

Shifted hybrid realization of nonminimal Higgs inflation in light of ACT DR6 and Planck data

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We investigate nonminimal Higgs inflation in a no-scale-inspired supergravity framework and confront its predictions with the latest cosmic microwave background (CMB) constraints from ACT DR6 and *Planck*. Working within a shifted hybrid inflation scenario, we construct an effective single-field description in which the grand unified theory Higgs direction serves as the inflaton after stabilization of the orthogonal scalar fields. We show that the inclusion of the leading nonrenormalizable operator in the superpotential induces a controlled deformation of the Starobinsky attractor, allowing the scalar spectral index to be shifted into the range favored by recent ACT and related CMB datasets while maintaining a small tensor-to-scalar ratio, $r \sim 10^{-3}$ – 10^{-2} . The resulting inflationary dynamics remain theoretically consistent, with controlled supergravity corrections and sub-Planckian inflaton field values. We perform a detailed numerical analysis of the model parameter space, including reheating and nonthermal leptogenesis, and identify regions that simultaneously satisfy current observational constraints and yield a viable postinflationary cosmological history.

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I. INTRODUCTION

The inflationary paradigm has emerged as the foundation of modern cosmology, elegantly addressing the horizon and flatness problems while providing a mechanism for the origin of large-scale structure [1–7]. By positing a period of rapid exponential expansion in the early Universe, inflation not only explains the observed homogeneity and isotropy on cosmic scales but also provides a quantum mechanical mechanism for generating primordial density perturbations, which serve as seeds for galaxy formation and the cosmic

microwave background (CMB) anisotropies. Single-field inflation is still the most researched and predictive framework among the many proposed realizations. Its attractiveness stems from its few assumptions and the fact that it yields distinct, verifiable predictions for the scalar spectral index (n_s) and the tensor-to-scalar ratio (r), which can be directly compared with accurate measurements of the CMB.

Advancements in CMB measurements have continually refined these predictions, imposing increasingly rigorous constraints on inflationary theories. The Planck 2018 results [8,9] results set benchmark values, with a $n_s = 0.9649 \pm 0.0042$ (68% CL) and an upper limit on $r < 0.056$ (95% CL) when paired with BICEP2/Keck Array data [10]. More recently, the Atacama Cosmology Telescope (ACT) Data Release 6 (DR6) has provided high-resolution CMB measurements that lead to a mildly higher preferred value of the scalar spectral index when combined with Planck data, $n_s = 0.9709 \pm 0.0038$ [11]. This upward shift becomes slightly more pronounced when the analysis is supplemented with DESI Year-1 BAO measurements,

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yielding $n_s = 0.9743 \pm 0.0034$ [11,12]. These updates are important for differentiating between inflationary models, as they investigate the form of the inflationary potential and the dynamics of the early universe with remarkable accuracy.

However, recent CMB measurements sharpen the preferred range of the scalar tilt and reveal mild tensions across data combinations. The SPT-3G D1 analysis finds that a CMB-only combination of Planck, ACT, and SPT, denoted the “CMB-SPA” likelihood, gives $n_s = 0.9679 \pm 0.0033$ [13], close to the original Planck constraint [8,9]. When DESI DR2 BAO data are added to this combination, the preferred value shifts upward to $n_s = 0.9726 \pm 0.0028$ [13,14]. This is consistent with ACT-based combinations, which also favor a slightly higher tilt when combined with DESI BAO [11,14]. Importantly, the SPT-3G D1 analysis reports a $\sim 2.8\sigma$ difference between CMB-SPA and DESI DR2 BAO within Λ CDM [13,14]. Therefore, constraints obtained by combining these data sets should be interpreted with some caution (see e.g., Ref. [15]). In what follows, we compare our predictions with both the conservative CMB-only band near $n_s \simeq 0.968$ and the DESI-informed band near $n_s \simeq 0.973$ [8,9,11,13,14].

On the theoretical side, plateau or attractor models provide some of the most successful and predictive realizations of inflation. Prominent examples include the $\mathcal{R}^2$ (Starobinsky) model [16] and its superconformal and no-scale supergravity extensions, which generically yield

$$n_s \simeq 1 - \frac{2}{N} + \mathcal{O}\left(\frac{1}{N^2}\right), \quad r \simeq \mathcal{O}\left(\frac{1}{N^2}\right). \quad (1)$$

For a typical number of e-folds $N \simeq 55$, these models predict $n_s \simeq 0.964$ and $r \simeq 4 \times 10^{-3}$ [17,18]. The same attractor behavior also emerges in nonminimal Higgs inflation and in a broad class of α -attractor models, making these predictions remarkably robust across diverse theoretical frameworks. While this universality is one of the major theoretical successes of attractor inflation, it also motivates the exploration of controlled departures from the strict attractor limit. Such deformations can arise from higher-dimensional operators, corrections to the Kähler potential, or small departures from exact no-scale structure in supergravity embeddings. Several supersymmetric realizations of Starobinsky-like inflation [19–26] have recently been proposed to reconcile attractor-based inflationary predictions with the latest observational data. For representative nonsupersymmetric realizations aimed at accommodating recent observational results, see [27–36]. A particularly attractive feature of these modifications is that they can increase the scalar spectral index while preserving the characteristic prediction of a small tensor-to-scalar ratio, typically in the range $r \sim \text{few} \times 10^{-3} - 10^{-2}$. Interestingly, this is precisely the regime that forthcoming CMB polarization experiments such as LiteBIRD and CMB-S4 are

expected to probe with unprecedented sensitivity [37,38]. These considerations provide a strong motivation for studying well-controlled deformations of attractor inflation within realistic particle-physics frameworks.

Supersymmetric hybrid inflation [39,40] provides one of the most compelling realizations of inflation within particle physics and grand unified theories [41]. In its standard version with minimal canonical Kähler potential, inflation is driven by a gauge-singlet superfield S , whose nearly flat potential is generated by supersymmetry and lifted by radiative, supergravity corrections and soft-supersymmetric breaking terms [42–44]. More recently, smooth hybrid inflation has been shown to accommodate the larger values of the scalar spectral index preferred by ACT- and SPT-informed datasets [45,46]. Depending on the structure of the Kähler potential, the resulting inflationary predictions can vary significantly [47,48]. In particular, logarithmic Kähler potentials inspired by no-scale supergravity naturally generate nonminimal couplings between the Higgs sector and gravity, thereby realizing nonminimal Higgs inflation within the supersymmetric hybrid inflation framework [49–52]. In the large- ξ limit, the Einstein-frame potential approaches the Starobinsky attractor form, yielding the characteristic predictions displayed in Eq. (1). While these predictions remain in excellent agreement with the original *Planck* results, the recent preference of ACT- and DESI-combined datasets for a slightly larger scalar spectral index motivates the exploration of well-motivated deviations from the attractor predictions.

In this paper, we investigate such a departure within the shifted realization of supersymmetric hybrid inflation. Unlike the standard hybrid inflation scenario, where the gauge symmetry remains unbroken during inflation and is subsequently broken at the waterfall transition, shifted hybrid inflation incorporates higher-dimensional superpotential operators that displace the inflationary trajectory away from the origin, allowing the gauge symmetry to be broken already during inflation [53–55]. This modification not only avoids the production of topological defects at the end of inflation but also introduces additional flexibility in the inflationary dynamics. By embedding the shifted hybrid inflation framework into nonminimal Higgs inflation with a logarithmic no-scale Kähler potential, we show that the resulting deformation of the Starobinsky attractor can naturally raise the scalar spectral index into the range favored by ACT DR6 and related data combinations while maintaining a tensor-to-scalar ratio in the experimentally accessible range $r \sim 10^{-3} - 10^{-2}$. We perform a comprehensive numerical analysis of the full Einstein-frame dynamics, including reheating and leptogenesis, and identify the regions of parameter space that simultaneously satisfy current observational constraints and provide a realistic cosmological history.

The remainder of this paper is organized as follows. In Sec. II, we introduce the model and its theoretical framework. In Sec. III, we develop its effective single-field

realization and derive analytical expressions for the inflationary observables in the slow-roll approximation. In Sec. IV A, we present the numerical analysis and discuss the regions of parameter space consistent with current observational constraints. Reheating and leptogenesis are studied in Sec. VI. Finally, our conclusions are summarized in Sec. VII.

II. R-SYMMETRIC HIGGS INFLATION

In order to investigate nonminimal Higgs inflation, we focus on a specific and phenomenologically realistic example built on the Pati-Salam gauge group $4_c \times 2_L \times 2_R \equiv \text{SU}(4)_c \times \text{SU}(2)_L \times \text{SU}(2)_R$ [56–58]. In supersymmetric constructions [54,59,60], shifted μ -hybrid inflation is realized through the following superpotential:

$$W = \kappa S(\bar{H}^c H^c - M^2) + \beta S \frac{(\bar{H}^c H^c)^2}{m_P^2} + \lambda S \mathfrak{h}^2 + a G H^c H^c + b G \bar{H}^c \bar{H}^c + \lambda_{ij} F_i^c F_j \mathfrak{h} + (\gamma_{ij}^1 F_i^c F_j^c + \gamma_{ij}^2 F_i F_j) \frac{H^c H^c}{m_P} + (\tilde{\gamma}_{ij}^1 F_i^c F_j^c + \tilde{\gamma}_{ij}^2 F_i F_j) \frac{\bar{H}^c \bar{H}^c}{m_P}, \quad (2)$$

where $\kappa, \lambda, \beta, a, b, \lambda_{ij}, \gamma_{ij}^1, \gamma_{ij}^2, \tilde{\gamma}_{ij}^1, \tilde{\gamma}_{ij}^2$ are real and positive dimensionless couplings, M denotes a superheavy mass scale, and, for simplicity, the cutoff scale is taken to be the reduced Planck mass, $m_P = 2.4 \times 10^{18}$ GeV. The matter and Higgs superfields of the supersymmetric (SUSY) Pati-Salam model, together with their representations, transformation properties under $4_c \times 2_L \times 2_R$, decompositions under Minimal Supersymmetric Standard Model (MSSM) gauge group $3_c \times 2_L \times 1_Y$, and $U(1)_R$ charge assignments, are summarized in Table I.

The third term, $\lambda S \mathfrak{h}^2$, yields the effective μ term. Assuming gravity-mediated SUSY breaking [61], the scalar component of S acquires a nonzero vacuum expectation value (VEV) proportional to the gravitino mass $m_{3/2}$. This induces an effective μ term, $\mu = -\lambda m_{3/2}/\kappa$, thereby resolving the MSSM μ problem [62].

The two superpotential terms involving the sextuplet superfield (G) are introduced to generate superheavy masses for d_H^c and $\bar{d}_H^c$. The λ_{ij} terms correspond to the Yukawa couplings and thus generate fermion masses, while the γ_{ij} terms give rise to large right-handed neutrino masses required for the see-saw mechanism.

The first three terms in the superpotential (W) are relevant for nonminimal Higgs inflation,

$$W \supset \kappa S(\bar{H}^c H^c - M^2) + \beta S \frac{(\bar{H}^c H^c)^2}{m_P^2}, \quad (3)$$

for which the global SUSY vacuum is given by

$$\langle S \rangle = 0, \quad \langle \mathfrak{h} \rangle = 0, \quad v^2 \equiv \langle H^c \bar{H}^c \rangle. \quad (4)$$

TABLE I. Superfields together with their decomposition under the SM and their R charge.

Superfields	$4_c \times 2_L \times 2_R$	$3_c \times 2_L \times 1_Y$	$q(R)$
F_i	$(4, 2, 1)$	$Q_{ia}(3, 2, \frac{1}{6})$ $L_i(1, 2, -\frac{1}{2})$	1
F_i^c	$(\bar{4}, 1, 2)$	$u_{ia}^c(\bar{3}, 1, -\frac{2}{3})$ $d_{ia}^c(\bar{3}, 1, \frac{1}{3})$ $\nu_i^c(1, 1, 0)$ $e_i^c(1, 1, 1)$	1
H^c	$(\bar{4}, 1, 2)$	$u_{Ha}^c(\bar{3}, 1, -\frac{2}{3})$ $d_{Ha}^c(\bar{3}, 1, \frac{1}{3})$ $\nu_H^c(1, 1, 0)$ $e_H^c(1, 1, 1)$	0
$\bar{H}^c$	$(4, 1, 2)$	$\bar{u}_{Ha}^c(3, 1, \frac{2}{3})$ $\bar{d}_{Ha}^c(3, 1, -\frac{1}{3})$ $\bar{\nu}_H^c(1, 1, 0)$ $\bar{e}_H^c(1, 1, -1)$	0
S	$(1, 1, 1)$	$S(1, 1, 0)$	2
G	$(6, 1, 1)$	$g_a(3, 1, -\frac{1}{3})$ $g^a(\bar{3}, 1, \frac{1}{3})$	2
$\mathfrak{h}$	$(1, 2, 2)$	$\mathfrak{h}_u(1, 2, \frac{1}{2})$ $\mathfrak{h}_d(1, 2, -\frac{1}{2})$	0

We note that the nonrenormalizable term proportional to β can also appear with other possible structures, such as $(H^c)^4$ and $(\bar{H}^c)^4$. However, these additional contributions vanish along the D -flat right-handed neutrino direction, $(H^c, \bar{H}^c) \supset (\nu_H^c, \nu_H)$, with $|\nu_H^c| = |\nu_H| = h/2$, which is relevant for the subsequent discussion.

The global SUSY VEV v of the Higgs fields can be expressed as

$$(v_{\pm}/M)^2 = \frac{-1 \pm \sqrt{1 + 4\xi_{\text{sh}}}}{2\xi_{\text{sh}}}, \quad (5)$$

where the shifted hybrid inflation parameter ξ_{sh} is defined as

$$\xi_{\text{sh}} \equiv \frac{\beta_{\kappa} M^2}{m_P^2}, \quad \beta_{\kappa} \equiv \frac{\beta}{\kappa}. \quad (6)$$

The dependence of the normalized Higgs VEV v/M on ξ_{sh} is shown in Fig. 1, where the solid and dashed curves correspond to the two analytic branches of the solution. The requirement that v be real imposes the constraint $\xi_{\text{sh}} \geq -1/4$, or equivalently,

$$\beta_{\kappa} \geq -\frac{m_P^2}{4M^2}. \quad (7)$$

This form of the superpotential has previously been employed in shifted μ -hybrid inflation models

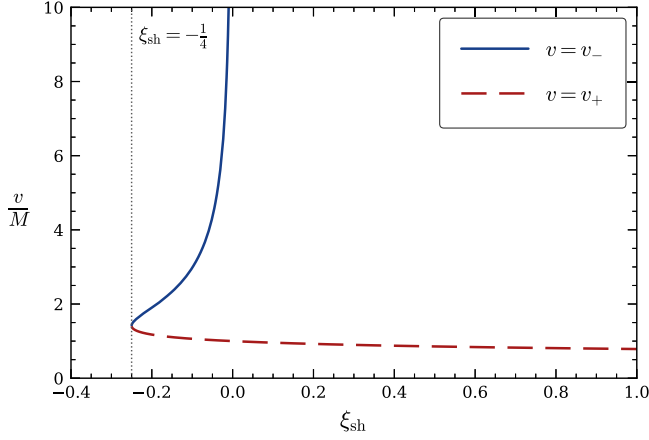


FIG. 1. Normalized Higgs VEV vs ξ_{sh} parameter. The ratio v/M as a function of the parameter ξ_{sh} . The solid (blue) and dashed (red) curves show the two analytic branches of v/M .

[43,54,59,63–65], typically with $\beta < 0$. In those scenarios, inflation is driven by the singlet field S , while the Higgs fields are stabilized at a local minimum along the D -flat direction with $|\nu_H^c| = |\nu_H|$ [52]. The parameter range $-1/4 < \xi_{\text{sh}} \leq -1/8$ is usually considered, corresponding to two distinct solutions for v . Beyond its inflationary relevance, the present framework is also connected to other testable aspects of high-scale particle physics. For example, closely related realizations have been shown to predict proton-decay rates that may be accessible to future experiments [66].

In contrast, in nonminimal Higgs inflation the roles of the fields are reversed; inflation is driven by the grand unified theory Higgs field h , while all other fields, including S , are stabilized at the origin. As will be discussed later, we instead focus on the regime $\xi_{\text{sh}} > 0$, which allows compatibility with recent CMB observations. For $\beta_\kappa \lesssim 1$ and $M \sim 10^{-2} m_P$, we find $0 < \xi_{\text{sh}} \ll 1$, leading to $v \lesssim M$. In the limit where the nonrenormalizable term is neglected ($\beta = 0$), one obtains $v = M$ [52]; however, the resulting predictions for CMB observables ($n_s \lesssim 0.965$ with $N_0 \lesssim 55$) in this limit are in tension with current data.

To realize nonminimal Higgs inflation we consider the Kähler potential K of the following form:

$$K = -3m_P^2 \ln \left(1 - \frac{1}{3m_P^2} (|S|^2 + |H^c|^2 + |\overline{H^c}|^2) \right) + \frac{\chi}{4m_P^2} (H^c \overline{H^c} + \text{H.c.}) + \frac{\gamma_4}{3m_P^4} |S|^4, \quad (8)$$

where all other fields are assumed to be stabilized at the origin. Here, χ and γ_4 are dimensionless parameters. The quartic term proportional to γ_4 is introduced to ensure the stabilization of the S field at the origin during Higgs inflation [67,68]. An alternative approach to stabilizing the S field, without invoking the quartic term in the Kähler

potential, is to consider a Kähler manifold of the form $U(1)_R \times (SU(2)/U(1))_S$ [69]. It is also worth emphasizing that all terms in the Kähler potential preserve the underlying R -symmetry. This differs from scenarios in which R -symmetry is explicitly broken at the nonrenormalizable level in the Kähler potential [70], for example, to facilitate the production of primordial black holes [71].

The scalar-gravity portion of the Lagrangian is expressed in the following manner:

$$\mathcal{L}_J = \sqrt{-g_J} \left[\frac{m_P^2}{2} \Omega^2 \mathcal{R}_J - g_J^{\mu\nu} \mathcal{G}_{ij} \partial_\mu z^i \partial_\nu z^{*j} - V_J \right], \quad (9)$$

where $\Omega^2 = e^{-K/3m_P^2}$, $\mathcal{R}_J$ is the Ricci scalar in the Jordan frame, $g_J^{\mu\nu}$ is the inverse of the Jordan frame spacetime metric $g_{J\mu\nu}$, $\mathcal{G}_{ij}$ is the field-space metric in the Jordan frame, V_J represents the scalar potential in the Jordan frame, and $z_i \in \{H, \overline{H}, S\}$. We then proceed with a conformal transformation of the spacetime metric, to work in the Einstein frame,

$$g_{\mu\nu}^E = \Omega^2 g_{\mu\nu}^J. \quad (10)$$

This transformation leads to the following Lagrangian in the Einstein frame with a canonical coupling to gravity:

$$\mathcal{L}_E = \sqrt{-g_E} \left[\frac{m_P^2}{2} \mathcal{R}_E - \frac{1}{2} G_{ij} g_E^{\mu\nu} \partial_\mu z^i \partial_\nu z^{*j} - V_E \right], \quad (11)$$

where G_{ij} is the field space metric in Einstein frame and $V_E = \Omega^{-4} V_J$ is the Einstein frame potential. Along the D -flat direction, the supergravity scalar potential in the Einstein frame, V_E , is defined as

$$V_E = e^{K/m_P^2} ((K^{-1})_{ij} D_{z_i} W D_{z_j}^* W^* - 3m_P^{-2} |W|^2), \quad (12)$$

where

$$D_{z_i} W \equiv \frac{\partial W}{\partial z_i} + m_P^{-2} \frac{\partial K}{\partial z_i} W, \quad K_{ij} \equiv \frac{\partial^2 K}{\partial z_i \partial z_j^*}, \quad (13)$$

$D_{z_i}^* W^* = (D_{z_i} W)^*$ and K given in (8). For the background equations and the relevant power spectra in a multifield framework, see Refs. [72–77].

III. MULTIFIELD ANALYSIS

We can write the complex fields (S , ν_H , ν_H^c) in terms of the canonically normalized real scalar fields (s , h) as

$$S = \frac{1}{\sqrt{2}} s e^{i\theta_s}, \quad \nu_H^c = \frac{h}{\sqrt{2}} e^{i\theta_h} \cos \varphi_h, \quad \overline{\nu_H^c} = \frac{h}{\sqrt{2}} e^{i\theta_h} \sin \varphi_h, \quad (14)$$

where θ_s , θ_h , and $\bar{\theta}_h$ are the respective phases. The D-Flat direction, $\nu_H^c = (\nu_H^c)^*$, implies the following conditions,

$$\varphi_h = \frac{\pi}{4}, \quad \theta_h = \bar{\theta}_h. \quad (15)$$

A detailed discussion of the stabilization of the fields s , θ_s , and θ_h is in order. For simplicity, we neglect contributions arising from the β coupling. Since inflation is expected to originate at very high energy scales, it is natural to consider Planckian initial conditions. Accordingly, we take the initial values of the scalar fields to satisfy $s \gtrsim m_P \gg M$ and $h \gtrsim m_P$. In the regime where $\chi h^2 \gg m_P^2$, the Einstein-frame scalar potential reduces to [78]

$$V_E \simeq \frac{4\kappa^2 m_P^4}{\chi^2 \cos^2(2\theta_h)} \left[1 - \frac{8}{\cos(2\theta_h)} \frac{m_P^2}{\chi h^2} + 2\gamma_4 \frac{s^2}{m_P^2} + \frac{16}{3} \frac{s^2}{\chi h^2} \left(\frac{1}{\cos(2\theta_h)} - \cos(2\theta_h) \right) + \frac{4}{3} \frac{s^2}{\chi h^2} (-\cos(2\theta_s) + 3 \tan(2\theta_h) \sin(2\theta_s)) \right]. \quad (16)$$

From this expression, it follows that the potential is minimized at $s = 0$, $\theta_s = 0$, and $\theta_h = 0$, indicating that these field directions are stabilized along the inflationary trajectory.

The exact Einstein-frame potential, expressed as a function of s and h , for $\beta_\kappa = 0$ and with all phases stabilized, takes the form

$$V_E(h, s) = \frac{\kappa^2}{16\Omega^4\Delta} \mathcal{F} \quad (17)$$

where

$$\begin{aligned} \Omega^2 &= 1 + \frac{\xi h^2}{m_P^2} - \frac{1}{6} \frac{s^2}{m_P^2} + \frac{\gamma_4}{12} \frac{s^4}{m_P^4}, \\ \Delta &= 1 + \xi(1 + 6\xi) \frac{h^2}{m_P^2} - 2\gamma_4 \frac{s^2}{m_P^2} + \frac{\gamma_4}{12} \frac{s^4}{m_P^4} \\ &\quad - 2\xi(1 + 6\xi)\gamma_4 \frac{h^2 s^2}{m_P^4}, \\ \mathcal{F} &= (h^2 - 4M^2)^2 \left[1 + \xi(1 + 6\xi) \frac{h^2}{m_P^2} \right] \\ &\quad - \frac{2}{3} [16M^4 - 2(3m_P^2 + 4(1 + 6\xi)M^2)h^2 \\ &\quad + (1 + 6\xi)h^4] \frac{s^2}{m_P^2} + \frac{\gamma_4}{12} [400M^4 \\ &\quad - 8(12m_P^2 + 5(5 + 24\xi)M^2)h^2 \\ &\quad + (25 + 144\xi)h^4] \frac{s^4}{m_P^4} + \frac{\gamma_4}{3} \frac{h^2 s^6}{m_P^6}. \end{aligned} \quad (18)$$

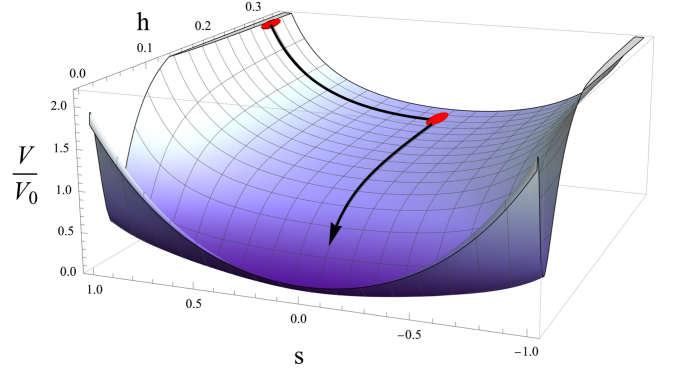


FIG. 2. Normalized Einstein-frame potential V_E/V_0 with $V_0 = \kappa^2 m_P^4 / (16\xi^2)$ as a function of s and h , for $M = 0.01$, $\xi = 10^3$, $\gamma_4 = 0.3$, and $\beta_\kappa = 0$. The black line with an arrow indicates the inflaton trajectory along the h direction, with blue dots marking key points.

The corresponding potential is illustrated in Fig. 2, where V_E/V_0 is shown as a function of the real scalar fields s and h . The potential exhibits a saddlelike structure, with a pronounced stabilization along the s -direction. In particular, the presence of the positive γ_4 term generates an effective mass for s , driving it rapidly to its minimum at $s = 0$. As a result, the inflationary dynamics effectively reduce to a single-field trajectory along the h -direction, as indicated by the black line in Fig. 2. The stabilization mechanism is robust, with the mass of the s field scaling approximately as $\sim \sqrt{\gamma_4 \xi h^2}$.

Setting $\beta_\kappa = 0$ simplifies the structure of the potential and highlights the dominant role of γ_4 in stabilizing the transverse direction. When a nonzero β_κ is introduced, it induces only small corrections to the potential, and the stabilization of the s -field remains intact, as verified through numerical analysis.

IV. EFFECTIVE SINGLE FIELD INFLATION

With an appropriate value of γ_4 , s stabilizes at the origin and the scalar potential in the Einstein frame takes the following form,

$$V_E = \kappa^2 M^4 \frac{\left(1 - \frac{h^2}{(2M)^2} - \beta_\kappa \frac{h^4}{(2M)^4} \left(\frac{M}{m_P} \right)^2 \right)^2}{\left(1 + \frac{\xi h^2}{m_P^2} \right)^2}, \quad (19)$$

where $\xi \equiv \frac{\chi}{8} - \frac{1}{6}$.

After performing the conformal rescaling, the canonically normalized inflaton field $\hat{h}(h)$ in the Einstein frame becomes a function of the original field h , with

$$J(h) \equiv \left(\frac{d\hat{h}}{dh} \right) = \sqrt{\frac{1}{\Omega^2} + \frac{3}{2} m_P^2 \left(\frac{d \ln \Omega^2}{dh} \right)^2}. \quad (20)$$

The slow-roll parameters can then be expressed in terms of h as

$$\begin{aligned}\epsilon(h) &= \frac{m_P^2}{2} \left(\frac{V'_E}{JV_E} \right)^2, & \eta(h) &= m_P^2 \left(\frac{V''_E}{J^2 V_E} - \frac{J' V'_E}{J^3 V_E} \right), \\ \zeta^2(h) &= m_P^4 \frac{V'_E}{JV_E^2} \left(\frac{V'''_E}{J^3} - \frac{3J' V''_E}{J^4} - \frac{J'' V'_E}{J^4} + \frac{3(J')^2 V'_E}{J^5} \right),\end{aligned}\quad (21)$$

where primes denote derivatives with respect to h .

To first order in the slow-roll approximation, the scalar spectral index n_s , the tensor-to-scalar ratio r and the running of the scalar spectral index are given by

$$\begin{aligned}n_s &\simeq 1 - 6\epsilon(h_0) + 2\eta(h_0), & r &\simeq 16\epsilon(h_0), \\ \alpha_s &\equiv \frac{dn_s}{d \ln k} \simeq 16\epsilon(h_0)\eta(h_0) - 24\epsilon^2(h_0) - 2\zeta^2(h_0),\end{aligned}\quad (22)$$

evaluated at the field value h_0 , which corresponds to a given number of e-folds. The number of e-folds before the end of inflation is

$$N_0 = \frac{1}{\sqrt{2}m_P} \int_{h_e}^{h_0} \frac{J(h)}{\sqrt{\epsilon(h)}} dh, \quad (23)$$

where h_e denotes the end of inflation, determined by the condition $\epsilon(h_e) = 1$. The field value h_0 is associated with the pivot scale k_0 , at which the amplitude of the scalar power spectrum is normalized by Planck [8] as

$$A_s(k_0) = \frac{V_E(h)}{24\pi^2 m_P^4 \epsilon(h)} \Big|_{h(k_0)=h_0} = 2.137 \times 10^{-9}, \quad (24)$$

with $k_0 = 0.05 \text{ Mpc}^{-1}$.

Using the mass matrix defined in [71], together with the Einstein-frame potential given in Eq. (19), the slow-roll parameters in the large- ξ limit and to leading order in β can be approximated as

$$\epsilon \simeq \frac{4}{3\psi^4} + \frac{2\beta_\kappa}{3} \frac{1}{\xi}, \quad \eta \simeq -\frac{4}{3\psi^2} + \frac{\beta_\kappa \psi^2}{3\xi}, \quad \zeta^2 \simeq \frac{16}{9\psi^4} + \frac{4\beta_\kappa}{9\xi}. \quad (25)$$

Here, $\psi \equiv \sqrt{\xi}h/m_P$ is a dimensionless field variable. At the pivot scale, its value is approximately $\psi_0 \approx \sqrt{4N_0/3}$, while at the end of inflation it is given by $\psi_e \approx (4/3)^{1/4} < \psi_0$. The corresponding expressions for the scalar spectral index n_s , the tensor-to-scalar ratio r , and the running of the scalar spectral index α_s are

$$\begin{aligned}n_s &\simeq 1 - \frac{2}{N_0} + \frac{8\beta_\kappa}{9\xi} N_0, & r &\simeq \frac{12}{N_0^2} + \frac{32\beta_\kappa}{3\xi}, \\ \alpha_s &\simeq -\frac{2}{N_0^2} - \frac{8\beta_\kappa}{9\xi}.\end{aligned}\quad (26)$$

These analytical expressions are valid in the limit $\beta_\kappa \ll \xi/N_0^2$, i.e., for sufficiently small β_κ , where terms of order $(\beta_\kappa N_0^2/\xi)^2$ and higher are neglected. They clearly demonstrate that both the scalar spectral index n_s and the tensor-to-scalar ratio r increase with increasing β_κ while the running of scalar spectral index α_s decreases with increasing β_κ .

A brief comment on the sign of β_κ is also in order. Within the effective theory considered here, the sign of β (and hence that of $\beta_\kappa = \beta/\kappa$ for $\kappa > 0$) is not fixed by the gauge and R symmetries of the model. The two signs, however, lead to qualitatively different large-field behavior of the inflationary potential. For $\beta_\kappa > 0$, the higher-dimensional correction raises the potential relative to the asymptotic plateau and increases its slope. In contrast, $\beta_\kappa < 0$ lowers the potential and reduces its slope, eventually generating a hilltop followed by an additional large-field minimum. Although the latter minimum typically occurs in the trans-Planckian regime, where the truncated effective theory cannot be reliably extrapolated, the preceding local maximum can occur at sub-Planckian field values. Inflation on the branch connected to the physical vacuum then requires appropriate initial conditions on the relevant side of the hilltop, introducing an additional initial-condition sensitivity for $\beta_\kappa < 0$.

The sign of β_κ is further constrained phenomenologically. It follows directly from Eq. (26) that $\beta_\kappa > 0$ shifts n_s upward relative to the standard Starobinsky/nonminimal-Higgs prediction, whereas $\beta_\kappa < 0$ shifts it downward. Since the recent ACT-based data combinations favor somewhat larger values of n_s , the positive- β_κ branch is phenomenologically preferred in the present analysis. We therefore focus on $\beta_\kappa > 0$, while emphasizing that its sign is not fixed by the low-energy symmetries of the model; a first-principles determination would require specifying the ultraviolet origin of the corresponding higher-dimensional operator.

A. ACT, SPT, P-ACT-LB, CMB-SPA+DESI

The latest ACT-DR6 release [11], especially when combined with DESI-DR1 [12], has sharpened constraints on the scalar tilt, pushing preferred values of the spectral index to the high side of the Planck 2018 [8] result $n_s = 0.9651 \pm 0.0044$. Earlier combined analyses already hinted at an upward shift, e.g., $n_s = 0.9683 \pm 0.0040$ [79]. The Planck + ACT + LB, DESI-DR1 combination (P-ACT-LB) now yields $n_s = 0.9743 \pm 0.0034$, differing from the Planck value by $\sim 2\sigma$ and, as emphasized in [11], placing classic plateaus such as Starobinsky R^2 inflation [16],

nonminimal Higgs inflation [52,71,80–82] in the large ξ limit, and many α -attractor realizations [18,83–85] under pressure. A complementary update from SPT-3G (CMB-SPA) [86] finds $n_s = 0.9679 \pm 0.0033$, close to Planck, but including DESI DR2 shifts the tilt to $n_s = 0.9726 \pm 0.0028$. Importantly, CMB-SPA and DESI DR2 show $\sim 2.8\sigma$ tension in Λ CDM [87], cautioning against firm model exclusions at the 2σ level based on combinations of datasets that are themselves in significant tension as discussed by the author of [87]. Nevertheless, the potential implications for B-mode targets remain notable: as emphasized by CMB-S4 and LiteBIRD forecasts [38,88], the current landscape effectively focuses attention on a small set of benchmark classes Starobinsky, Higgs, T -model α -attractors (continuous α), and Poincaré-disk α -attractors [18,89–91] with discrete α whose (n_s, r) predictions span the experimentally relevant discovery space.

V. EXACT NUMERICAL ANALYSIS

We now examine the inflationary predictions of the model and compare them with the latest cosmological constraints. In the limit $\beta = 0$, the model reduces to standard nonminimal Higgs inflation. For sufficiently large values of the nonminimal coupling, $\xi \gg 1$, the Einstein-frame potential develops a Starobinsky-like plateau, yielding the universal attractor predictions

$$(n_s, r, \alpha_s) \simeq \left(1 - \frac{2}{N_0}, \frac{12}{N_0^2}, -\frac{2}{N_0^2}\right). \quad (27)$$

Numerically, these become

$$\begin{aligned} (n_s, r, \alpha_s) &\simeq (0.960, 0.005, -8 \times 10^{-4}) \quad \text{for } N_0 = 50, \\ &\simeq (0.967, 0.003, -6 \times 10^{-4}) \quad \text{for } N_0 = 60. \end{aligned} \quad (28)$$

These are the well-known predictions of Starobinsky $\mathcal{R}^2$ inflation and are represented by the dashed curves in Figs. 3–5. The corresponding inflaton field value at horizon exit,

$$\frac{h_0}{m_P} \simeq \sqrt{\frac{4N_0}{3\xi}}, \quad (29)$$

remains sub-Planckian for $\psi_0 \gg 1$ or $\xi \gg N_0$.

In contrast, for $\xi \ll 1$ the conformal flattening becomes ineffective and the model approaches the quartic-potential limit,

$$(n_s, r, \alpha_s) \simeq \left(1 - \frac{3}{N_0}, \frac{16}{N_0}, -\frac{3}{N_0^2}\right), \quad (30)$$

yielding

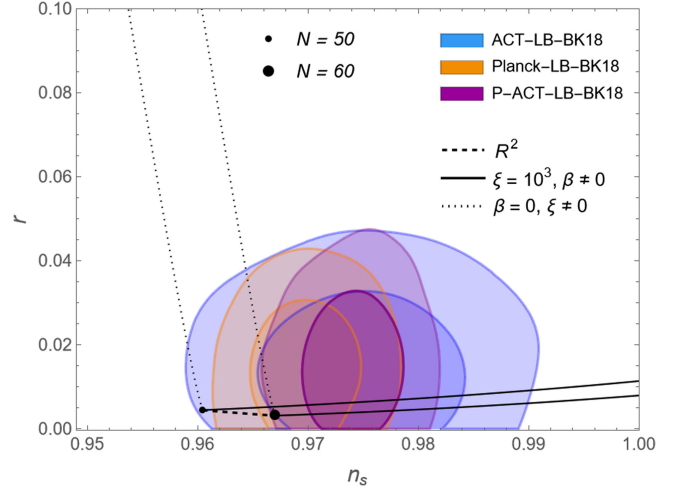


FIG. 3. Predictions in the (n_s, r) plane together with the latest ACT-LB-BK18, Planck-LB-BK18, and P-ACT-LB constraints. The dashed, dotted, and solid curves correspond to the Starobinsky $\mathcal{R}^2$ model, nonminimal Higgs inflation with $\beta = 0$, and nonminimal Higgs inflation with positive $\beta > 0$ and $\xi = 10^3$, respectively. In each case, the predictions are displayed for $N_0 = 50$ and $N_0 = 60$, assuming a fixed symmetry-breaking scale $M/m_P = 0.01$.

$$\begin{aligned} (n_s, r, \alpha_s) &\simeq (0.940, 0.320, -1.20 \times 10^{-3}) \quad \text{for } N_0 = 50, \\ &\simeq (0.950, 0.267, -8.33 \times 10^{-4}) \quad \text{for } N_0 = 60. \end{aligned} \quad (31)$$

These predictions lie well outside the current CMB bounds, as illustrated by the dotted curves in Figs. 3 and 4. The observational contours shown in these figures were obtained from publicly available Markov chain Monte Carlo chains released through the LAMBDA (Legacy Archive for Microwave Background Data Analysis) archive and analyzed using the GetDist package [92]. In the quartic-potential limit, the inflaton field value at horizon exit approaches $h_0/m_P \simeq \sqrt{8N_0}$, corresponding to trans-Planckian field excursions. Adopting the conservative requirement $h_0 < m_P$ therefore favors values of ξ larger than a few hundred.

The effect of the higher-dimensional operator becomes apparent once $\beta \neq 0$. In the large- ξ regime, the additional h^4 contribution arising from the shifted-hybrid-inflation sector modifies the shape of the Einstein-frame plateau. The resulting predictions for n_s , r , and α_s are given analytically in Eq. (26). In particular, the positive correction proportional to β_k increases the scalar spectral index while producing only a modest enhancement of the tensor-to-scalar ratio. Consequently, the model naturally accommodates the larger values of n_s favored by the ACT data while maintaining r well below current observational limits. The corresponding predictions for $\xi = 10^3$ are shown by the solid curves in Figs. 3 and 4 for $N_0 = 50$ and 60 . Values

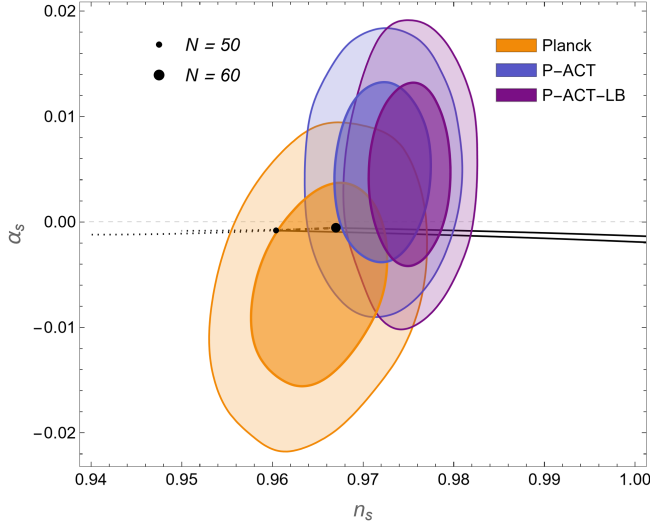


FIG. 4. Predicted values of the running of the scalar spectral index, $\alpha_s \equiv dn_s/d \ln k$, in the (n_s, α_s) plane. The shaded contours show the marginalized 68% (1σ) and 95% (2σ) confidence regions obtained from the Planck (orange), P-ACT (blue), and P-ACT-LB (purple) data combinations. The dashed, dotted, and solid curves correspond to the Starobinsky $\mathcal{R}^2$ model, non-minimal Higgs inflation with $\beta = 0$, and nonminimal Higgs inflation with positive $\beta > 0$ and $\xi = 10^3$, respectively. In each case, the predictions are displayed for $N_0 = 50$ and $N_0 = 60$, assuming a fixed symmetry-breaking scale $M/m_P = 0.01$.

$\beta_\kappa \gtrsim 0.3$ typically lead to $n_s \gtrsim 0.99$, which lies outside the preferred ACT-Planck region.

The ACT DR6 analysis [11] allows for a scale-dependent spectral index and finds

$$\alpha_s = 0.0062 \pm 0.0052 \quad (32)$$

from the P-ACT-LB data combination, consistent with no statistically significant evidence for running. Although Planck alone mildly preferred negative running, the inclusion of ACT shifts the posterior toward slightly positive values while remaining compatible with $\alpha_s \sim 0$. Throughout the phenomenologically viable region of parameter space, our model predicts a small negative running, typically $|\alpha_s| \lesssim 10^{-3}$, in agreement with current observational constraints. Moreover, the predicted tensor-to-scalar ratio remains in the range $r \simeq 0.003$ – 0.01 , making the scenario potentially testable with future CMB B-mode polarization experiments such as LiteBIRD and CMB-S4.

Figure 5 displays the variation of κ , β_κ , and the normalized inflaton field value h_0/m_P as functions of n_s for $\xi = 10^3$, $M/m_P = 0.01$, and $N_0 = 50$ and 60 . For completeness, both positive and negative values of β are shown. Negative values of β shift the predictions toward smaller values of n_s , thereby worsening agreement with the ACT-preferred region, consistent with the analytical expressions in Eq. (26). Restricting to the 1σ region of

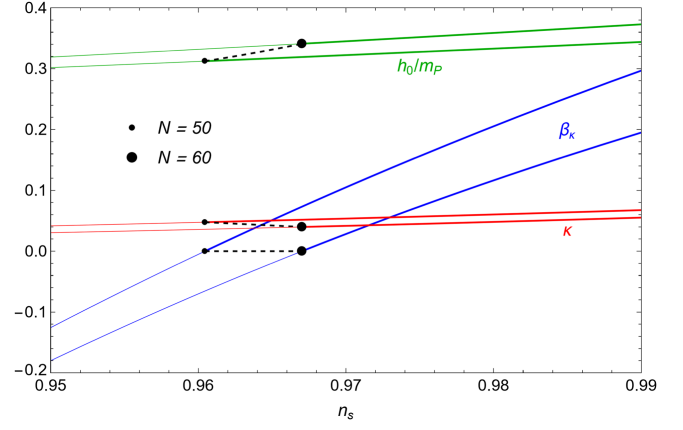


FIG. 5. Predicted values of the parameters β_κ and κ , together with the normalized inflaton field value at the pivot scale, h_0/m_P , as functions of the scalar spectral index n_s for fixed $\xi = 10^3$ and $M/m_P = 0.01$. The results are shown for $N_0 = 50$ and 60 e-folds. The green, blue, and red curves correspond to h_0/m_P , $\beta_\kappa = \beta/\kappa$, and κ , respectively, while the markers indicate the predictions in the nonminimal Higgs-inflation limit ($\beta = 0$). For each quantity, the thicker and thinner curves correspond to the $\beta > 0$ and $\beta < 0$ branches, respectively.

the P-ACT-LB-BK18 data set yields the approximate parameter ranges,

$$0.0006 \lesssim \beta \lesssim 0.01, \quad 0.04 \lesssim \kappa \lesssim 0.06, \quad (33)$$

$$0.01 \lesssim \beta_\kappa \lesssim 0.2, \quad 0.32 \lesssim \frac{h_0}{m_P} \lesssim 0.36. \quad (34)$$

These results demonstrate that even a modest higher-dimensional correction can significantly modify the inflationary predictions. Conversely, large values of β_κ drive the model away from the observationally allowed region, highlighting the importance of nonrenormalizable operators in determining the phenomenology of nonminimal Higgs inflation unless they are suppressed by an underlying ultraviolet completion.

The present framework should be distinguished from the pole-induced Higgs-inflation models of Ref. [93], which employ the same shifted-hybrid-inflation superpotential but within a different theoretical setting. In that work, the observationally viable low- r branch is realized by parametrizing the coefficient of the nonrenormalizable superpotential term as $\beta_\kappa = 1 + \delta_\lambda$, with $\delta_\lambda \sim 10^{-5}$, and by exploiting a hyperbolic Kähler geometry characterized by a curvature parameter N . The role of the higher-dimensional operator in the present model is fundamentally different. Rather than being introduced to cancel the effects of a Kähler pole, β_κ modifies the slope and curvature of the Einstein-frame inflationary plateau. As a result, the same class of R -symmetric Higgs-sector operators that appears in pole-induced Higgs inflation serves here as the key

ingredient for shifting the inflationary predictions toward the higher scalar spectral index favored by the ACT data, while maintaining a tensor-to-scalar ratio within the reach of forthcoming CMB observations.

To sum up, the recent ACT-DR6, Planck, and DESI data favor a higher scalar spectral index, $n_s \simeq 0.9743 \pm 0.0034$ (68% CL), while maintaining a stringent upper bound $r < 0.038$ (95% CL). These results place mild tension on the Starobinsky $\mathcal{R}^2$ model, whose predictions lie near the lower edge of the preferred n_s region. In contrast, the present model naturally accommodates the higher ACT-preferred values of n_s through the leading nonrenormalizable contribution parametrized by β_κ , while preserving a small tensor-to-scalar ratio within current observational bounds and potentially accessible to future CMB experiments.

A. Scanning the ξ – β_κ parameter space for compatibility with ACT-DR6 and complementary datasets

The analytical results presented in Sec. III provide useful insight into the dependence of the inflationary observables on the model parameters. To identify the regions favored by current cosmological data, we performed a comprehensive numerical scan over the nonminimal coupling ξ and the parameter $\beta_\kappa = \beta/\kappa$. Throughout the analysis we fixed the symmetry-breaking scale to $M/m_P = 0.01$ and adopted a reheating temperature $T_r = 10^8$ GeV, corresponding to $N_0 \simeq 53$ e-folds. For each parameter point, the background evolution was solved numerically using the canonically normalized inflaton field, while the coupling κ was determined by requiring consistency with the observed amplitude of scalar perturbations at the pivot scale $k_0 = 0.05$ Mpc $^{-1}$. As a theoretical prior, we imposed a sub-Planckian inflaton value at horizon exit, $h_0 \leq m_P$.

The results of the scan are summarized in Figs. 6–8. Figure 6 displays the predicted scalar spectral index n_s in the (ξ, β_κ) plane, with each point color-coded according to its n_s value using the CMasher colormap (cmr.ocean) [94]. The white (red) circles denote raw parameter points lying within the 1σ (2σ) regions of the scan, while the shaded bands represent smooth interpolations of the corresponding confidence regions. Superimposed contours indicate the preferred values of n_s inferred from the ACT-DR6, Planck, DESI, and CMB-SPA analyses, enabling a direct comparison between the model predictions and current observational constraints. Figures 7 and 8 show the corresponding predictions for the tensor-to-scalar ratio r and the running of the scalar spectral index $\alpha_s \equiv dn_s/d\ln k$, respectively.

Several important features emerge from the scan. First, increasing β_κ at fixed ξ systematically shifts the prediction toward larger values of n_s , in agreement with the analytical expressions derived in Sec. III. This allows the model to interpolate smoothly between the scalar tilt preferred by the CMB-SPA analysis,

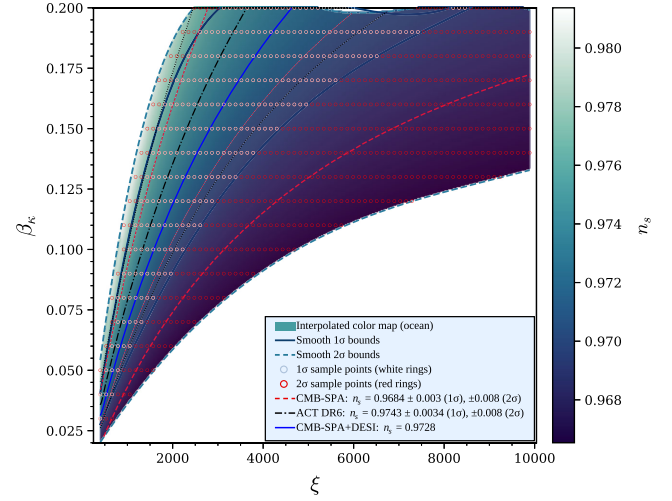


FIG. 6. Color-coded map of model predictions in the (ξ, β_κ) plane, where the color bar indicates the scalar index n_s . The interpolated 1σ (ocean) band is drawn above the 2σ (rainforest) band for clear separation; unfilled white rings (red rings) mark raw 1σ (2σ) samples. External CMB bounds appear as iso- n_s contours with shaded/hatched $\pm 1\sigma$ and $\pm 2\sigma$ intervals (SPT-(CMB-SPA): dashed crimson; ACT DR6(P-ACT-LB): dash-dot black). the blue line shows the combine (CMB-SPA+DESI). All displayed points satisfy the prior $h_0 \leq m_P$.

$$n_s = 0.9679 \pm 0.0033, \quad (35)$$

and the higher value favored by the ACT-DR6 + Planck + DESI combination,

$$n_s = 0.9743 \pm 0.0034. \quad (36)$$

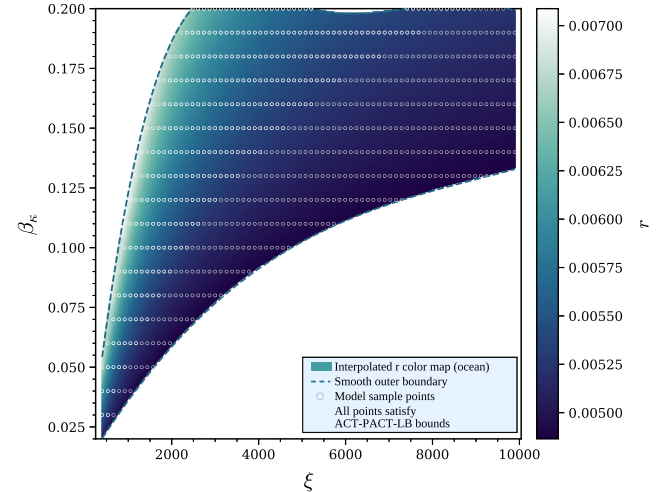


FIG. 7. Color-coded map of the model predictions in the (ξ, β_κ) plane, where the color scale represents the tensor-to-scalar ratio r . The displayed parameter points satisfy the sub-Planckian field condition $h_0 \leq m_P$. See Fig. 6 for a detailed description of the contours, sample points, and observational constraints.

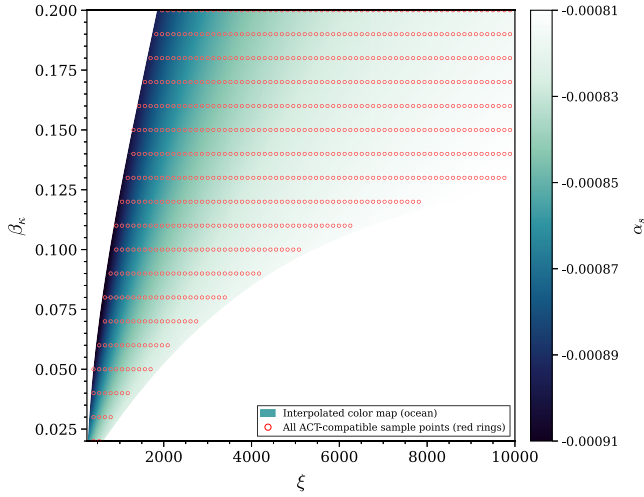


FIG. 8. Color-coded map of the model predictions in the (ξ, β_κ) plane, where the color scale indicates the running of the scalar spectral index α_s . The displayed parameter points satisfy the sub-Planckian field condition $h_0 \leq m_P$. See Fig. 6 for a detailed description of the contours, sample points, and observational constraints.

Second, the parameter combinations consistent with the ACT-preferred region form a broad corridor extending over approximately

$$10^2 \lesssim \xi \lesssim 10^4, \quad 0.02 \lesssim \beta_\kappa \lesssim 0.2, \quad (37)$$

while maintaining a sub-Planckian inflaton field value. Along this corridor, the coupling constants typically lie in the ranges

$$0.002 \lesssim \kappa \lesssim 0.06, \quad 10^{-5} \lesssim \beta \lesssim 10^{-2}, \quad (38)$$

demonstrating that the required deformation of the inflationary potential originates from a relatively modest higher-dimensional contribution.

The tensor-to-scalar ratio remains small throughout the phenomenologically viable region, as shown in Fig. 7. We find

$$0.004 \lesssim r \lesssim 0.007, \quad (39)$$

well below current observational limits but within the projected sensitivity of future CMB B -mode polarization experiments such as LiteBIRD and CMB-S4. Consequently, the model retains the prospect of an observable primordial gravitational-wave signal despite accommodating the larger scalar tilt favored by ACT.

The running of the scalar spectral index, shown in Fig. 8, remains negative throughout the allowed parameter space,

$$-10^{-3} \lesssim \alpha_s \lesssim -8 \times 10^{-4}, \quad (40)$$

with typical values of order 10^{-3} . These predictions are fully consistent with current observational constraints and constitute a robust prediction of the model.

Overall, the scan reveals a sizable and continuous region of parameter space simultaneously compatible with ACT-DR6, Planck, DESI, and CMB-SPA data. The model successfully accommodates the higher scalar spectral index preferred by ACT through the parameter β_κ , while preserving a small but potentially observable tensor-to-scalar ratio, a negative running of order 10^{-3} , and sub-Planckian inflaton field values. This demonstrates that the scenario remains both theoretically well-controlled and observationally testable in forthcoming CMB measurements.

A brief discussion of the sensitivity of the predictions of the present model to higher-dimensional operators is in order. The higher-dimensional operator proportional to β in Eq. (3) represents the leading nonrenormalizable correction to the superpotential. In general, the symmetries also allow higher-order operators of the form

$$\beta_p \bar{S} \bar{H}^c H^c \left(\frac{\bar{H}^c H^c}{m_P^2} \right)^p, \quad p = 2, 3, \dots \quad (41)$$

Successive terms are suppressed by additional powers of $\bar{H}^c H^c / m_P^2$. In the large- ξ regime, Eq. (29) gives $h_0 / m_P \simeq \sqrt{4N_0 / (3\xi)}$, and hence, along the D -flat direction,

$$\frac{\bar{H}^c H^c}{m_P^2} \simeq \frac{h_0^2}{4m_P^2} \simeq \frac{N_0}{3\xi}. \quad (42)$$

Thus, for $\xi \gtrsim 10^3$ and $N_0 \simeq 50$ – 60 , each additional power is suppressed at the percent level, provided the corresponding dimensionless coefficients are not parametrically enhanced. The effective-field-theory truncation is therefore well controlled in this sub-Planckian, large- ξ region. Toward the lower- ξ boundary, where h_0 approaches m_P , higher-dimensional operators can become more important and the predictions are correspondingly more sensitive to the ultraviolet completion.

VI. REHEATING AND LEPTOGENESIS

Following the end of inflation, the inflaton field oscillates about the minimum of its potential and transfers its energy to the thermal bath through its decay products, thereby reheating the Universe. In the present framework, an important decay channel is the production of the lightest right-handed neutrino (RHN), which subsequently enables nonthermal leptogenesis. The corresponding partial decay width is given by [95]

$$\Gamma_N = \frac{1}{64\pi} \left[\frac{M_N \Omega_0^{3/2}}{M J_0} \left(1 - \frac{12\xi M^2}{M_P^2} \right) \right]^2 m_{\text{inf}} \sqrt{1 - \frac{4M_N^2}{m_{\text{inf}}^2}}, \quad (43)$$

where

$$m_{\text{inf}} = \frac{\sqrt{2}\kappa M}{\Omega_0 J_0} \quad (44)$$

denotes the inflaton mass. The quantities

$$J_0 = \sqrt{\frac{1}{\Omega_0} + \frac{6\xi^2 M^2}{M_P^2 \Omega_0^2}}, \quad \Omega_0 = 1 + \frac{4\xi M^2}{M_P^2}, \quad (45)$$

are evaluated at the vacuum, and M_N represents the mass of the lightest RHN. For the parameter region of interest, we typically find $M_N \sim 10^{10}$ – 10^{12} GeV and $m_{\text{inf}} \sim 10^{12}$ – 10^{13} GeV.

In addition to RHN production, the inflaton can decay through Yukawa interactions. In supergravity realizations, the dominant contribution arises from the top-Yukawa operator $y_{33}^{u,\nu} Q_3 L_3 H_u$ [96]. For no-scalelike supergravity models, the corresponding decay width is given by [51]

$$\Gamma_y = \frac{3}{128\pi^3} \left(\frac{6\xi y \Omega_0^{3/2}}{J_0} \right)^2 \left(\frac{M}{M_P} \right)^2 \left(\frac{m_{\text{inf}}}{M_P} \right)^2 m_{\text{inf}}, \quad (46)$$

where $y \equiv y_{33}$ denotes the top Yukawa coupling.

The reheating temperature is determined by the total inflaton decay width,

$$T_r = \left(\frac{72}{5\pi^2 g_*} \right)^{1/4} \sqrt{\Gamma M_P}, \quad \Gamma = \Gamma_N + \Gamma_y, \quad (47)$$

with $g_* \simeq 228.75$ corresponding to the MSSM particle content. Assuming a standard postinflationary thermal history, the number of e-folds between horizon exit and the end of inflation can be approximated by [97]

$$N_0 = 53 + \frac{1}{3} \ln \left(\frac{T_r}{10^9 \text{ GeV}} \right) + \frac{2}{3} \ln \left(\frac{\sqrt{\kappa} M}{10^{15} \text{ GeV}} \right). \quad (48)$$

For our benchmark choice

$$M = 0.01 m_P, \quad \kappa = 0.0054, \quad \xi = 10^3, \quad (49)$$

we obtain

$$\Omega_0 \simeq 1.4, \quad J_0 \simeq 17.5, \quad m_{\text{inf}} \simeq 7 \times 10^{12} \text{ GeV}. \quad (50)$$

Taking a representative value $y \equiv y_{33} = 0.5$, the resulting RHN decay widths, reheating temperatures, and corresponding values of N_0 are summarized below:

M_N [GeV]	Γ_N [GeV]	T_r [GeV]	N_0
1.0×10^{10}	2.31×10^{-6}	2.91×10^8	52.960
1.0×10^{11}	2.31×10^{-4}	2.92×10^8	52.966
1.0×10^{12}	2.2×10^{-2}	2.98×10^8	52.975

These results indicate a relatively weak dependence of the reheating temperature on the RHN mass over the range $10^{10} \lesssim M_N \lesssim 10^{12}$ GeV. We find $T_r \simeq 2.91$ – 2.98×10^8 GeV, while the corresponding number of e-folds remains close to $N_0 \simeq 52.9$. Although the top-Yukawa channel provides the dominant contribution to the total decay width for the benchmark point, RHN production remains sufficiently efficient to generate the observed baryon asymmetry through nonthermal leptogenesis.

An important consistency requirement in supergravity inflationary models is the avoidance of gravitino overproduction [98]. The resulting bound constrains the reheating temperature and depends on the gravitino mass and its stability. For unstable gravitinos with masses above $\mathcal{O}(10)$ TeV, the constraint becomes relatively insensitive to $m_{3/2}$, whereas for stable gravitinos one typically requires $T_r \lesssim 10^9$ GeV [99–101]. The reheating temperatures obtained above, $T_r \simeq (2.91$ – $2.98) \times 10^8$ GeV, comfortably satisfy these bounds.

The generated RHNs decay out of equilibrium and produce a lepton asymmetry, a fraction of which is subsequently converted into the observed baryon asymmetry through electroweak sphaleron processes [102–104]. We, therefore, investigate the viability of nonthermal leptogenesis [41,42,67]. The resulting lepton asymmetry is approximately

$$\frac{n_L}{s} \simeq \frac{3T_r}{2m_{\text{inf}}} \left(\frac{\Gamma_N}{\Gamma} \right) \epsilon_1, \quad (51)$$

where ϵ_1 denotes the CP asymmetry generated in the decay of the lightest RHN. Assuming a hierarchical neutrino spectrum, one obtains [105]

$$\epsilon_1 \simeq -\frac{3m_{N_3} M_N}{8\pi v_u^2} \delta_{\text{eff}}, \quad (52)$$

where v_u is the vacuum expectation value of the up-type Higgs field, m_{N_3} is the heaviest light-neutrino mass, and δ_{eff} parametrizes the effective CP -violating phase [106]. Substituting this expression yields

$$\frac{n_L}{s} \simeq b \left(\frac{\Gamma_N}{\Gamma} \right) \left(\frac{T_r}{m_{\text{inf}}} \right) \left(\frac{M_N}{10^6 \text{ GeV}} \right) \left(\frac{m_{N_3}}{0.05 \text{ eV}} \right) \delta_{\text{eff}}, \quad (53)$$

where $b = 3.9 \times 10^{-10}$.

Throughout our analysis we adopt $m_{N_3} = 0.05$ eV and assume a hierarchical RHN spectrum,

$$M_N \ll M_{N_2} \ll M_{N_3}. \quad (54)$$

The lightest RHN mass is scanned subject to the requirements of nonthermal production, $M_N > T_r$, and kinematic accessibility, $M_N < m_{\text{inf}}/2$. Demanding successful leptogenesis while imposing the theoretical bound $|\delta_{\text{eff}}| \leq 1$,

TABLE II. Nonthermal leptogenesis: Minimal benchmark point satisfying $M_N > T_r$ and $M_N < m_{\text{inf}}/2$ with $|\delta_{\text{eff}}| \leq 1$.

Quantity	Value
M_N^*	9.5×10^{11} GeV
T_r^*	2.97×10^8 GeV
M_N^*/T_r^*	3.18×10^3
m_{inf}	7.0×10^{12} GeV
Γ_N^*	2.0×10^{-2} GeV
Γ_y^*	4.42×10^{-1} GeV
$n_L/s _{\delta_{\text{eff}}=1}$	6.39×10^{-10}
N_0^*	52.98
δ_{eff} for $n_L/s = 6.37 \times 10^{-10}$	0.96

we find the minimal viable solution

$$M_N^* \simeq 9.5 \times 10^{11} \text{ GeV}, \quad (55)$$

for which

$$\left. \frac{n_L}{s} \right|_{\delta_{\text{eff}}=1} \simeq 6.4 \times 10^{-10}, \quad (56)$$

in agreement with the observed baryon asymmetry [107]. The corresponding benchmark values are summarized in Table II. For this benchmark, reproducing the observed asymmetry requires $\delta_{\text{eff}} \simeq 0.96$, demonstrating that successful nonthermal leptogenesis can be realized within the allowed parameter space of the model.

VII. CONCLUSION

In this work, we have investigated a nonminimally coupled, R -symmetric Higgs inflation scenario embedded within the supersymmetric Pati-Salam framework. Employing the full Einstein-frame dynamics and fixing the inflationary scale through the observed amplitude of scalar perturbations, we performed a numerical analysis of the inflationary evolution without relying on slow-roll approximations. The postinflationary evolution was treated consistently through reheating and leptogenesis, allowing a unified study of inflationary predictions and their cosmological implications.

A comprehensive scan of the (ξ, β_κ) parameter space reveals an extended region that remains compatible with current cosmological observations. In particular, we find that for $\xi \sim 3 \times 10^2$ – 10^5 , $\beta_\kappa \sim 0.01$ – 0.2 , the model naturally predicts a scalar spectral index in the range preferred by the latest ACT-DR6 analyses, $n_s \simeq 0.971$ – 0.974 , while yielding a tensor-to-scalar ratio $r \simeq (4.6$ – $9) \times 10^{-3}$

throughout the central region of parameter space. Larger values, $r \lesssim 2 \times 10^{-2}$, are obtained near the boundaries of the observationally allowed region. We further find that the mild differences between the preferred spectral tilts inferred from ACT and SPT data can be accommodated through the intrinsic ξ – β_κ parameter degeneracy, without requiring any modification of the underlying framework.

The reheating dynamics were analyzed by including inflaton decays into right-handed neutrinos and Yukawa-coupled matter fields. For representative benchmark points, the resulting reheating temperature is of order $T_r \sim 10^8$ GeV, which is consistent with gravitino constraints and leads to a number of e-folds $N_0 \simeq 53$. Incorporating the neutrino sector naturally associated with the Pati-Salam symmetry, we showed that successful nonthermal leptogenesis can be realized while satisfying both the kinematic requirements for inflaton decay and the observed baryon asymmetry of the Universe. The inflationary and baryogenesis sectors therefore remain mutually consistent throughout the phenomenologically relevant parameter space.

Overall, the model provides a coherent and predictive realization of Higgs inflation within a supersymmetric grand unified framework, linking the inflationary epoch, reheating, neutrino physics, and the generation of the cosmic matter-antimatter asymmetry. The predicted tensor-to-scalar ratio lies within the projected sensitivity of upcoming CMB polarization experiments such as LiteBIRD, CMB-S4, and the Simons Observatory. Consequently, the parameter regions identified in this work will be subject to stringent experimental scrutiny in the near future. Improved measurements of primordial gravitational waves, together with progress in neutrino physics and baryogenesis studies, will provide important tests of the scenario and may offer valuable insight into the connection between inflation and grand unification.

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DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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