



Minimal warm higgs inflation in palatini R^2 gravity

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Received: 25 May 2025 / Accepted: 30 August 2025
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Abstract We investigate a novel realization of warm inflation driven by the Standard Model (SM) Higgs field within the Palatini formulation of gravity, supplemented by an R^2 curvature correction. By incorporating dissipative dynamics and a thermal radiation bath, our model eliminates the need for a separate reheating phase and naturally flattens the inflationary potential. We derive the slow-roll dynamics in the strong dissipation regime, compute the resulting inflationary observables, and demonstrate compatibility with the latest CMB constraints from Planck, BICEP/Keck, and BAO data. Specifically, we show that the scalar spectral index n_s lies in the range 0.963–0.965, and the tensor-to-scalar ratio r is strongly suppressed, $r \sim \mathcal{O}(10^{-3}–10^{-2})$, for Higgs quartic couplings $\lambda \sim 10^{-9}$. These predictions remain robust under two distinct growing-mode functions. Furthermore, we show that the model is consistent with the refined de Sitter swampland conjecture due to entropy production during inflation. This work represents the first comprehensive study of warm minimal Higgs inflation in Palatini R^2 gravity, unifying quantum gravity criteria, thermal field dynamics, and observational consistency within a single inflationary framework.

1 Introduction

One of the most minimalistic frameworks for explaining the early Universe's exponential expansion is to assume that the

Standard Model (SM) Higgs boson itself plays the role of the inflaton [1,2]. Such a scenario is particularly attractive as it does not require the introduction of new scalar fields or exotic physics beyond what is already part of the SM. This idea has been explored extensively since its inception and continues to be of great interest in both theoretical and observational cosmology. Traditionally, Higgs inflation models invoke a non-minimal coupling between the Higgs field ϕ and the Ricci scalar R in the form $\xi\phi^2 R$, see Refs. [2–7]. Additionally, the implementation has been successfully employed to the composite-field scenarios, see [8–10]. This interaction flattens the Higgs potential at high field values, enabling slow-roll inflation and ensuring consistency with current cosmological observations. However, this setup typically assumes the metric formulation of General Relativity (GR) and follows the so-called “cold inflation” paradigm [11–15], where inflation ends in a phase of reheating that bridges the transition to the radiation-dominated Universe [16–18].

Yet, there exists greater freedom in both the gravitational framework and the inflationary mechanism itself. In particular, the Palatini formulation of gravity [19] treats the metric $g_{\mu\nu}$ and connection Γ as independent variables, leading to modified dynamics when compared to the standard metric formulation. This distinction becomes particularly important in theories that involve extended gravity sectors, such as those with higher-order curvature terms or non-minimal couplings to matter fields. One especially promising modification is the inclusion of an R^2 term in the gravitational action, characterized by a dimensionless parameter α . When such a term is included in the Palatini formulation [20,21], it has been shown that the SM Higgs boson can act as the inflaton even in

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the absence of the $\xi\phi^2 R$ term. This avoids issues of unitarity violation and non-renormalizability typically associated with large non-minimal couplings. Moreover, both $\xi\phi^2 R$ and αR^2 terms can arise naturally as quantum corrections in curved spacetime, providing theoretical motivation for their inclusion [22].

An additional and compelling extension of the minimal scenario is to consider warm inflation rather than the conventional cold inflation picture. In warm inflation, dissipative effects are present during the inflationary phase itself, leading to continuous production of radiation [23–29]. This removes the need for a separate reheating phase and allows the Universe to smoothly transition to the radiation-dominated epoch. It also has notable implications for the dynamics of inflation and its observational predictions. For instance, thermal fluctuations, rather than quantum ones, can dominate the generation of the primordial curvature perturbations, and the inflationary slow-roll conditions can be relaxed due to the presence of dissipation.

In the context of Palatini gravity with an αR^2 term, incorporating warm inflation can address several challenges found in the cold version of the scenario. The presence of a thermal bath may alleviate the severe fine-tuning required for the Higgs self-coupling and can significantly modify the tensor-to-scalar ratio and spectral index, potentially bringing predictions into better agreement with CMB data. Furthermore, warm inflation opens the door for richer particle physics connections, especially relevant for the SM Higgs field, which is inherently coupled to other fields in the thermal plasma. While previous works have examined Higgs inflation in the Palatini formalism with R^2 corrections [20, 21, 30], and others have studied warm inflation dynamics separately [23–29, 31–37], a combined treatment has not yet been explored. Our work is the first to integrate warm inflation with a minimal Higgs potential in Palatini gravity including an R^2 term, thereby synthesizing three powerful frameworks into a single inflationary model. This synthesis enables us to address several challenges simultaneously: reducing the required Higgs self-coupling, suppressing tensor modes, ensuring a smooth transition to the radiation-dominated era, and satisfying swampland constraints via entropy generation. This positions our model as a theoretically robust and observationally viable alternative to conventional Higgs inflation.

The structure in the present work is the following: Sect. 2.1, we revisit the model of Higgs-driven inflation in the Palatini- R^2 framework. We then extend the model of Higgs-driven inflation in the Palatini- R^2 framework by embedding it in the warm inflationary paradigm in Sect. 2.2. In the section, we incorporate the warm inflationary paradigm into this setup by introducing dissipative effects and a thermal radiation bath, and derive the modified slow-roll dynamics and inflationary observables in the strong dissipation regime. In Sect. 3, we examine the compatibility of the model with the

refined de Sitter swampland conjecture by analyzing entropy production during inflation and its impact on quantum gravity constraints. Finally, in Sect. 4, we summarize our results and discuss the theoretical and observational implications of our findings, highlighting the viability of warm minimal Higgs inflation in the Palatini- R^2 framework as a consistent alternative to conventional inflationary models.

2 Warm inflation with R^2 term

2.1 A review on minimal Higgs- R^2 Palatini inflation

We begin by presenting the model under consideration, closely following Ref. [30]. The action in the Jordan frame is given by:

$$S_J = \int d^4x \sqrt{-g} \left(\frac{1}{2} M_P^2 R + \frac{\alpha}{4} R^2 - \frac{1}{2} \nabla_\mu \phi \nabla^\mu \phi - V(\phi) \right), \quad (1)$$

where g is the determinant of the metric $g_{\mu\nu}$, ϕ is the Standard Model Higgs field with potential $V(\phi)$, M_P is the reduced Planck mass, and $R = g^{\mu\nu} R_{\mu\nu}(\Gamma)$ is the Ricci scalar, which depends on both the metric and the independent connection Γ , and α is a dimensionless parameter. The covariant derivative ∇_μ is taken with respect to this connection. We assume the connection is torsion-free, i.e., $\Gamma_{\alpha\beta}^\lambda = \Gamma_{\beta\alpha}^\lambda$. While in GR the connection becomes the Levi-Civita connection, rendering the metric and Palatini formalisms equivalent, this is not the case when the gravity sector includes additional terms or matter fields couple non-minimally.

The action can be rewritten using an auxiliary field z [38]:

$$S_J = \int d^4x \sqrt{-g} \left(\frac{1}{2} M_P^2 (1 + \alpha z^2) R - \frac{\alpha}{4} z^4 - \frac{1}{2} \nabla_\mu \phi \nabla^\mu \phi - V(\phi) \right). \quad (2)$$

Performing a Weyl transformation given by

$$g_{\mu\nu} \rightarrow \Omega g_{\mu\nu}, \quad \Omega = 1 + \frac{\alpha z^2}{M_P^2}, \quad (3)$$

the action in the Einstein frame becomes

$$S_E = \int d^4x \sqrt{-g} \left(\frac{1}{2} M_P^2 R - \frac{1}{2\Omega} \nabla_\mu \phi \nabla^\mu \phi - \frac{V(\phi)}{\Omega^2} \right). \quad (4)$$

Taking the variation of action (4) with respect to the field z yields an algebraic constraint equation—rather than a dynamical equation involving derivatives of z . Similar to the Palatini

formalism, this indicates that the field is non-dynamical-and can be solved explicitly for [20, 21, 39]:

$$\frac{z^2}{M_P^2} = \frac{4V(\phi) + \nabla_\mu \phi \nabla^\mu \phi}{M_P^4 - \alpha \nabla_\mu \phi \nabla^\mu \phi}. \quad (5)$$

Substituting this into Eq. (4) gives

$$S_E \simeq \int d^4x \sqrt{-g} \left(\frac{1}{2} M_P^2 R - \frac{\nabla^\mu \phi \nabla_\mu \phi}{1 + 4\alpha \frac{V(\phi)}{M_P^4}} - \frac{V(\phi)}{1 + 4\alpha \frac{V(\phi)}{M_P^4}} \right), \quad (6)$$

where we have assumed $\dot{\phi}^2 \ll V(\phi)$, i.e., the field is in the slow-roll regime, and that spatial gradients of the field vanish throughout inflation. We now consider the Higgs potential during inflation to be $V(\phi) = \lambda \phi^4/4$, which is valid for large field values $\phi \gg v$, where $v = 246$ GeV represents the Higgs vacuum expectation value at zero temperature. A suitable field redefinition $\phi = \phi(\chi)$ is then introduced, determined by

$$\frac{d\phi}{d\chi} = \sqrt{1 + \lambda\alpha \left(\frac{\phi}{M_P} \right)^4}. \quad (7)$$

The action becomes:

$$S_E = \int d^4x \sqrt{-g} \left(\frac{1}{2} M_P^2 R - \frac{1}{2} \partial_\mu \chi \partial^\mu \chi - U(\chi) \right), \quad (8)$$

where the effective potential in terms of χ is:

$$U(\chi) = \frac{\lambda \phi^4(\chi)}{4 \left(1 + \lambda\alpha \frac{\phi^4(\chi)}{M_P^4} \right)} = \frac{\tilde{U}}{1 + 4\alpha \tilde{U}/M_P^4}. \quad (9)$$

where $\tilde{U}$ the potential in the case of $\alpha = 0$. As demonstrated in Ref. [21], the dynamics of inflation can be described using the conventional slow-roll parameters:

$$\epsilon \equiv \frac{1}{2} M_P^2 \left(\frac{U'}{U} \right)^2 = \frac{1}{1 + \alpha \lambda \left(\frac{\bar{\phi}}{M_P} \right)^4} \bar{\epsilon}, \quad (10)$$

$$\eta \equiv M_P^2 \frac{U''}{U} = \bar{\eta} - 3 \frac{\alpha \lambda \left(\frac{\bar{\phi}}{M_P} \right)^4}{1 + \alpha \lambda \left(\frac{\bar{\phi}}{M_P} \right)^4} \bar{\epsilon}, \quad (11)$$

$$\zeta = \frac{U'}{U} \frac{U'''}{U} = \bar{\zeta} + 12 \frac{\left(\alpha \lambda \left(\frac{\bar{\phi}}{M_P} \right)^4 \right)^2}{\left(1 + \alpha \lambda \left(\frac{\bar{\phi}}{M_P} \right)^4 \right)^2} \bar{\epsilon}^2 - 8 \frac{\alpha \lambda \left(\frac{\bar{\phi}}{M_P} \right)^4}{1 + \alpha \lambda \left(\frac{\bar{\phi}}{M_P} \right)^4} \bar{\epsilon} \bar{\eta}, \quad (12)$$

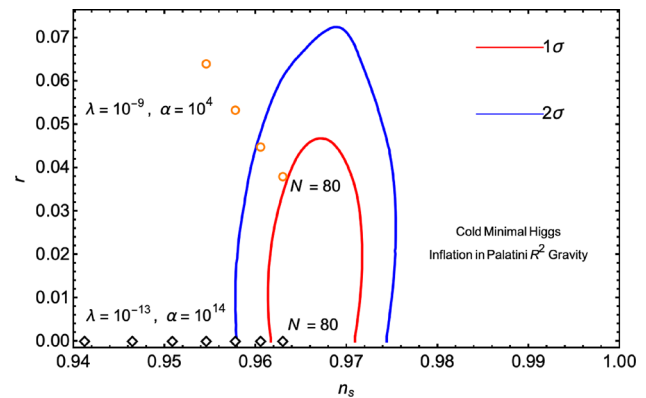


Fig. 1 We compare the theoretical predictions of (r, n_s) with Planck'18 results for TT, TE, EE, +lowE+lensing+BK15+BAO for cold Minimal Higgs Inflation in Palatini R^2 Gravity. For the plots, we have considered two set of parameters: (1) $\lambda = 10^{-9}$, $\alpha = 10^4$ and (2) $\lambda = 10^{-13}$, $\alpha = 10^{14}$ and $N \in [50, 80]$

where derivatives are taken with respect to χ , and $\bar{\epsilon}$, $\bar{\eta}$ represent the slow-roll parameters in the limit $\alpha = 0$, corresponding to the field $\bar{\phi}$. Since all expressions are formulated in terms of the standard $\lambda \bar{\phi}^4/4$ inflation model (hereafter referred to as $\lambda \bar{\phi}^4$ inflation), the familiar results of this model can be applied to extract leading-order expressions for key inflationary observables like the spectral index and tensor-to-scalar ratio. The number of e -folds N between horizon exit and the end of inflation is given by:

$$N = \frac{1}{M_P^2} \int_{\chi_f}^{\chi_i} \frac{U}{U'} d\chi = \int_{\bar{\phi}_f}^{\bar{\phi}_i} \frac{d\bar{\phi}}{M_P} \frac{1}{\sqrt{2\bar{\epsilon}}}, \quad (13)$$

where $\bar{\phi}_i$ and $\bar{\phi}_f$ denote the field values at horizon exit and end of inflation, respectively, with $\bar{\epsilon}(\bar{\phi}_f) = 1$. For the quartic potential $V = \lambda \bar{\phi}^4/4$, one finds $\bar{\epsilon} = 8M_P^2/\bar{\phi}_i^2 = 1/(N+1)$ and $\bar{\eta} = 3\bar{\epsilon}/2$. The leading-order results for the spectral index and tensor-to-scalar ratio become:

$$n_s^{\text{cold}} = 1 - 6\epsilon + 2\eta = 1 - \frac{3}{N+1}, \quad (14)$$

$$r^{\text{cold}} = 16\epsilon = \frac{16}{(N+1)(1 + 64\alpha\lambda(N+1)^2)}. \quad (15)$$

Notably, the expression for n_s remains identical to that of standard $\lambda \bar{\phi}^4$ inflation. We illustrate the theoretical predictions of the scalar spectral index n_s^{cold} and the tensor-to-scalar ratio r^{cold} within the framework of cold minimal Higgs inflation in Palatini R^2 gravity displayed in Fig. 1. These predictions are plotted against the most recent cosmological observational constraints from Planck 2018, including temperature and polarization anisotropies (TT, TE, EE), lowE, lensing, BICEP/Keck 2015 (BK15), and baryon acoustic oscillations (BAO). The figure highlights a fine-tuning issue: to reduce r enough to match observations, one must take extremely small $\lambda \sim 10^{-13}$, which is unnatural for the Standard Model

Higgs. As elaborated in Ref. [21], quantities derived from the power spectrum's derivatives-like the spectral index and its running-are unchanged at leading order, as the curvature power spectrum remains unaltered. Its amplitude is given by [40,43]:

$$\mathcal{P}_\zeta = \frac{1}{24\pi^2 M_p^4} \frac{U}{\epsilon} = \frac{1}{24\pi^2 M_p^4} \frac{\bar{U}}{\bar{\epsilon}} = \frac{2(N+1)^3 \lambda}{3\pi^2}, \quad (16)$$

with the observed value $\mathcal{P}_\zeta = 2.1 \times 10^{-9}$ [41]. The Planck satellite analysis (68% confidence level) provides:

$$n_s = 0.9625 \pm 0.0048, \quad (17)$$

using the TT,TE,EE+lowE+lensing dataset. The combined Planck and BICEP2/Keck Array data (95% confidence level) constrain:

$$r < 0.06. \quad (18)$$

These constraints apply at the pivot scale $k_* = 0.05 \text{ Mpc}^{-1}$. From Eq. (16), it's evident that α does not influence the amplitude $\mathcal{P}_\zeta$ at tree-level, confirming consistency with the $\lambda\phi^4$ result. For $N \gtrsim 50$, the quartic coupling must satisfy $\lambda \lesssim \mathcal{O}(10^{-12})$, which poses fine-tuning challenges, particularly for the SM Higgs. Tighter constraints on r with have been suggested, with $r < 0.036$ in Ref. [44]. In Ref. [39], the constraint $\sqrt{\lambda\alpha} \leq 10^{-2}$ was assumed. However, if $\alpha < 1$, the corresponding scale $M > M_p$ may conflict with quantum gravity. Therefore, in this work, we take $\alpha > 1$.

2.2 Slow-roll dynamics in warm minimal Higgs- R^2 Palatini inflation

The dynamics of warm inflation are described by a set of coupled differential equations: one governing the background evolution of the inflaton field, $\chi(t)$, and the other tracking the radiation energy density, $\rho_r(t)$. Within the Einstein frame and assuming a flat FLRW metric, the Friedmann equation during warm inflation is expressed as

$$H^2 = \frac{1}{3M_p^2} (\rho_\chi + \rho_r) = \frac{1}{3M_p^2} \left(\frac{1}{2} \dot{\chi}^2 + U(\chi) + \rho_r \right), \quad (19)$$

where $\dot{\chi} = d\chi/dt$, and ρ_r represents the energy density of the radiation fluid, whose equation of state is given by $w_r = 1/3$. Observational bounds, such as those from Planck 2018 combined with BK15, place an upper limit on the Hubble scale during inflation: $H_*/M_p < 2.5 \times 10^{-5}$ at 95% confidence level [41].

The inflaton field evolves according to the following equation:

$$\ddot{\chi} + 3H\dot{\chi} + U'(\chi) = -\Gamma\dot{\chi}, \quad (20)$$

with $U'(\chi) = dU/d\chi$. This is equivalent to the continuity equation for the inflaton energy density ρ_χ :

$$\dot{\rho}_\chi + 3H(\rho_\chi + p_\chi) = -\Gamma(\rho_\chi + p_\chi), \quad (21)$$

where the pressure is defined as $p_\chi = \dot{\chi}^2/2 - U(\chi)$, and the combination $\rho_\chi + p_\chi = \dot{\chi}^2$ corresponds to the kinetic contribution. The right-hand side acts as a source term accounting for energy transfer from the inflaton to radiation. For the radiation component, its energy density evolves as

$$\dot{\rho}_r + 4H\rho_r = \Gamma\dot{\chi}^2. \quad (22)$$

A necessary condition for warm inflation is that $\rho_r^{1/4} > H$, ensuring that dissipative effects influence both the inflaton dynamics and the primordial perturbation spectrum. In the context of warm inflation operating under the slow-roll approximation, the Friedmann equation and the dynamical equations governing the inflaton field and the radiation component can be expressed in the following simplified form:

$$H^2 \approx \frac{1}{3M_p^2} U(\chi), \quad (23)$$

$$\dot{\chi} \approx -\frac{U'(\chi)}{3H(1+Q)}, \quad Q \equiv \frac{\Gamma}{3H}, \quad (24)$$

$$\rho_r \approx \frac{\Gamma\dot{\chi}^2}{4H}, \quad \rho_r = C_r T^4, \quad (25)$$

where Q denotes a dissipative coefficient and $C_r = g_* \pi^2/30$. To obtain the above expressions, we have used the following approximations:

$$\rho_r \ll \rho_\chi, \quad \rho_\chi = \frac{1}{2} \dot{\chi}^2 + U, \quad (26)$$

$$\dot{\chi}^2 \ll U(\chi), \quad (27)$$

$$\ddot{\chi} \ll 3H(1+Q)\dot{\chi}, \quad (28)$$

$$\dot{\rho}_r \ll 4H\rho_r, \quad (29)$$

According to Refs. [45,46], the dissipation coefficient is generally taken as

$$\Gamma = C_m \frac{T^m}{\chi^{m-1}}, \quad (30)$$

where m is an integer, and C_m encapsulates the microphysical details of the dissipation process. Various choices of m correspond to different theoretical frameworks. For example, $m = 1$ is associated with high-temperature scenarios [28,47,55], while $m = 3$ arises in supersymmetric models [28,46,51]. Minimal models of warm inflation have also been explored [60–62]. Notably, thermal effects can greatly suppress the tensor-to-scalar ratio r , and can lead to distinctive non-Gaussian signatures. Besides the Hubble friction term,

the additional damping from Γ plays a crucial role in warm inflation. In the slow-roll approximation, the dynamics simplify considerably. Notably, the temperature can be expressed in terms of the scalar field, χ . In this study, we consider the temperature for $Q \gg 1$ to be represented in the following form:

$$T = \left(\frac{U'^2 \chi^{m-1}}{4H C_m C_r} \right)^{\frac{1}{4+m}}, \quad (31)$$

with m being any integer. Various choices for m have been explored in the literature [45, 46, 52]. Specifically: (1) $m = 1$, which corresponds to the high-temperature regime, as discussed in [28, 47, 55]; and (2) $m = 3$, motivated by supersymmetric models [28, 46, 51], and featured in minimal warm inflation scenarios [60–62]. In this work we consider solely $m = 1$. Note that, a high-temperature realization ($m = 1$) can arise when the inflaton couples directly to fermionic fields via a Yukawa-like interaction $L_{int} = y\chi\bar{\psi}\psi$, which leads to $\Gamma \sim T$ under the assumption of a large thermal bath and on-shell decay processes. Such models are discussed in [16, 19, 49] and can be aligned with SM Yukawa couplings, particularly involving top quarks or electroweak gauge bosons during high-temperature epochs [28, 47, 55]. In warm inflation, the slow-roll parameters are slightly modified and they take the form

$$\epsilon = \frac{1}{2} M_P^2 \left(\frac{U'}{U} \right)^2 = \frac{1}{1 + \alpha\lambda \left(\frac{\bar{\phi}}{M_P} \right)^4} \bar{\epsilon}, \quad (32)$$

$$\eta = M_P^2 \frac{U''}{U} = \bar{\eta} - 3 \frac{\alpha\lambda \left(\frac{\bar{\phi}}{M_P} \right)^4}{1 + \alpha\lambda \left(\frac{\bar{\phi}}{M_P} \right)^4} \bar{\epsilon}, \quad (33)$$

$$\beta = M_P^2 \left(\frac{U' \Gamma'}{U \Gamma} \right) = M_P \frac{\bar{\Gamma}'}{\bar{\Gamma}} \sqrt{\bar{\epsilon}}. \quad (34)$$

Similar to cold inflation, inflationary phase of the universe in warm inflation takes place when the slow-roll parameters satisfy the following conditions:

$$\epsilon \ll 1 + Q, \quad \eta \ll 1 + Q, \quad \beta \ll 1 + Q. \quad (35)$$

Additionally, the number of e-foldings, N , in warm inflation gets modified and it can be written for $Q \gg 1$ as

$$N = \frac{1}{M_P^2} \int_{\chi_{\text{end}}}^{\chi_N} \frac{Q U}{U'} d\chi = \int_{\bar{\phi}_f}^{\bar{\phi}_i} \frac{d\bar{\phi}}{M_P} \frac{Q}{\sqrt{2\bar{\epsilon}}}. \quad (36)$$

The power spectrum of the warm inflation was calculated in Refs. [26, 27, 53, 54, 56–59] and it reads

$$\Delta_{\mathcal{R}} = \left(\frac{H_N^2}{2\pi \dot{\chi}_N} \right)^2 \left(1 + 2n_N + \left(\frac{T_N}{H_N} \right) \frac{2\sqrt{3}\pi Q_N}{\sqrt{3+4\pi Q_N}} \right) G(Q_N), \quad (37)$$

Here, the subscript “ N ” denotes quantities evaluated at the Hubble horizon crossing, and $n = 1/(\exp H/T - 1)$ represents the Bose-Einstein distribution function. Additionally, the function $G(Q_N)$ captures the coupling between the inflaton and radiation within the heat bath, which results in a growing mode in the fluctuations of the inflaton field, as initially explored in Ref. [53], with subsequent implications [56]. In addition, the scalar spectral index is defined as

$$n_s - 1 = \frac{d \ln \Delta_{\mathcal{R}}}{d \ln k} \Big|_{k=k_N} = \frac{d \ln \Delta_{\mathcal{R}}}{d N}, \quad (38)$$

with $\ln k \equiv a H = N$. The tensor-to-scalar ratio of the perturbation, r , can be determined using the following formula:

$$r = \frac{\Delta_T}{\Delta_{\mathcal{R}}}, \quad (39)$$

where Δ_T represents the power spectrum of the tensor perturbation and follows the same form as in standard (cold) inflation, i.e., $\Delta_T = 2H^2/\pi^2 M_P^2 = 2U(\chi)/3\pi^2 M_P^4$. Thus, the power spectrum in Eq. (37) can be rewritten as follows:

$$\Delta_{\mathcal{R}} = \frac{(1 + Q_N)^2 \bar{U}}{24\pi^2 M_P^4 \bar{\epsilon}} \left(1 + 2n_N + \left(\frac{T_N}{H_N} \right) \frac{2\sqrt{3}\pi Q_N}{\sqrt{3+4\pi Q_N}} \right) G(Q_N). \quad (40)$$

The growth rate of the inflaton field fluctuation, resulting from the coupling between the inflaton and the radiation fluid in the thermal bath, is described by the function $G(Q_N)$ [53]. For Higgs-like and plateau-like potentials, the growing mode function was proposed in Ref. [54] and is expressed as follows:

$$G_1(Q_N) \simeq 1 + 0.18 Q_N^{1.4} + 0.01 Q_N^{1.8}, \quad (41)$$

whereas, the original growing mode function for warm little inflation was expressed as follows [55]:

$$G_2(Q_N) \simeq 1 + 0.335 Q_N^{1.364} + 0.0185 Q_N^{2.315}. \quad (42)$$

The function $G(Q)$ captures the growth of inflaton fluctuations due to coupling with radiation and must be basically determined. As noted in Ref. [54], this function also shows a slight dependence on the form of the scalar potential. For a quartic potential scenario, $G(Q)$ is provided in Eq. (42), while for Higgs-like and plateau-like potentials, it is given in Eq. (41). Note that at the thermalized inflaton fluctuation limit, $1 + 2n_N \simeq 2 T_N/H_N$ and $T_N/H_N = 3 Q_N/C_1$, one can re-write the power spectrum in the following form [54],

$$\Delta_{\mathcal{R}} \simeq \frac{5 C_1^3}{12\pi^4 g_* Q_N^2} \left(1 + \frac{\sqrt{3}\pi Q_N}{\sqrt{3+4\pi Q_N}} \right) G(Q_N), \quad (43)$$

where the approximation $\rho_r/V(\phi) = \epsilon Q/2(1+Q)^2$ has been used to obtain above relation. We note that the above power spectrum in this limit is in-explicitly dependent on the

inflaton potential [54]. Then, the tensor-scalar ratio r parameter in this case can be obtained by using Eqs. (39) and (43). It takes the form:

$$r = \frac{\Delta_T}{\Delta_{\mathcal{R}}} = 16\epsilon \left[\frac{6Q_N^3}{C_1} \left(1 + \frac{\sqrt{3}\pi Q_N}{\sqrt{3+4\pi Q_N}} \right) G(Q_N) \right]^{-1} \\ = 16 \frac{1}{1 + \alpha\lambda \left(\frac{\bar{\phi}}{M_P} \right)^4 \bar{\epsilon}} \\ \times \left[\frac{6Q_N^3}{C_1} \left(1 + \frac{\sqrt{3}\pi Q_N}{\sqrt{3+4\pi Q_N}} \right) G(Q_N) \right]^{-1}. \quad (44)$$

The spectral index of the power spectrum, using the growing mode function in Eq. (41), is given by [48,54,56]

$$n_s = 1 + \frac{Q_N}{3+5Q_N} \frac{(6\epsilon-2\eta)}{\Delta_{\mathcal{R}}} \frac{d\Delta_{\mathcal{R}}}{dQ_N}, \quad (45) \\ \frac{d\Delta_{\mathcal{R}}}{dQ_N} = \frac{5C_1^3}{12\pi^4 g_*} \left[\frac{1}{Q_N^2} \left(1 + \frac{\sqrt{3}\pi Q_N}{\sqrt{3+4\pi Q_N}} \right) \frac{dG(Q_N)}{dQ_N} \right. \\ \left. - \frac{1}{Q_N^3} \left(2 + \frac{\sqrt{3}\pi Q_N}{\sqrt{3+4\pi Q_N}} + \frac{2\sqrt{3}\pi^2 Q_N^2}{(3+4\pi Q_N)^{3/2}} \right) G(Q_N) \right].$$

We first take the growing mode function proposed in Ref. [54] $G(Q_N) \simeq 1 + 0.18 Q_N^{1.4} + 0.01 Q_N^{1.8}$, and consider $m = 1$. We then calculate Γ given in Eq. (30) to obtain

$$\Gamma = \frac{10\sqrt{3}C_1}{\sqrt{2}} \left(\frac{M_P^6}{C_1 C_r \lambda \bar{\phi}^6} \left(\frac{\lambda M_P^2 \bar{\phi}^4}{\alpha \lambda \bar{\phi}^4 + M_P^4} \right)^{5/2} \right)^{1/5}. \quad (46)$$

Using Eqs. (24) and (46), we find Q :

$$Q = \frac{2^{4/5} \lambda M_P^{10} \bar{\phi}^2}{3^{2/5} C_r} \frac{1}{\left(\frac{M_P^6}{C_1 C_r \lambda \bar{\phi}^6} \right)^{4/5} (\lambda M_P^2 \bar{\phi}^4)^2}. \quad (47)$$

When inflation ends, one finds using a condition $\epsilon_{\text{end}}(\phi_e) \approx Q_{\text{end}}(\phi_e)$ that

$$\frac{1}{1 + \alpha\lambda \left(\frac{\bar{\phi}}{M_P} \right)^4 \bar{\epsilon}} \approx \frac{2^{4/5} \lambda M_P^{10} \bar{\phi}^2}{3^{2/5} C_r} \frac{1}{\left(\frac{M_P^6}{C_1 C_r \lambda \bar{\phi}^6} \right)^{4/5} (\lambda M_P^2 \bar{\phi}^4)^2}. \quad (48)$$

Using the perturbation trick, $\phi_e = \phi_0 + g\psi$, with $g \equiv \alpha\lambda \ll 1$, we can solve Eq. (48) to obtain at the first-order approximation:

$$\frac{\phi_e}{M_P} \simeq \frac{4 \cdot 2^{3/4} \sqrt{3} C_r^{1/4} \lambda^{1/4}}{C_1}. \quad (49)$$

Moreover, the inflaton field at the Hubble horizon crossing in the strong regime, ϕ_N , can be determined using Eq. (36) to approximately yield

$$\frac{\phi_N}{M_P} \simeq \frac{16\sqrt{3} C_r^{1/4} \lambda^{1/4} N^{5/4}}{5\sqrt{5} C_1}. \quad (50)$$

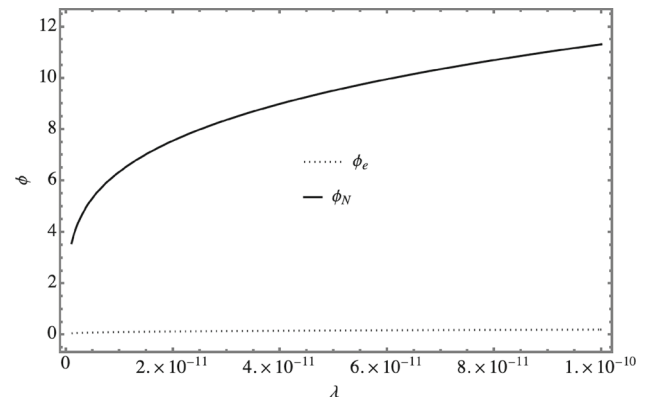


Fig. 2 We compare the amplitude of the scalar field at two distinct epochs of inflation-namely, at the end of inflation (ϕ_e) and 60 e-folds before the end ($\phi_{N=60}$)-as functions of the coupling constant λ . Here we have used $C_1 = 0.5$, $C_r = 70$, $\alpha = 10^4$

It is noticed that the inflaton field at the Hubble horizon crossing does not depend on a parameter α . A comparison of ϕ_N and ϕ_e is displayed in Fig. 2. More concretely, Fig. 2 illustrates the behavior of the scalar field amplitude at two distinct epochs of inflation-namely, at the end of inflation (ϕ_e/M_P) and 60 e-folds before the end ($\phi_{N=60}/M_P$)-as functions of the coupling constant λ . The plots provide valuable insights into how the dynamics of the inflaton field evolve with respect to changes in λ , a parameter often associated with the strength of self-interaction in scalar field potentials.

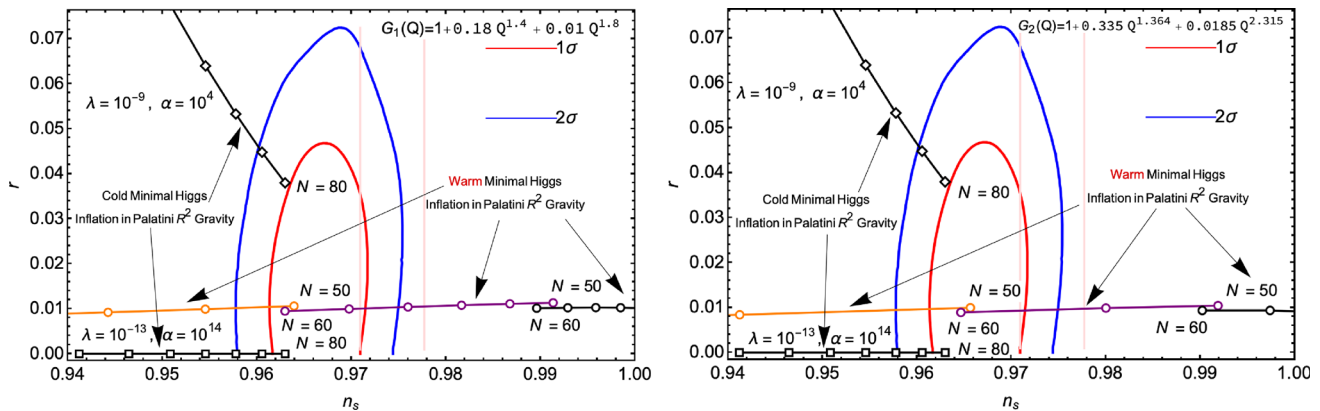
Figure 3 presents the theoretical predictions of the tensor-to-scalar ratio r versus the scalar spectral index n_s in the strong dissipation regime $Q \gg 1$, compared against the most recent observational constraints from Planck 2018 (TT, TE, EE + lowE + lensing) combined with BK15 and BAO data [41]. The left and right panels correspond to two different growing mode functions, $G_1(Q_N)$ and $G_2(Q_N)$, which model the enhancement of inflaton perturbations due to coupling with the radiation bath.

These plots demonstrate that warm inflation in the Palatini- R^2 framework is capable of producing viable predictions for a wide range of parameters. Notably, the values of (n_s, r) shift with varying λ and N , moving the predictions into regions of parameter space that are consistent with observations. For both growing mode functions, increasing the number of e-folds N from 50 to 60 leads to a decrease in both n_s and r , as expected from slow-roll dynamics. The differences between $G_1(Q_N)$ and $G_2(Q_N)$ are modest but non-negligible, with $G_2(Q_N)$ yielding slightly lower values of r at fixed n_s , reflecting a stronger suppression of tensor modes, see also Table 1.

Overall, the results in Fig. 3 confirm that warm Higgs inflation in the Palatini formulation with an R^2 term can generate inflationary predictions well within current observational bounds. The model is especially appealing as it addresses

Table 1 Predictions for the spectral index n_s and tensor-to-scalar ratio r for different values of λ and e-folds N . We take $C_1 = 0.5$, $C_r = 70$. The predictions depend weakly on the α parameter and we can fix it for $\alpha \gg 1$

Model	λ	$N = 50$	$N = 55$	$N = 60$
$G_1(Q)$	$10^{-8.87}$	$n_s = 0.964$, $r = 0.0106$	$n_s = 0.939$, $r = 0.0089$	$n_s = 0.908$, $r = 0.0072$
	$10^{-9.07}$	$n_s = 0.991$, $r = 0.0113$	$n_s = 0.979$, $r = 0.0106$	$n_s = 0.963$, $r = 0.0095$
$G_2(Q)$	$10^{-8.87}$	$n_s = 0.966$, $r = 0.0100$	$n_s = 0.941$, $r = 0.0084$	$n_s = 0.911$, $r = 0.0069$
	$10^{-9.07}$	$n_s = 0.992$, $r = 0.0105$	$n_s = 0.980$, $r = 0.0100$	$n_s = 0.965$, $r = 0.0089$

**Fig. 3** We compare the theoretical predictions of (r, n_s) with Planck'18 results for TT, TE, EE, +lowE+lensing+BK15+BAO [41] in the strong limit $Q \gg 1$ for a linear form of the growing mode function $G_1(Q_N)$ (left panel) and $G_2(Q_N)$ (right panel). For the plots, we have used $C_1 = 0.5$, $C_r = 70$, $\lambda = [10^{-8.87}, 10^{-9.07}, 10^{-9.27}]$, and from left to right $N = [50, 60]$. The predictions depend weakly on the α

parameter and we can fix it for $\alpha \gg 1$. For a cold scenario Eqs. (14) and (15), we have considered two set of parameters: (1) $\lambda = 10^{-9}$, $\alpha = 10^4$ and (2) $\lambda = 10^{-13}$, $\alpha = 10^{14}$ and $N \in [50, 80]$. The two-vertical bars denote 1σ confident region of the Atacama Cosmology Telescope (ACT) data [42], for which $n_s = 0.9743 \pm 0.0034$

both the fine-tuning problems of cold Higgs inflation and the constraints from the swampland conjectures, while maintaining consistency with the measured power spectrum of the Cosmic Microwave Background (CMB).

To further highlight the merits of warm inflation within the Palatini- R^2 framework, we compare our results with the corresponding predictions from the cold inflation scenario, governed by Eqs. (14) and (15). These expressions show that the scalar spectral index n_s^{cold} depends solely on the number of e-folds N and remains independent of both λ and α , resulting in a narrow prediction band. For instance, for $N = 60$, one finds $n_s^{\text{cold}} \approx 0.9508$. The tensor-to-scalar ratio r^{cold} is suppressed for large α but remains highly sensitive to λ , which must be fine-tuned to values as small as $\mathcal{O}(10^{-13})$ to meet current observational constraints [30].

In contrast, the warm inflation predictions depicted in Fig. 3 exhibit several improvements. First, the scalar spectral index n_s can attain higher values, typically in the range $n_s = 0.963\text{--}0.965$, depending on the choice of the Higgs quartic coupling λ , the number of e-folds N , and the specific form of the growing mode function $G(Q_N)$. This flexibility allows warm inflation to better match the Planck 2018 data within a broader region of parameter space. Second,

the tensor-to-scalar ratio r is significantly suppressed due to the presence of strong dissipation ($Q \gg 1$), with values as low as $r \approx 0.0069\text{--}0.0113$ even for moderate values of $\lambda \sim 10^{-9}$. This suppression arises from the thermal damping effects intrinsic to warm inflation, allowing it to evade the tight observational bounds on r without requiring extreme flattening of the inflaton potential. Third, although the difference between the two growing mode functions is modest, the function $G_2(Q_N)$ yields slightly lower values of r at a fixed n_s compared to $G_1(Q_N)$, reflecting a more efficient suppression of tensor modes due to enhanced thermal back-reaction.

These features make warm inflation more flexible and observationally viable than its cold counterpart. The warm scenario accommodates steeper potentials and larger couplings while satisfying Planck and BK15 constraints, thereby alleviating the extreme fine-tuning typically required in cold Higgs inflation. Moreover, it would be beneficial to compare the predictions with the most recent results from ACT Planck 2018 and DESI, for which $n_s = 0.9743 \pm 0.0034$. The vertical bars in both plots correspond to the 1σ confidence region from the ACT data. We observe from Fig. 3 that the predictions for warm minimal Higgs inflation in the

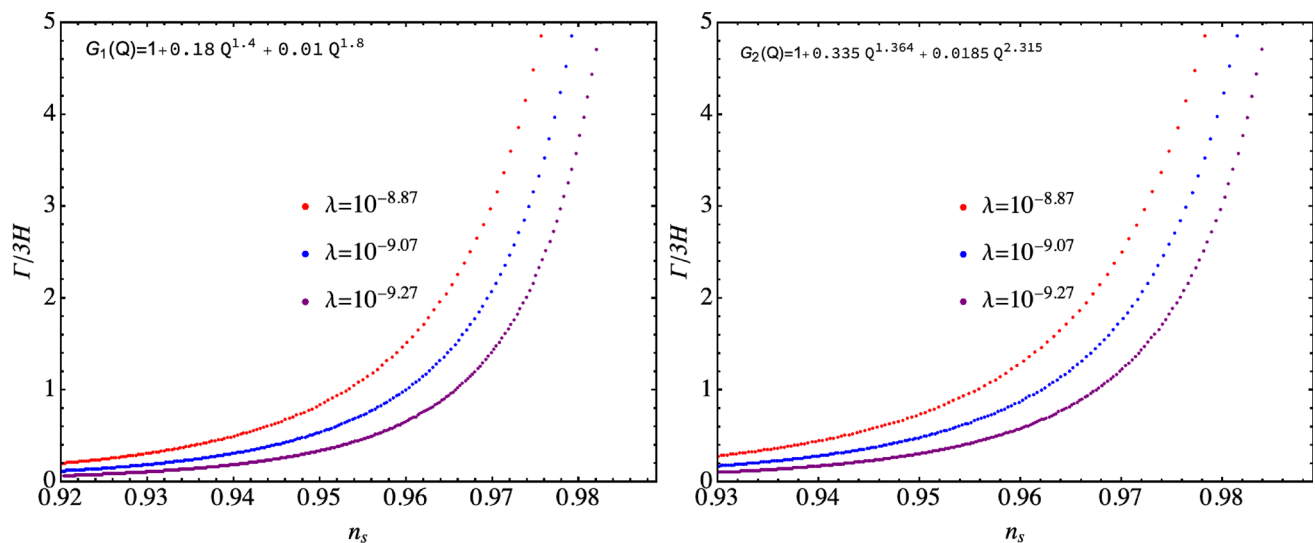


Fig. 4 The plots display the evolution of the ratio $Q = \Gamma/3H$ versus the primordial tilt n_s in the interpolation between the weak and strong dissipative regimes for the warm minimal Higgs inflation in Palatini R^2

for G_1 (left panel), and G_2 (right panel). In both panels, we have used three different values of the parameter $\lambda = [10^{-8.87}, 10^{-9.07}, 10^{-9.27}]$, and $C_1 = 0.7$, $C_r = 70$ and $\alpha = 10^4$

Palatini R^2 framework lie comfortably within the allowed Planck and ACT contours, particularly for lower values of λ and for N . More specifically, taking $G(Q) = G_1(Q)$ and using $\lambda = 10^{-9.07}$, and $C_1 = 0.5$, $C_r = 70$ and $\alpha = 10^4$, the results fit well within the 1σ confident region of ACT data. Similar behaviors have been observed for $G(Q) = G_2(Q)$. The inclusion of thermal fluctuations via the function $G(Q)$ significantly shifts the predictions towards higher n_s values while keeping r suppressed, a characteristic consistent with observational constraints. We now analyze the interpolation between the weak and strong dissipative regimes following Ref. [65]. In Fig. 4, presents the evolution of the dissipation ratio $Q = \Gamma/3H$ as a function of the scalar spectral index n_s in the context of warm minimal Higgs inflation in Palatini R^2 gravity. This analysis interpolates between the weak and strong dissipative regimes using two representative forms of the dissipation enhancement function $G(Q)$, denoted as $G_1(Q)$ and $G_2(Q)$. For this analysis, we fix the parameter $m = 1$, which characterizes the temperature dependence in the dissipation coefficient. In the left panel, corresponding to $G_1(Q) = 1 + 0.18Q^{1.4} + 0.01Q^{1.8}$, we observe a smooth and monotonic increase in the dissipation ratio Q with respect to n_s , reflecting the transition from the weak to the strong dissipative regime. Three different values of the self-coupling parameter λ are considered: $\lambda = 10^{-8.87}$, $10^{-9.07}$, and $10^{-9.27}$, represented by red, blue, and purple curves, respectively. Smaller values of λ lead to a slower increase in Q , indicating that weaker couplings delay the onset of strong dissipation for a given n_s . The right panel displays the same observable using the alternative dissipation enhancement function $G_2(Q) = 1 + 0.335Q^{1.364} + 0.0185Q^{2.315}$.

A qualitatively similar behavior is observed, with the dissipation ratio increasing more steeply at larger values of n_s . The variation in λ again controls the amplitude and sharpness of the transition, reinforcing the consistency of the warm inflation dynamics across different dissipation profiles. These results highlight the sensitivity of the dissipative dynamics to both the form of the $G(Q)$ function and the coupling parameter λ . Importantly, the plots confirm that in both models, a transition from weak ($Q \ll 1$) to strong ($Q \gg 1$) dissipation can naturally occur within the allowed Planck 2018 range of the scalar spectral index $n_s = 0.9625 \pm 0.0048$. This supports the viability of the warm minimal Higgs inflation scenario in interpolating regimes and sets the stage for further constraints from tensor-to-scalar ratio r , field excursions, and reheating dynamics.

3 de Sitter swampland conjecture

In this section, we build on previous findings that highlighted the resilience of warm inflation (WI) in satisfying the de Sitter swampland criteria, especially with regard to entropy effects [50]. In the WI paradigm, the inflaton field interacts with other matter fields, inducing dissipation and sustaining a radiation bath during inflation. This contrasts with Cold Inflation (CI), where expansion leads to cooling. The resulting inflaton dynamics are governed by

$$\ddot{\chi} + (3H + \Gamma)\dot{\chi} + U' = 0, \quad (51)$$

where Γ is the dissipation coefficient and H is the Hubble parameter. The radiation energy density evolves via

$$\dot{\rho}_r + 4H\rho_r = \Gamma\dot{\chi}^2, \quad (52)$$

with ρ_r as the radiation density. Entropy density s is related by $Ts = (1 + \omega_r)\rho_r$, where $\omega_r = 1/3$, and $\rho_r = \alpha T^4$, with $\alpha = \pi^2 g_*/30$ and g_* denoting the effective relativistic degrees of freedom. Under slow-roll, WI yields

$$3H(1 + Q)\dot{\chi} \approx -U', \quad \rho_r \approx \frac{3Q}{4}\dot{\chi}^2, \quad Ts \approx Q\dot{\chi}^2, \quad (53)$$

where $Q = \Gamma/3H$. Given $U \gg \rho_r$, we approximate the Hubble rate as

$$H^2 \approx \frac{U}{3M_p^2}. \quad (54)$$

This leads to a temperature expression:

$$\alpha T^4 \approx \frac{(U')^2}{12H^2} \frac{Q}{(1 + Q)^2}. \quad (55)$$

In light of the Bousso entropy bound, we incorporate thermal entropy into the inequality:

$$N\gamma(M_{Pl}R)^\delta + \frac{4}{3}\pi s R^3 \leq 8\pi^2(M_{Pl}R)^2, \quad (56)$$

implying the constraint on degrees of freedom:

$$N\gamma M_{Pl}^\delta \leq R^{3-\delta} \left(8\pi^2 R^{-1} M_{Pl}^2 - \frac{4}{3}\pi s \right). \quad (57)$$

Differentiating the Friedmann equation gives the refined swampland condition:

$$\frac{|U'|}{U} > c_2 + \delta c_2 \equiv \tilde{c}_2, \quad (58)$$

where $c_2 \sim O(1)$, and the entropy correction δc_2 modifies it to

$$\tilde{c}_2 = c_2 \left| 1 - \frac{1}{3-\delta} \frac{1 - \frac{\dot{s}}{6\pi M_{Pl}^2 \dot{H}}}{1 + \frac{\dot{s}}{6\pi M_{Pl}^2 H}} \right|^{-1}. \quad (59)$$

We estimate the entropy effect by examining s/H , leading to

$$\frac{s}{H} \approx \frac{Q}{1+Q} \frac{2M_p^2 \epsilon_H}{H T}, \quad \epsilon_H \equiv \frac{\epsilon_U}{1+Q}, \quad \epsilon_U \approx \frac{M_{Pl}^2}{2} \left(\frac{U'}{U} \right)^2, \quad (60)$$

and thus,

$$\frac{s}{6\pi M_{Pl}^2 H} \ll 1. \quad (61)$$

This is valid in both weak and strong dissipative limits. The derivative term becomes

$$\frac{\dot{s}}{6\pi M_{Pl}^2 \dot{H}} = \frac{s}{6\pi M_{Pl}^2 H} \frac{s'H}{sH'}, \quad (62)$$

where $s'H/sH' \approx 6TU/(TU')$, leading to

$$\tilde{c}_2 \simeq c_2 \left| 1 - \frac{1}{3-\delta} \left[1 + \kappa \left(1 - 6 \frac{T'U}{TU'} \right) + \mathcal{O}(\kappa^2) \right] \right|^{-1}, \quad (63)$$

with

$$\kappa = \frac{s}{6\pi M_{Pl}^2 H}. \quad (64)$$

To examine this effect in WI, we take the generic form of the dissipation coefficient [28, 29, 45, 51, 52]:

$$\Gamma(T, \chi) = CT^n \chi^p M^{1-n-p}, \quad (65)$$

where C is a dimensionless constant, M is a mass scale, and n and p are numerical powers. In the strong regime ($Q \gg 1$), temperature and its derivative yield:

$$T(\chi) \simeq \left[\frac{\sqrt{3}M_p}{4\alpha} \frac{(U')^2}{C\chi^p U^{1/2} M^{1-n-p}} \right]^{\frac{1}{4+n}}. \quad (66)$$

Taking $p = 0$, we find

$$\begin{aligned} \frac{T'U}{TU'} \Big|_{Q>1} &\simeq \frac{1}{4+n} \left[2 \frac{UU''}{(U')^2} - \frac{1}{2} \right] \\ &= \frac{1}{4+n} \left[\frac{\alpha\lambda\bar{\phi}^4(\bar{\eta} - 3\bar{\epsilon}) + \bar{\eta}M_p^4}{M_p^4\bar{\epsilon}} - \frac{1}{2} \right], \end{aligned} \quad (67)$$

with $\bar{\phi}/M_p \sim \sqrt{8N}$. Using $\bar{\epsilon}$ and $\bar{\eta}$ in Sect. (2.1), we find:

$$\frac{T'U}{TU'} \Big|_{Q\gg 1} \sim \frac{1}{5} \left(1 - 96\alpha\lambda N^2 \right). \quad (68)$$

Equation (68) captures the behavior of the ratio $\frac{T'U}{TU'}$ in the strong dissipation limit $Q \gg 1$, highlighting how temperature derivatives respond to the evolution of the inflaton field and the form of the potential. Specifically, the expression $\frac{T'U}{TU'} \Big|_{Q\gg 1}$ demonstrates that this ratio becomes negative for sufficiently large values of the product $\alpha\lambda N^2$, implying a non-trivial interplay between thermal and potential gradients in the warm inflation regime. The suppression of this ratio with increasing $\alpha\lambda$ is essential because it modulates the entropy correction to the refined de Sitter swampland condition via the entropy flow term $\frac{\dot{s}}{H}$.

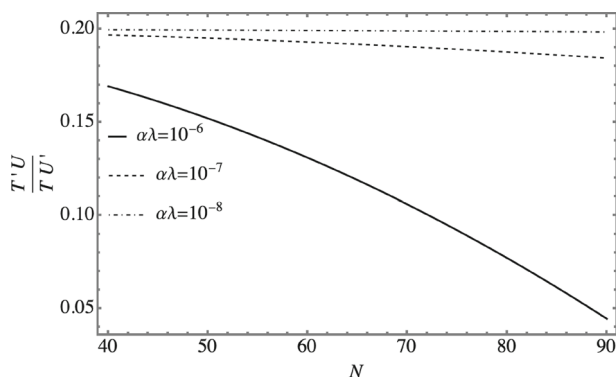


Fig. 5 A plot of the ratio $\frac{T'U}{TU'}|_{Q \gg 1}$ as a function of the number of e-folds N , evaluated under the strong dissipative regime ($Q \gg 1$) of warm inflation. The plot illustrates the impact of varying the product $\alpha\lambda$ on this ratio, with three representative values: 10^{-6} (solid line), 10^{-7} (dashed line), and 10^{-8} (dash-dotted line). As N increases, the ratio decreases for larger $\alpha\lambda$, indicating stronger sensitivity to dissipation in this regime

Figure 5 graphically illustrates the dependence of $\frac{T'U}{TU'}|_{Q \gg 1}$ on the number of e-folds N for different values of the product $\alpha\lambda$. As shown, increasing $\alpha\lambda$ suppresses the value of the ratio more rapidly as N increases, with the transition to negative values occurring at smaller N . This behavior is significant because it affects the entropy-induced shift in the swampland criterion (Eq. (58)). Therefore, warm inflation in the Palatini- R^2 framework provides a natural mechanism to address swampland constraints, especially in scenarios with sizable dissipation and thermal backreaction.

4 Conclusion

In this work, we have developed a comprehensive framework for warm inflation driven by the Standard Model (SM) Higgs field, embedded in the Palatini formulation of gravity with an R^2 correction. This model addresses several theoretical and observational challenges associated with conventional cold Higgs inflation, notably the issues of unitarity violation, reheating, and the fine-tuning of the Higgs self-coupling λ . By introducing dissipative effects during the inflationary epoch, our model maintains a thermal radiation bath, thereby removing the need for a separate reheating phase and enabling a smooth transition to the radiation-dominated Universe. The R^2 term in the Palatini formalism naturally flattens the effective potential, suppresses the tensor-to-scalar ratio r , and allows for viable inflationary dynamics with moderate values of λ .

We derived the modified slow-roll equations in the strong dissipation regime ($Q \gg 1$) and calculated the key inflationary observables. Our analysis reveals that the scalar spectral index n_s can attain values in the range 0.963–0.965,

while the tensor-to-scalar ratio is significantly suppressed to $r \sim 10^{-3}$ – 10^{-2} , all in excellent agreement with the latest Planck 2018 and BICEP/Keck observational bounds. These results hold for different forms of the growing mode function $G(Q)$, showing the robustness of the model under variations in the dissipation mechanism. Importantly, we have also explored the implications of this framework for the refined de Sitter swampland conjecture. By accounting for entropy production from the radiation bath, we showed that the model naturally satisfies the swampland criteria, even in the presence of steep potentials. The inclusion of thermal backreaction and entropy flow significantly modifies the constraints, highlighting a synergy between warm inflation and quantum gravity considerations.

Indeed, throughout our analysis, we have focused on the strong dissipative regime ($Q \gg 1$), where analytical approximations can be reliably applied to simplify the dynamics and obtain closed-form expressions for observables such as the scalar spectral index and tensor-to-scalar ratio. The intermediate regime ($Q \sim 1$), however, presents additional challenges. In this region, neither the weak ($Q \ll 1$) nor strong ($Q \gg 1$) approximations hold, and the dissipation coefficient cannot be treated as a dominant or subdominant perturbative quantity. As a result, analytical treatment becomes significantly more involved, requiring the use of exact expressions without dropping higher-order terms. For most models, especially with temperature-dependent dissipative coefficients, numerical methods are typically required to accurately capture the transition dynamics across regimes. Therefore, to maintain clarity and focus, we restricted our study to the $Q \gg 1$ regime. Nonetheless, we acknowledge that extending the analysis to include the intermediate dissipative regime would provide a more complete understanding of the model, and we plan to investigate this in future work using a numerical approach. A key advantage of embedding the Higgs inflationary scenario within the warm inflation framework, particularly in the strong dissipative regime $Q \gg 1$, is its ability to alleviate the severe fine-tuning required for the Higgs self-coupling λ . In the conventional cold inflation scenario within Palatini R^2 gravity, achieving consistency with both the observed scalar amplitude $\mathcal{P}_\zeta \sim 2.1 \times 10^{-9}$ and the stringent upper bound on the tensor-to-scalar ratio $r < 0.036$ necessitates an extremely small value of the quartic coupling, typically $\lambda \sim 10^{-13}$. Such a small coupling is difficult to justify within the Standard Model and poses a naturalness problem. In contrast, the warm inflationary dynamics significantly alter the generation of primordial perturbations. The presence of thermal fluctuations and dissipation leads to an enhancement of the scalar power spectrum through the growing mode function $G(Q)$ and thermal factors, while simultaneously suppressing the tensor amplitude due to additional friction and thermal damping. This dual effect allows successful inflationary predictions to be achieved with larger

values of $\lambda \sim 10^{-9}$, as demonstrated in Table 1 and Fig. 3 of our analysis. Quantitatively, this implies that warm Palatini-Higgs inflation permits a relaxation of the fine-tuning of λ by approximately four orders of magnitude compared to the cold case. This substantial improvement not only enhances the theoretical viability of Higgs inflation but also strengthens its connection to Standard Model physics by avoiding unnatural parameter choices. These findings highlight a central benefit of the warm inflation framework and underscore its significance in resolving one of the persistent challenges in minimal Higgs inflation models.

In summary, warm minimal Higgs inflation in Palatini- R^2 gravity provides a theoretically consistent and phenomenologically viable model for the early Universe. It unifies the benefits of thermal field theory, extended gravity, and particle physics minimalism, while remaining compatible with both CMB observations and swampland constraints. Future extensions of this work may involve constructing a UV-complete particle physics realization of the dissipation mechanism, exploring non-Gaussianities and isocurvature perturbations, and confronting predictions with upcoming precision cosmology experiments such as LiteBIRD [63] and CMB-S4 [64].

Acknowledgements This work was supported by the Ministry of Science and Higher Education of the Republic of Kazakhstan, Grant No. AP26101889.

Data Availability Statement This manuscript has no associated data. [Authors' comment: Data sharing not applicable to this article as no datasets were generated or analysed during the current study.]

Code Availability Statement This manuscript has no associated code/software. [Authors' comment: Code/Software sharing not applicable to this article as no code/software was generated or analysed during the current study.]

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References

- D.S. Salopek, J.R. Bond, J.M. Bardeen, Phys. Rev. D **40**, 1753 (1989)
- F.L. Bezrukov, M. Shaposhnikov, Phys. Lett. B **659**, 703–706 (2008)
- F.L. Bezrukov, A. Magnin, M. Shaposhnikov, Phys. Lett. B **675**, 88–92 (2009)
- S.C. Park, S. Yamaguchi, JCAP **08**, 009 (2008)
- F. Bezrukov, M. Shaposhnikov, JHEP **07**, 089 (2009)
- F. Bezrukov, A. Magnin, M. Shaposhnikov, S. Sibiryakov, JHEP **01**, 016 (2011)
- M.P. Hertzberg, JHEP **11**, 023 (2010)
- P. Channuie, J.J. Joergensen, F. Sannino, JCAP **05**, 007 (2011)
- F. Bezrukov, P. Channuie, J.J. Joergensen, F. Sannino, Phys. Rev. D **86**, 063513 (2012)
- P. Channuie, J.J. Joergensen, F. Sannino, Phys. Rev. D **86**, 125035 (2012)
- A.A. Starobinsky, Phys. Lett. B **91**, 99–102 (1980)
- K. Sato, Mon. Not. R. Astron. Soc. **195**(3), 467–479 (1981)
- A.H. Guth, Phys. Rev. D **23**, 347–356 (1981)
- A.D. Linde, Phys. Lett. B **108**, 389–393 (1982)
- A. Albrecht, P.J. Steinhardt, Phys. Rev. Lett. **48**, 1220–1223 (1982)
- A.D. Linde, Contemp. Concepts Phys. **5**, 1–362 (1990)
- A. Albrecht, P.J. Steinhardt, M.S. Turner, F. Wilczek, Phys. Rev. Lett. **48**, 1437 (1982)
- L.F. Abbott, E. Farhi, M.B. Wise, Phys. Lett. B **117**, 29 (1982)
- F. Bauer, D.A. Demir, Phys. Lett. B **665**, 222–226 (2008)
- I. Antoniadis, A. Karam, A. Lykkas, K. Tamvakis, JCAP **11**, 028 (2018)
- V.M. Enckell, K. Enqvist, S. Rasanen, L.P. Wahlman, JCAP **02**, 022 (2019)
- A. Salvio, A. Mazumdar, Phys. Lett. B **750**, 194–200 (2015)
- A. Berera, L.Z. Fang, Phys. Rev. Lett. **74**, 1912–1915 (1995)
- A. Berera, Phys. Rev. D **55**, 3346–3357 (1997)
- A. Berera, Nucl. Phys. B **585**, 666–714 (2000)
- A.N. Taylor, A. Berera, Phys. Rev. D **62**, 083517 (2000)
- L.M.H. Hall, I.G. Moss, A. Berera, Phys. Rev. D **69**, 083525 (2004)
- A. Berera, I.G. Moss, R.O. Ramos, Rep. Prog. Phys. **72**, 026901 (2009)
- S. Bartrum, M. Bastero-Gil, A. Berera, R. Cerezo, R.O. Ramos, J.G. Rosa, Phys. Lett. B **732**, 116–121 (2014)
- T. Tenkanen, Phys. Rev. D **99**(6), 063528 (2019)
- J. Yuennan, P. Channuie, D. Momeni, Nucl. Phys. B **1012**, 116810 (2025)
- T. Eadkhong, P. Dam-O, P. Channuie, D. Momeni, Nucl. Phys. B **994**, 116289 (2023)
- A. Payaka, W. Amaek, P. Channuie, Nucl. Phys. B **986**, 116052 (2023)
- W. Amaek, A. Payaka, P. Channuie, Phys. Rev. D **105**(8), 083501 (2022)
- D. Samart, P. Ma-adlerd, P. Koad, P. Channuie, Eur. Phys. J. C **82**(6), 504 (2022)
- D. Samart, P. Ma-adlerd, P. Channuie, Eur. Phys. J. C **82**(2), 122 (2022)
- C. Dioguardi, A. Racioppi, E. Tomberg, JHEP **06**, 106 (2022)
- T.P. Sotiriou, V. Faraoni, Rev. Mod. Phys. **82**, 451–497 (2010)
- I. Antoniadis, A. Karam, A. Lykkas, T. Pappas, K. Tamvakis, JCAP **03**, 005 (2019)
- D.H. Lyth, A. Riotto, Phys. Rep. **314**, 1–146 (1999)
- Y. Akrami et al. [Planck], Astron. Astrophys. **641**, A10 (2020)
- T. Louis et al. [ACT], arXiv:2503.14452 [astro-ph.CO]
- N. Aghanim et al. [Planck], Astron. Astrophys. **641**, A6 (2020) [Erratum: Astron. Astrophys. **652**, C4 (2021)]
- R.L. Workman et al. [Particle Data Group], PTEP **2022**, 083C01 (2022)
- Y. Zhang, JCAP **03**, 023 (2009)
- M. Bastero-Gil, A. Berera, R.O. Ramos, JCAP **07**, 030 (2011)
- G. Panotopoulos, N. Videla, Eur. Phys. J. C **75**(11), 525 (2015)
- M. Benetti, R.O. Ramos, Phys. Rev. D **95**(2), 023517 (2017)
- W. Cheng, X.W. Chen, R. Zhou, J.J. Jiang, X.R. Dai, Z.H. Zhang, T. Qin, Phys. Rev. D **109**(8), 083509 (2024)

50. R. Brandenberger, V. Kamali, R.O. Ramos, *JHEP* **08**, 127 (2020)
51. M. Bastero-Gil, A. Berera, R.O. Ramos, *JCAP* **09**, 033 (2011)
52. M. Bastero-Gil, A. Berera, R.O. Ramos, J.G. Rosa, *JCAP* **01**, 016 (2013)
53. C. Graham, I.G. Moss, *JCAP* **07**, 013 (2009)
54. M. Bastero-Gil, A. Berera, R. Hernández-Jiménez, J.G. Rosa, *Phys. Rev. D* **98**(8), 083502 (2018)
55. M. Bastero-Gil, A. Berera, R.O. Ramos, J.G. Rosa, *Phys. Rev. Lett.* **117**(15), 151301 (2016)
56. M. Bastero-Gil, A. Berera, *Int. J. Mod. Phys. A* **24**, 2207–2240 (2009)
57. R.O. Ramos, L.A. da Silva, *JCAP* **03**, 032 (2013)
58. H.P. De Oliveira, S.E. Joras, *Phys. Rev. D* **64**, 063513 (2001)
59. L. Visinelli, *JCAP* **07**, 054 (2016)
60. K.V. Berghaus, P.W. Graham, D.E. Kaplan, *JCAP* **03**, 034 (2020)
61. M. Laine, S. Procacci, *JCAP* **06**, 031 (2021)
62. M. Motaharfar, R.O. Ramos, *Phys. Rev. D* **104**(4), 043522 (2021)
63. E. Allys et al. [LiteBIRD], *PTEP* **2023**(4), 042F01 (2023)
64. K. Abazajian, G. Addison, P. Adshead, Z. Ahmed, S.W. Allen, D. Alonso, M. Alvarez, A. Anderson, K.S. Arnold, C. Baccigalupi et al., [arXiv:1907.04473](https://arxiv.org/abs/1907.04473) [astro-ph.IM]
65. R. Herrera, M. Olivares, N. Videla, *Phys. Rev. D* **88**, 063535 (2013)