phys.* **08** (2025) 034.
- [84] T. M. C. Abbott *et al.* (DES Collaboration), Dark Energy Survey: A 2.1% measurement of the angular baryonic acoustic oscillation scale at redshift $z_{\text{eff}} = 0.85$ from the final dataset, *Phys. Rev. D* **110**, 063515 (2024).
- [85] V. Duret *et al.* (Euclid Collaboration), Euclid preparation. BAO analysis of photometric galaxy clustering in configuration space, [arXiv:2503.11621](https://arxiv.org/abs/2503.11621).