CP-violating inflation and its cosmological imprints

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We study models with several $SU(2)$ scalar doublets where the inert doublets have a nonminimal coupling to gravity and play the role of the inflaton. We allow for this coupling to be complex, thereby introducing *CP* violation—a necessary source of the baryon asymmetry—in the Higgs-inflaton couplings. We investigate the inflationary dynamics of the model and discuss how the *CP* violation of the model is imprinted on the particle asymmetries after inflation in the hot big bang universe.

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I. INTRODUCTION

The Standard Model (SM) of particle physics has been extensively tested and is in great agreement with experimental data, with its last missing particle the Higgs boson discovered by ATLAS and CMS experiments at the CERN Large Hadron Collider (LHC) [1,2]. Although the properties of the observed scalar are in agreement with those of the SM-Higgs boson, it may just be one member of an extended scalar sector. Even though so far no signs of new physics have been detected, it is well understood that the SM of particle physics is incomplete.

Cosmological and astrophysical observations imply a large dark matter (DM) component in the energy budget of the universe. Within the particle physics setting, this would be a particle which is stable on cosmological timescales, cold, nonbaryonic, neutral and weakly interacting [3]. A particle with such characteristics does not exist in the SM. Another shortcoming of the SM is the lack of an explanation for the origin of the observed matter-antimatter asymmetry in the universe. One of the most promising baryogenesis scenarios is electroweak baryogenesis (EWBG) [4], which produces the baryon excess during the electroweak phase transition (EWPT). Although the SM in principle contains all required ingredients for EWBG, it is unable to explain the observed baryon excess due to its insufficient amount of *CP* violation [5–7] and the lack of a first-order phase transition [8].

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Furthermore, in its current form, the SM fails to incorporate cosmic inflation in a satisfactory manner. Inflation is a well-motivated theory predicting a period of exponential expansion in the early universe which explains the generation of primordial density fluctuations seeding structure formation, flatness, homogeneity and isotropy of the universe [9–12]. The simplest models of inflation in best agreement with observations are those driven by a scalar field, the *inflaton*, with a standard kinetic term, slowly rolling down its smooth potential. At the end of inflation, the inflaton which naturally is assumed to have couplings with the SM-Higgs, dumps its energy into the SM bath during the *reheating* process which populates the universe with SM particles.

Scalars with nonminimal couplings to gravity are well-motivated inflaton candidates since they acquire fluctuations proportional to the inflationary scale and can drive the inflation process in the early universe, as in the Higgs-inflation model [13] where the SM-Higgs plays the role of the inflaton, and *s*-inflation models [14,15] where the SM is extended by a singlet scalar. Extensive studies have been carried out in simple one singlet or one doublet scalar extensions of the SM (see e.g., [16–19] and references therein). These models, however, by construction can only partly provide a solution to the main drawbacks of the SM. For example, to incorporate both *CP* violation and DM into the model one has to go beyond simple scalar extensions of the SM [20]; see also e.g., [21–26].

It is therefore theoretically appealing to have a more coherent setting where different motivations of beyond SM frameworks could be simultaneously investigated. For example, in nonminimal Higgs frameworks with conserved discrete symmetries one can accommodate stabilized DM candidates. Moreover, the extended scalar potential could provide new sources of *CP* violation and accommodate a strong first order phase transition [27]. Collider searches can constrain these model frameworks by excluding or discovering the existence of the spectrum of new states.

In this paper we introduce a model where a source of CP violation originates from the couplings of the inflation. Through the process of reheating this is transmitted to an asymmetry within the SM and can furthermore seed the generation of an excess of matter over antimatter during the evolution of the early universe. We describe these dynamics in the context of a Z_2 symmetric 3-Higgs doublet model (3HDM) with a CP-violating extended dark sector, which also provides a viable DM candidate, new sources of CP violation and a strong first-order EWPT [20–25]. We study the inflationary dynamics of this setup and outline its main consequences. In a future work we aim to continue to complement this study by more thorough analysis of EWBG and DM observables as well as a phenomenological analysis towards LHC searches for new physics.

The paper is organized as follows. In Sec. II we present the scalar potential and explore the inflationary dynamics. In Sec. III, we discuss the inflationary imprints of our novel CP violating inflation phenomena. In Sec. IV, we discuss the inflaton decay into the SM particles and possible consequences. In Sec. V we draw our conclusions and discuss the outlook for further work.

II. THE SCALAR POTENTIAL

A. General definitions

A 3HDM scalar potential which is symmetric under a group G of phase rotations can be written as the sum of two parts: V_0 with terms symmetric under any phase rotation, and V_G with terms symmetric under G [28,29]. As a result, a Z_2-symmetric 3HDM can be written as[1]

$$\begin{aligned} V &= V_0 + V_{Z_2},\\ V_0 &= -\mu_1^2(\phi_1^\dagger\phi_1) - \mu_2^2(\phi_2^\dagger\phi_2) - \mu_3^2(\phi_3^\dagger\phi_3) + \lambda_{11}(\phi_1^\dagger\phi_1)^2 + \lambda_{22}(\phi_2^\dagger\phi_2)^2 + \lambda_{33}(\phi_3^\dagger\phi_3)^2 + \lambda_{12}(\phi_1^\dagger\phi_1)(\phi_2^\dagger\phi_2)\\ &\quad + \lambda_{23}(\phi_2^\dagger\phi_2)(\phi_3^\dagger\phi_3) + \lambda_{31}(\phi_3^\dagger\phi_3)(\phi_1^\dagger\phi_1) + \lambda'_{12}(\phi_1^\dagger\phi_2)(\phi_2^\dagger\phi_1) + \lambda'_{23}(\phi_2^\dagger\phi_3)(\phi_3^\dagger\phi_2) + \lambda'_{31}(\phi_3^\dagger\phi_1)(\phi_1^\dagger\phi_3),\\ V_{Z_2} &= -\mu_{12}^2(\phi_1^\dagger\phi_2) + \lambda_1(\phi_1^\dagger\phi_2)^2 + \lambda_2(\phi_2^\dagger\phi_3)^2 + \lambda_3(\phi_3^\dagger\phi_1)^2 + \text{H.c.}, \end{aligned} \tag{1}$$

where the three Higgs doublets, ϕ_1, ϕ_2, ϕ_3, transform under the Z_2 group, respectively, as

$$g_{Z_2} = \mathrm{diag}(-1, -1, +1). \tag{2}$$

The parameters of the V_0 part of the potential are real by construction. We allow for the parameters of V_{Z_2} to be complex, using the following notation throughout the paper:

$$\lambda_j = |\lambda_j| e^{i\theta_j} \quad (j = 1, 2, 3), \quad \text{and} \quad \mu_{12}^2 = |\mu_{12}^2| e^{i\theta_{12}}. \tag{3}$$

The composition of the doublets is as follows:

$$\phi_1 = \begin{pmatrix} H_1^+ \\ \frac{H_1 + iA_1}{\sqrt{2}} \end{pmatrix}, \qquad \phi_2 = \begin{pmatrix} H_2^+ \\ \frac{H_2 + iA_2}{\sqrt{2}} \end{pmatrix},$$
$$\phi_3 = \begin{pmatrix} G^+ \\ \frac{v + h + iG^0}{\sqrt{2}} \end{pmatrix}, \tag{4}$$

where ϕ_1 and ϕ_2 are the Z_2-odd *inert* doublets, $\langle\phi_1\rangle = \langle\phi_2\rangle = 0$, and ϕ_3 is the one Z_2-even *active* doublet, which at low energy attains a vacuum expectation value (VEV) $\langle\phi_3\rangle = v/\sqrt{2} \neq 0$. The doublet ϕ_3 plays the role of the SM Higgs doublet, with h being the SM Higgs boson and $G^\pm$, G^0 the would-be Goldstone bosons. Note that according to the Z_2 generator in Eq. (2) the symmetry of the potential is respected by the vacuum $(0, 0, v/\sqrt{2})$. In this paper we consider the scenario where the components of the inert doublets act as inflation candidates and reheat the universe at the end of inflation through their interactions with the SM-Higgs and gauge bosons. Note that at the scales relevant for inflation we can take the VEV of the active doublet to be zero, $\langle\phi_3\rangle = 0$.

Furthermore, CP violation is only introduced in the *inert* sector which is forbidden from mixing with the *active* sector by the conservation of the Z_2 symmetry. As a result, the amount of CP violation is not limited by electric dipole moments [21]. The lightest particle amongst the CP-mixed neutral fields from the inert doublets is a viable DM candidate and stable due to the unbroken Z_2 symmetry. In this paper, we focus on the inflationary dynamics of the model and shall not discuss DM implications of the model any further.

B. Potential for the inflaton

We start by rewriting the doublets in the unitary gauge and ignore the charged scalars (since they do not affect the inflationary dynamics):

$$\phi_1 = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ h_1 + i\eta_1 \end{pmatrix}, \qquad \phi_2 = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ h_2 + i\eta_2 \end{pmatrix},$$
$$\phi_3 = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ h_3 \end{pmatrix}. \tag{5}$$

[1] We ignore additional Z_2-symmetric terms that can be added to the potential, e.g., $(\phi_3^\dagger\phi_1)(\phi_2^\dagger\phi_3)$, $(\phi_1^\dagger\phi_2)(\phi_3^\dagger\phi_3)$, $(\phi_1^\dagger\phi_2)(\phi_1^\dagger\phi_1)$ and $(\phi_1^\dagger\phi_2)(\phi_2^\dagger\phi_2)$, as they do not change the phenomenology of the model [23].

The action of the model in the Jordan frame is

$$
\begin{aligned}
S_J = \int d^4x\sqrt{-g}\Big[&-\frac{1}{2}M_{\rm pl}^2 R - D_\mu\phi_1^\dagger D^\mu\phi_1 \\
&- D_\mu\phi_2^\dagger D^\mu\phi_2 - D_\mu\phi_3^\dagger D^\mu\phi_3 \\
&- V(\phi_1,\phi_2,\phi_3) - (\xi_1|\phi_1|^2 + \xi_2|\phi_2|^2 \\
&+ \xi_3|\phi_3|^2 + \xi_4(\phi_1^\dagger\phi_2) + \xi_4^*(\phi_2^\dagger\phi_1))R\Big], \qquad (6)
\end{aligned}
$$

where R is the Ricci scalar, $M_{\rm pl}$ is the reduced Planck mass and the parameters ξ_i are dimensionless couplings of the scalar doublets to gravity. Note that, in principle, ξ_4 could be a complex parameter for which we use the notation

$$\xi_4 = |\xi_4|e^{i\theta_4}. \qquad (7)$$

In Eq. (6) the covariant derivative, D_μ, contains couplings of the scalars with the gauge bosons. However, for the dynamics during the inflation, the covariant derivative is reduced to the normal derivative $D_\mu \to \partial_\mu$. The minus sign in the kinetic terms follows the metric convention of $(-,+,+,+)$.

Since we identify the two inert doublets with inflaton, we assume that the energy density of ϕ_3 is subdominant during inflation. Therefore, the part of the potential relevant for inflation is

$$
\begin{aligned}
V = &-\mu_1^2(\phi_1^\dagger\phi_1) - \mu_2^2(\phi_2^\dagger\phi_2) \\
&+ \lambda_{11}(\phi_1^\dagger\phi_1)^2 + \lambda_{22}(\phi_2^\dagger\phi_2)^2 \\
&+ \lambda_{12}(\phi_1^\dagger\phi_1)(\phi_2^\dagger\phi_2) + \lambda'_{12}(\phi_1^\dagger\phi_2)(\phi_2^\dagger\phi_1) \\
&- \mu_{12}^2(\phi_1^\dagger\phi_2) + \lambda_1(\phi_1^\dagger\phi_2)^2 + \text{H.c.} \qquad (8)
\end{aligned}
$$

Due to local SU(2) invariance, we can rotate away one of the CP-odd fields, say η_2. Such a transformation is equivalent to taking the $\eta_2 \to 0$ limit, and we assume this limit to be taken when writing the fields in terms of components in Eq. (5).

To facilitate the analysis, we apply a conformal transformation from the Jordan frame, which contains terms with scalar-gravity quadratic couplings, to the Einstein frame with no explicit couplings to gravity [30]. Physical observables are invariant under this frame transformation. The two frames are equivalent after the end of inflation when the transformation parameter equals unity.

The action in the Einstein frame can be written as

$$S_E = \int d^4x\sqrt{-\tilde{g}}\Big[-\frac{1}{2}M_{\rm pl}^2\tilde{R} - \frac{1}{2}\tilde{g}^{\mu\nu}G_{ij}\partial_\mu\varphi_i\partial_\nu\varphi_j - \tilde{V}\Big], \qquad (9)$$

where $\tilde{V} = V/\Omega^4$ is the potential in the Einstein frame following the conformal transformation

$$
\begin{aligned}
\tilde{g}_{\mu\nu} &= \Omega^2 g_{\mu\nu}, \\
G_{ij} &= \frac{1}{\Omega^2}\delta_{ij} + \frac{3}{2}\frac{M_{\rm pl}^2}{\Omega^4}\frac{\partial\Omega^2}{\partial\varphi_i}\frac{\partial\Omega^2}{\partial\varphi_j}, \qquad (10)
\end{aligned}
$$

where $\varphi_k = h_1, h_2, \eta_1$, and the transformation parameter

$$
\begin{aligned}
\Omega^2 = 1 &+ \frac{\xi_1}{M_{\rm pl}^2}(h_1^2 + \eta_1^2) + \frac{\xi_2}{M_{\rm pl}^2}h_2^2 \\
&+ \frac{2|\xi_4|}{M_{\rm pl}^2}(h_1h_2c_{\theta_4} + \eta_1h_2s_{\theta_4}) \qquad (11)
\end{aligned}
$$

using the shorthand notation $c_{\theta_k} = \cos\theta_k$ and $s_{\theta_k} = \sin\theta_k$ throughout the paper.

The prefactor G_{ij} in Eq. (10) leads to mixed kinetic terms. We introduce the reparametrization

$$A = \sqrt{\frac{3}{2}}M_{\rm pl}\log(\Omega^2) \quad \text{with} \quad \frac{\partial\Omega^2}{\partial\varphi_k} = \sqrt{\frac{2}{3}}\frac{\Omega^2}{M_{\rm pl}}\frac{dA}{d\varphi_k} \qquad (12)$$

which reduces the kinetic terms to the diagonal form:

$$
\begin{aligned}
\tilde{g}_{\mu\nu}G_{ij}\partial_\mu\varphi_i\partial_\nu\varphi_j &= \Omega^2 g_{\mu\nu}\left(\frac{\delta_{ij}}{\Omega^2} + \frac{\partial A}{\partial\varphi_i}\frac{\partial A}{\partial\varphi_j}\right)\partial_\mu\varphi_i\partial_\nu\varphi_j \\
&= \partial_\mu\varphi_i\partial_\mu\varphi_i + \Omega^2\partial_\mu A\partial_\mu A. \qquad (13)
\end{aligned}
$$

To write the potential in the Einstein frame, we keep only terms in the potential in Eq. (8) which are quartic in $h_{1,2}$ and η_1. This reduces the potential to

$$
\begin{aligned}
\tilde{V} \approx \frac{1}{4\Omega^4}[&\lambda_{11}(h_1^2 + \eta_1^2)^2 + \lambda_{22}h_2^4 + (\lambda_{12} + \lambda'_{12})(h_1^2 + \eta_1^2)h_2^2 \\
&+ 2|\lambda_1|(c_{\theta_1}(h_2^2(h_1^2 - \eta_1^2)) + 2s_{\theta_1}h_2^2h_1\eta_1)], \qquad (14)
\end{aligned}
$$

where θ_1 is the CP-violating phase of the λ_1 parameter.

Further, we introduce another reparametrization,

$$\eta_1 = \beta_1 h_1, \qquad h_2 = \beta_2 h_1, \qquad (15)$$

with β_1, β_2 as field dependent values, to rewrite the potential as

$$
\begin{aligned}
\tilde{V} \approx \frac{h_1^4}{4\Omega^4}[&\lambda_{11}(1 + \beta_1^2)^2 + \lambda_{22}\beta_2^4 + ((\lambda_{12} + \lambda'_{12})(1 + \beta_1^2) \\
&+ 2|\lambda_1|(c_{\theta_1}(1 - \beta_1^2) + 2s_{\theta_1}\beta_1))\beta_2^2]. \qquad (16)
\end{aligned}
$$

Using this reparametrization, one can also simplify the Ω^2 parameter in Eq. (11) as

$$\Omega^2 = 1 + \left(\frac{\xi_1}{M_{\rm pl}^2}(1+\beta_1^2) + \frac{\xi_2}{M_{\rm pl}^2}\beta_2^2 + \frac{2|\xi_4|}{M_{\rm pl}^2}\beta_2(c_{\theta_4} + \beta_1 s_{\theta_4})\right)$$

$$h_1^2 \equiv 1 + \frac{B}{M_{\rm pl}^2} h_1^2. \tag{17}$$

From Eq. (12), recall that $\Omega^2 = \exp(\tilde{A})$ using the shorthand notation $\tilde{A} = \sqrt{\frac{2}{3}}\frac{A}{M_{\rm pl}}$. One can then write the field h_1 in terms of the reparametrized field $\tilde{A}$:

$$h_1^2 = \frac{M_{\rm pl}^2}{B}(e^{\tilde{A}} - 1). \tag{18}$$

Therefore, expressing h_1^2 and Ω^2 in terms of $\tilde{A}$ allows us to write the potential in Eq. (16) in the form

$$\tilde{V} \sim (1 - e^{-\tilde{A}})^2 X(\beta_1, \beta_2). \tag{19}$$

We will be interested in the effect of the nonminimal coupling ξ_4 and the associated phase θ_4. Therefore, we will set $\xi_1 = \xi_2 = 0$ and assume that the initial field values are such that $\Omega^2 > 0$ is guaranteed. Therefore, with these assumptions, the potential in Eq. (16) can be written as

$$\tilde{V} = \left(\frac{M_{\rm pl}^2}{2|\xi_4|}\right)^2 (1 - e^{-\tilde{A}})^2 X(\beta_1, \beta_2), \tag{20}$$

where

$$X(\beta_1, \beta_2) = \frac{\lambda_{11}(1+\beta_1^2)^2 + \lambda_{22}\beta_2^4 + ((\lambda_{12} + \lambda'_{12})(1+\beta_1^2) + 2|\lambda_1|(c_{\theta_1}(1-\beta_1^2) + 2 s_{\theta_1}\beta_1))\beta_2^2}{4\beta_2^2(c_{\theta_4} + \beta_1 s_{\theta_4})^2}. \tag{21}$$

Following the procedure in [16], to find the direction of inflation, we first minimize the $X(\beta_1, \beta_2)$ function with respect to β_2 which occurs at

$$\frac{\partial X(\beta_1, \beta_2)}{\partial \beta_2} = 0 \Rightarrow \beta_2^2 = \sqrt{\frac{\lambda_{11}}{\lambda_{22}}}(1+\beta_1^2). \tag{22}$$

The second order derivative at this point is

$$\frac{\partial^2 X(\beta_1, \beta_2)}{\partial \beta_2^2} = \frac{2\lambda_{22}}{(c_{\theta_4} + \beta_1 s_{\theta_4})^2} \tag{23}$$

which is always positive provided $\lambda_{22} > 0$, as shown in the left panel in Fig. 1.

Using the β_2 value in Eq. (22), we can write the $X(\beta_1, \beta_2)$ function solely in terms of β_1,

$$X(\beta_1) = \frac{(1+\beta_1^2)\Lambda + 2((1-\beta_1^2)c_{\theta_1} + 2\beta_1 s_{\theta_1})|\lambda_1|}{4(c_{\theta_4} + \beta_1 s_{\theta_4})^2} \tag{24}$$

with $\Lambda = \lambda_{12} + \lambda'_{12} + 2\sqrt{\lambda_{11}\lambda_{22}}$.

We repeat the same treatment and minimize the $X(\beta_1)$ function with respect to β_1.

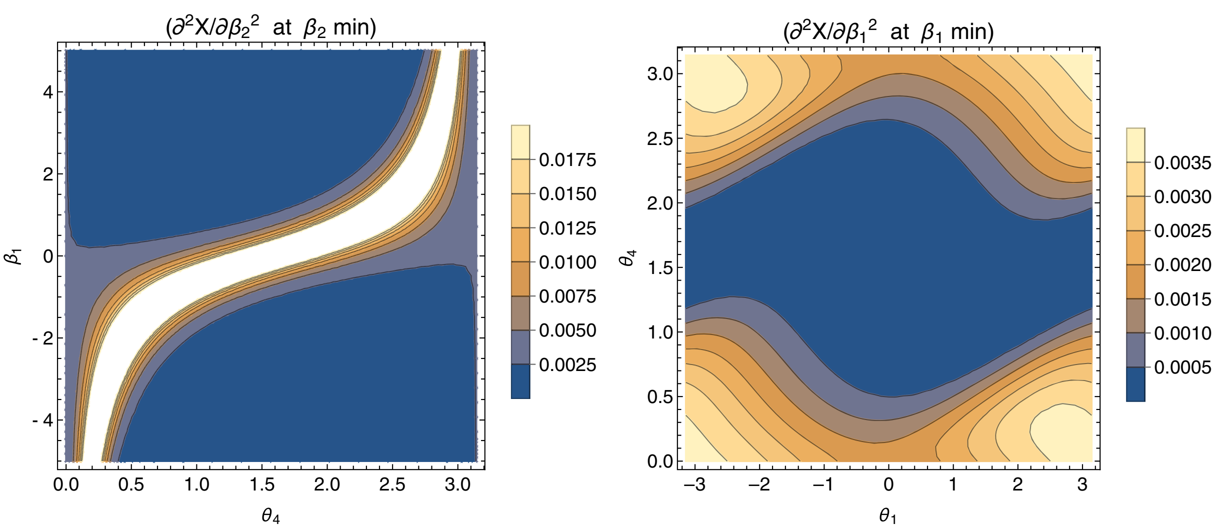

FIG. 1. The second order derivative of the function $X(\beta_1, \beta_2)$ with respect to β_2 at the minimum ($\partial X/\partial\beta_2 = 0$) on the left and the second order derivative of the function $X(\beta_1)$ with respect to β_1 at the minimum ($\partial X/\partial\beta_1 = 0$) on the right (all $\lambda_i \sim 0.001$). The white area on the left panel corresponds to where the denominator in Eq. (23) becomes zero.

$$\frac{\partial X(\beta_1)}{\partial \beta_1} = 0 \Rightarrow \beta_1 = \frac{(\Lambda + 2|\lambda_1|c_{\theta_1})s_{\theta_4} - 2|\lambda_1|c_{\theta_4}s_{\theta_1}}{(\Lambda - 2|\lambda_1|c_{\theta_1})c_{\theta_4} - 2|\lambda_1|s_{\theta_4}s_{\theta_1}}. \tag{25}$$

We check the positivity of the second order derivative at the minimum point which is satisfied for all θ_1, θ_4 values as shown in the right panel of Fig. 1.

Replacing the β_1 value which minimizes the $X(\beta_1)$ function back into the $X(\beta_1)$ function itself, yields the form of X independent of β_1 and β_2 with only θ_1 and θ_4 as variables:

$$X(\theta_1, \theta_4) = \frac{\frac{1}{4}\Lambda^2 - \lambda_1^2}{\Lambda - 2\lambda_1 \cos(\theta_1 - 2\theta_4)}. \tag{26}$$

The left panel in Fig. 2 shows the $X(\theta_1, \theta_4)$ function for allowed values of θ_1 and θ_4. At each point in the plots, one can derive the values of β_1 and consequently β_2 using Eq. (22) for given values of θ_1 and θ_4. The right panel in Fig. 2 shows the values of β_1 for varying values of θ_1 and θ_4.

III. INFLATIONARY DYNAMICS

With the procedure used in the previous section, the dynamics is essentially that of a single field inflation. The full inflationary potential in Eq. (20) can be written as

$$\tilde{V} = \left(\frac{M_{\rm pl}^2}{2|\xi_4|}\right)^2 (1 - e^{-\tilde{A}})^2 X(\theta_1, \theta_4). \tag{27}$$

Figure 3 shows the inflationary potential for different values of θ_1 and θ_4. Note that the potential is almost flat at high field values which ensures a slow roll inflation.

For the usual slow roll parameters in this case the function X is irrelevant, since it cancels in the expressions for ϵ and η, which are

$$\epsilon = \frac{1}{2} M_{\rm pl}^2 \left(\frac{1}{\tilde{V}} \frac{d\tilde{V}}{dA}\right)^2 = \frac{4}{3(1 - e^{\tilde{A}})^2}, \tag{28}$$

$$\eta = M_{\rm pl}^2 \frac{1}{\tilde{V}} \frac{d^2\tilde{V}}{dA^2} = \frac{4(2 - e^{\tilde{A}})}{3(1 - e^{\tilde{A}})^2}. \tag{29}$$

For field values $A \gg M_{\rm pl}$ (or equivalently $\tilde{A} \gg 1$), both parameters $\epsilon, \eta \ll 1$ which satisfies the slow roll condition. Inflation ends when $\epsilon \simeq 1$. To calculate the values of A at the beginning and end of inflation, A_i and A_f respectively, one needs to calculate the number of e-folds N_e, i.e., the number of times the universe expanded by e times its own size. N_e is calculated to be

$$N_e = \frac{1}{M_{\rm pl}^2} \int_{A_f}^{A_i} \frac{\tilde{V}}{\tilde{V}'} dA = \frac{3}{4}[\tilde{A}_f - \tilde{A}_i - e^{\tilde{A}_f} + e^{\tilde{A}_i}], \tag{30}$$

where $\tilde{V}' = \frac{d\tilde{V}}{dA}$ and A_i ($\tilde{A}_i$) is the value of A ($\tilde{A}$) at the beginning of inflation and A_f ($\tilde{A}_f$) is the value of A ($\tilde{A}$) at the end of the inflation. Since inflation ends when $\epsilon \simeq 1$, one can calculate A_f, which yields

$$e^{\tilde{A}_f} = \exp\left(\sqrt{\frac{2}{3}} \frac{A_f}{M_{\rm pl}}\right) \simeq 2.1547$$

$$\Rightarrow \tilde{A}_f = \sqrt{\frac{2}{3}} \frac{A_f}{M_{\rm pl}} \simeq 0.7676. \tag{31}$$

To calculate A_i, one could plug in the A_f value into Eq. (30) assuming $N_e = 60$, which results in

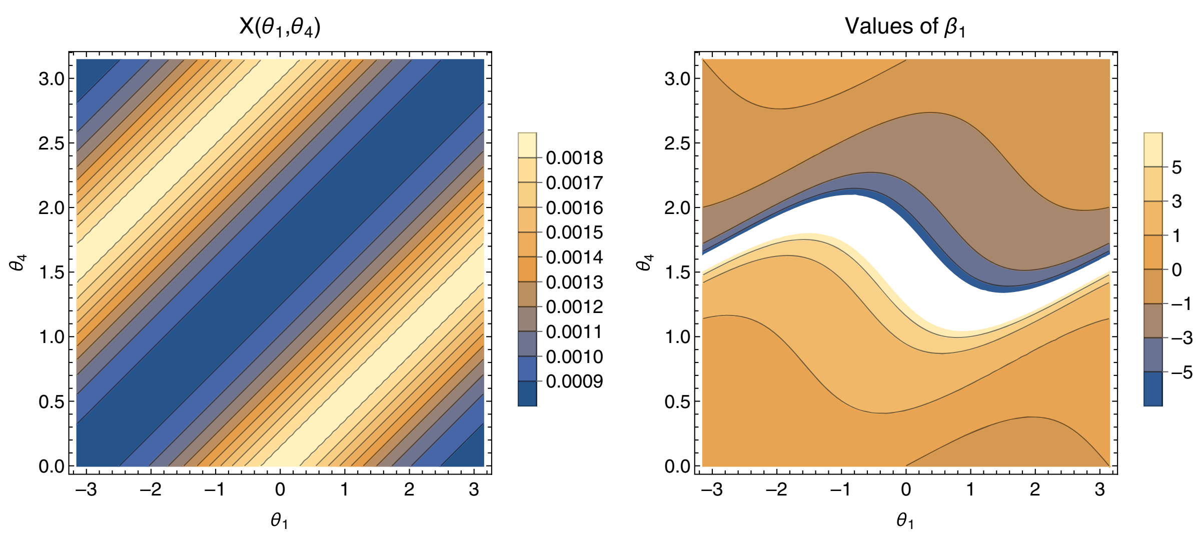

FIG. 2. The $X(\theta_1, \theta_4)$ function on the left and the values of β_1 on the right for varying values of θ_1 and θ_4 (all $\lambda_i \sim 0.001$). The white region in the right panel shows a discontinuity where β_1 values tend to plus infinity approaching from the bottom and to minus infinity approaching from the top of the plot.

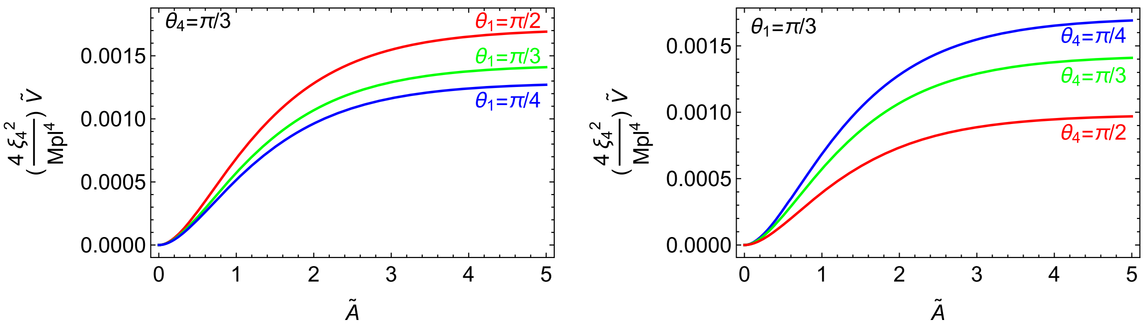

FIG. 3. The inflationary potential for different values of θ_1 and θ_4 (all $\lambda_i \sim 0.001$).

$$\frac{3}{4}[-\tilde{A}_i + e^{\tilde{A}_i}] - 1.0403 = 60,$$

$$\Rightarrow \tilde{A}_i = \sqrt{\frac{2}{3}}\frac{A_i}{M_{\rm pl}} \approx 4.4524. \quad (32)$$

At this point we can also check the field values in terms of the original field h_1 using Eq. (18). This gives

$$h_{1f} = \frac{1.85\times 10^{18}}{\sqrt{|\xi_4|\beta_2(c_{\theta_4}+\beta_1 s_{\theta_4})}},$$
$$h_{1i} = \frac{1.59\times 10^{19}}{\sqrt{|\xi_4|\beta_2(c_{\theta_4}+\beta_1 s_{\theta_4})}}. \quad (33)$$

In the case of Higgs inflation where the nonminimal coupling to gravity, ξ, is forced to be of the order $\sim 10^4$ GeV, the h field values during inflation are as large as 10^{16} GeV or so. In our case the situation is similar.

Having fixed N_e to 60, and calculated the A field value at the start of inflation, we can derive the scalar power spectrum, P_s, the tensor to scalar ratio r and the spectral index n_s as follows:

$$P_s = \frac{1}{12\pi^2 M_{\rm pl}^6}\frac{(\tilde{V})^3}{(\tilde{V}')^2} = \left(\frac{(1-e^{\tilde{A}})^4}{128\pi^2 e^{2\tilde{A}}}\right)\frac{X(\theta_1,\theta_4)}{|\xi_4|^2} = 5.565\times\frac{X(\theta_1,\theta_4)}{|\xi_4|^2}, \quad (34)$$

$$r = 16\epsilon = 0.00296, \quad (35)$$

$$n_s = 1 - 6\epsilon + 2\eta = 0.9678, \quad (36)$$

where $\tilde{V}'$ is the derivative of $\tilde{V}$ with respect to A and both $\tilde{V}$ and $\tilde{V}'$ are calculated at the A_i. Figure 4 shows the slow roll parameters N_e, n_s and r with respect to $\tilde{A}$ with the grid lines highlighting the $55 < N_e < 65$ values. We show the inflationary parameters over a range of N_e, since there is no reason for N_e to be precisely 60. The values of r and n_s are well within the Plank bounds of $n_s = 0.9677 \pm 0.0060$ at the 1σ level and $r < 0.11$ at 95% confidence level [31]. Note that the spectral index and the tensor to scalar ratio are in agreement with the Planck bounds over the full range of N_e. Figure 5 shows the 1σ and 2σ regions allowed by Planck observations in the $r - n_s$ plane and the theoretical predictions of our framework for N_e values of 55 and 65.

Observations from WMAP7 [32] constrain the scalar power spectrum which put a bound on the ξ_4 coupling and angles θ_1, θ_4,

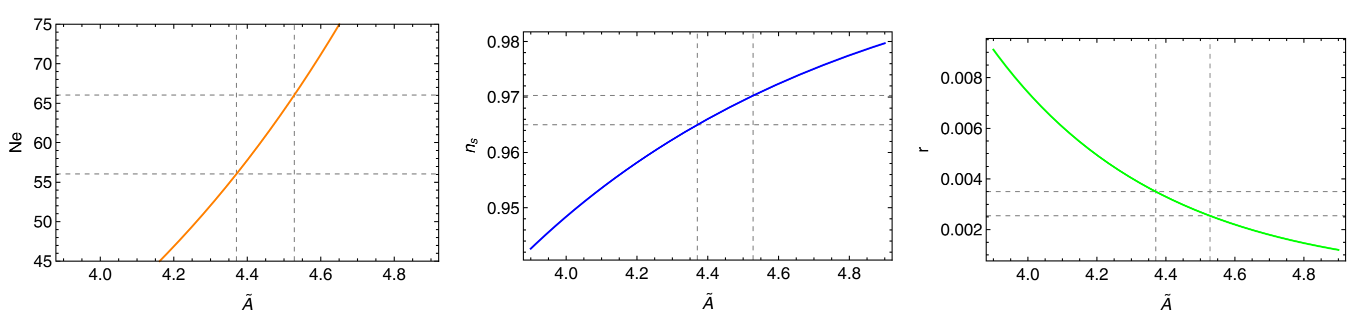

FIG. 4. The slow roll parameters: the number of e-folds N_e (left), spectral index n_s (center) and tensor to scalar ratio r (right) as a function of $\tilde{A}$ with the grid lines highlighting the $55 < N_e < 65$ values.

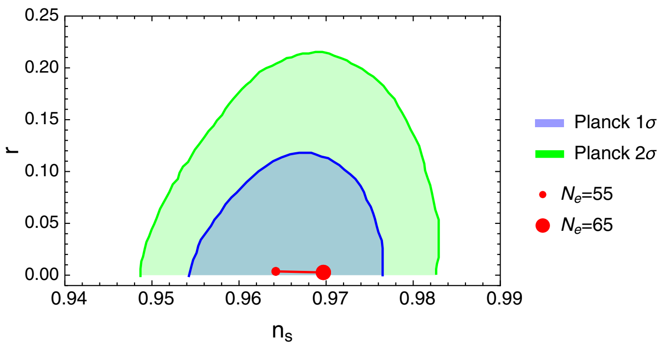

FIG. 5. The 1σ and 2σ regions for n_s and r from Planck observation compared to the theoretical prediction of our framework.

$$P_s = (2.430 \pm 0.091) \times 10^{-9} = 5.565 \times \frac{X(\theta_1, \theta_4)}{|\xi_4|^2}. \quad (37)$$

In the left panel of Fig. 6, we show P_s values for the fixed $\theta_1 = \pi/3$ angle and varying values of ξ_4 and θ_4 up to 3σ standard deviation from the central value in Eq. (37). In the right panel, we fix P_s to the WMAP7 central value for fixed values of $\lambda_i \sim 0.001$ to get

$$|\xi_4| = 4.785 \times 10^4 \sqrt{X(\theta_1, \theta_4)} \quad (38)$$

and show contours of ξ_4 for varying values of θ_1 and θ_4. Note that every point in the plot yields the exact P_s central value.

This is a very important feature of our framework. To satisfy the bounds on the scalar power spectrum, the function $X(\theta_1, \theta_4)$ allows for a wide range of ξ_4 values as shown in Fig. 6. This is in contrast to the Higgs-inflation models where $P_s \propto \lambda/\xi^2$ with λ the Higgs self-coupling which is fixed to be ~ 0.12 at the electroweak scale. Thus, for P_s to agree with observations at the inflationary scale, ξ will have to be very large $\mathcal{O}(10^4)$. In our setup, a combination of parameters λ_1, λ_{11}, λ_{22}, λ_{12}, λ'_{12} appears in the $X(\theta_1, \theta_4)$ function. The only constraint limiting these parameters is the stability of the potential requiring

$$\lambda_{ii} > 0, \qquad \lambda_{ij} + \lambda'_{ij} > -2\sqrt{\lambda_{ii}\lambda_{jj}},$$
$$|\lambda_i| \le |\lambda_{ii}|, |\lambda_{ij}|, |\lambda'_{ij}|, \qquad i \neq j = 1, 2, 3, \quad (39)$$

which allows for very small values of $\lambda_i \sim 0.001$ which, in turn, allows for much smaller values of ξ_4, at least 1 order of magnitude than the ξ value in Higgs-inflation models.

IV. REHEATING AND SCALAR ASYMMETRIES

At the end of inflation, the energy stored in the inflaton disperses as the inflaton decays/annihilates into the SM particles through processes mediated by the SM-Higgs and gauge bosons in our case, during the so-called reheating phase [33]. There are numerous details on how the inflaton decays and creates the initial condition for the conventional hot early universe. Here our main interest is to discuss how the *CP* asymmetry originating from the nonminimal coupling is transferred to the SM degrees of freedom.

For the discussion of the scalar asymmetries, let us focus on the neutral components of the ϕ_1 doublets acquiring an initial nonvanishing expectation value at the exit from inflaton. We write the field fluctuations around the initial conditions as

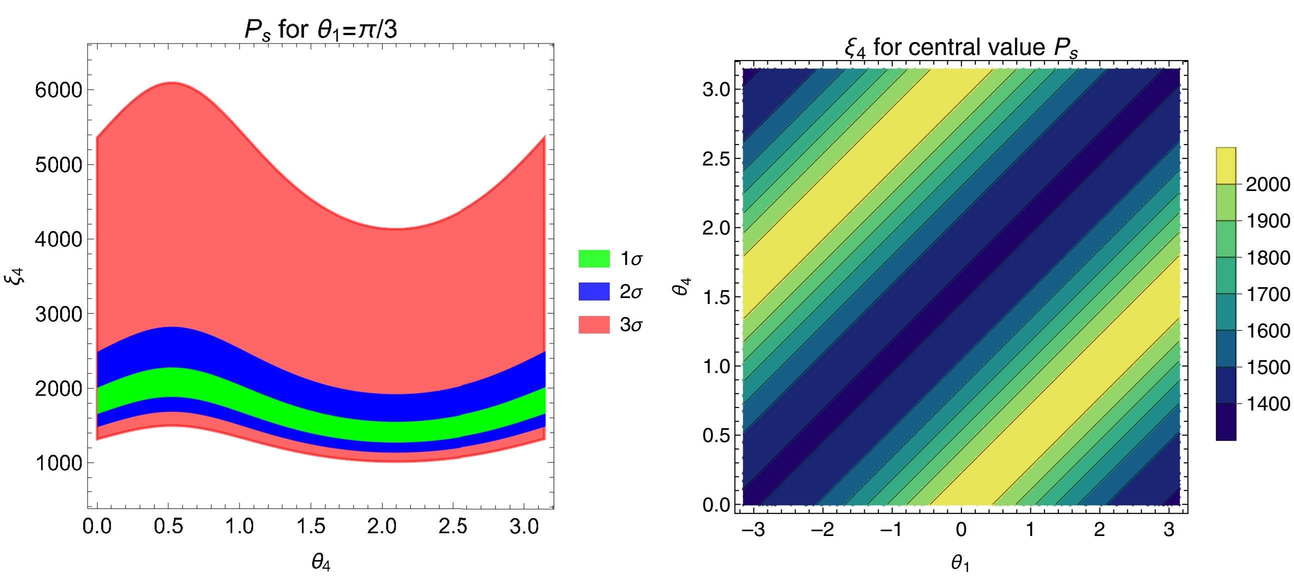

FIG. 6. Left panel: P_s values for the fixed $\theta_1 = \pi/3$ angle and varying values of ξ_4 and θ_4 up to 3σ standard deviation from the observed central value. Right panel: contours of ξ_4 in the $\theta_1 - \theta_4$ plane which lead to P_s central values (all $\lambda_i \sim 0.001$).

$$\begin{aligned} &\phi_1 \to \phi_1 - a_1 e^{i\alpha}, \qquad \phi_1^\dagger \to \phi_1^* - a_1 e^{-i\alpha} \\ &\phi_2 \to \phi_2 - a_2, \qquad \phi_2^\dagger \to \phi_2^* - a_2 \\ &\phi_3 \to \phi_3 - a_3, \qquad \phi_3^\dagger \to \phi_3^* - a_3. \end{aligned} \tag{40}$$

The phase α here is related to the CP-violating phases of inflation. Note that at the end of inflation the h_1 field has taken the value h_{1f} according to Eq. (33) which is dependent on the inflationary dynamics, namely θ_4, β_1 and β_2 which are dependent on θ_1. Since h_1 is the real part of the complex field ϕ_1, its value is what feeds the $a_1 \cos\alpha$ component of fluctuations in Eq. (40). The imaginary part of ϕ_1, represented by η_1, takes a value proportional to $\eta_{1f} = \beta_1 h_{1f}$ as shown in Eq. (15), and feeds the $a_1 \sin\alpha$ component of the field fluctuations. Recall that one can obtain the values of β_1 and β_2 for any given value of θ_1 and θ_4 from Eqs. (22) and (25). Explicitly, one can write

$$\begin{aligned} \tan\alpha &= \frac{a_1 \sin\alpha}{a_1 \cos\alpha} = \frac{\eta_{1f}}{h_{1f}} = \beta_1 \\ &= \frac{(\Lambda + 2|\lambda_1| c_{\theta_1}) s_{\theta_4} - 2|\lambda_1| c_{\theta_4} s_{\theta_1}}{(\Lambda - 2|\lambda_1| c_{\theta_1}) c_{\theta_4} - 2|\lambda_1| s_{\theta_4} s_{\theta_1}}, \end{aligned} \tag{41}$$

with $\Lambda = \lambda_{12} + \lambda'_{12} + 2\sqrt{\lambda_{11}\lambda_{22}}$ as mentioned before. However, to keep the present discussion more transparent, we retain a generic phase α here.

To discuss the consequences of this complex phase, we now assume instant reheating. Since the field ϕ_3 is light with respect to the inflaton degrees of freedom, we expect the latter to quickly decay to ϕ_3. The asymmetry arising from the values of the fields in Eq. (40) will manifest in creation of an unequal number of ϕ_3 and ϕ_3^* quanta as follows.

Let us study the decay process $\phi_1 \to \phi_3^* \phi_3^*$ in detail. From the potential in Eq. (1), the amplitude of the tree-level process is proportional to

$$\begin{aligned} &\mathcal{M}_{(\phi_1 \to \phi_3^* \phi_3^*)} \propto -2 a_1 \lambda_3 e^{i(\alpha+\theta_3)} \quad \text{and} \\ &\mathcal{M}_{(\phi_1^* \to \phi_3 \phi_3)} \propto -2 a_1 \lambda_3 e^{-i(\alpha+\theta_3)}. \end{aligned} \tag{42}$$

The generation of the asymmetry is sensitive to the interference between the tree and loop diagrams [34,35]. Hence, we need to sketch what happens at loop level. At one loop level, there are many diagrams that contribute to this decay process. For the purpose of demonstration, we consider the bubble diagrams which convert ϕ_1 to ϕ_3 with only ϕ_1 and ϕ_1^* in the loop, as shown in Fig. 7. Clearly one needs to take into account all diagrams contributing to this decay process, especially since there may be interferences canceling the CP asymmetry. However, since all triple scalar couplings in the potential can be different, one can ensure that such cancellation does not occur. More careful analysis of these effects is deferred to a future work.

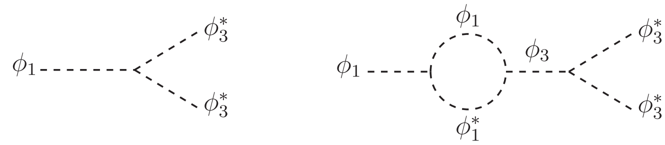

FIG. 7. The tree-level decay process $\phi_1 \to \phi_3^* \phi_3^*$ and the one-loop bubble diagram with ϕ_1 and ϕ_1^* in the loop.

The amplitude of the loop process with ϕ_1 and ϕ_1^* running in the loop is proportional to

$$\mathcal{M}_{(\phi_1 \to \phi_3 \to \phi_3^* \phi_3^*)} \propto -4 a_1 a_3^2 \lambda_{11} \lambda_{33} (\lambda_{31} + \lambda'_{31}) e^{-i\alpha}, \tag{43}$$

$$\mathcal{M}_{(\phi_1^* \to \phi_3^* \to \phi_3 \phi_3)} \propto -4 a_1 a_3^2 \lambda_{11} \lambda_{33} (\lambda_{31} + \lambda'_{31}) e^{i\alpha}. \tag{44}$$

Due to the interference of the tree and loop diagrams, the decay processes are CP violating and result in an unequal number of ϕ_3 and ϕ_3^* states. Consequently, we define the asymmetry A_{CP}^1 as the difference between the ϕ_1 decay rate and its conjugate, and we find

$$\begin{aligned} A_{CP}^1 &= \Gamma^{\text{tree+loop}}_{(\phi_1 \to \phi_3^* \phi_3^*)} - \Gamma^{\text{tree+loop}}_{(\phi_1^* \to \phi_3 \phi_3)} \\ &= -\frac{1}{16\sqrt{3}\pi^2} a_1^2 a_3^2 \lambda_3 \lambda_{11} \lambda_{33} (\lambda_{31} + \lambda'_{31}) \sin(2\alpha + \theta_3). \end{aligned} \tag{45}$$

This asymmetry in the scalar sector is then transferred to the fermion sector through the couplings of the Higgs field (the ϕ_3 doublet) with the fermions. For example, assuming the existence of right-handed neutrinos, the Yukawa interactions between neutrinos and ϕ_3 will generate an asymmetry between ν_L and $\bar{\nu}_R$, which would be further translated into baryon asymmetry by the electroweak sphalerons.

V. CONCLUSION AND OUTLOOK

Scalar fields which have nonminimal couplings to gravity are well-motivated inflaton candidates. Paradigmatic examples are the Higgs-inflation [13] and s-inflation models [15]. In this paper we have considered a scenario where several nonminimally coupled scalars contribute to the inflationary dynamics. In particular we investigated a model where these scalars are electroweak doublets and therefore generalize the Higgs inflation. We focused on a setting where the dominant nonminimal coupling is allowed to be complex and investigated the effect that this would have on CP violation in our universe. We determined the inflationary dynamics in the regime where the model essentially conforms to the predictions of single field inflation. The essential difference is that the inflaton obtains a nonzero phase representing a possible source of CP violation for subsequent postinflationary evolution. At the end of inflation, the inflaton particle which is naturally assumed to have couplings with the SM Higgs, dumps its energy into the SM particle bath through the process of reheating, which populates the universe with the SM particles. We sketched how the complex value of the inflaton field leads to an asymmetry in the scalar sector decays, and how this asymmetry will further be transmitted to the fermion sector. There are numerous details in our scenario which can be investigated in more detail. These include the multifield dynamics during the inflation as well as the details of reheating and subsequent particle decays. Also the detailed analysis of the effects on the generation of baryon asymmetry needs to be addressed in more detail. We will consider these in future work on the model introduced in this paper.

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