PHYSICAL REVIEW D **92,** 121305(R) (2015)

Kinetically modified nonminimal Higgs inflation in supergravity

Constantinos Pallis*

Department of Physics, University of Cyprus, P.O. Box 20537, Nicosia 1678, Cyprus

(Received 18 August 2015; published 30 December 2015)

We consider models of chaotic inflation driven by the real parts of a conjugate pair of Higgs superfields involved in the spontaneous breaking of a grand unification symmetry at a scale assuming its supersymmetric value. We combine a superpotential, which is uniquely determined by applying a continuous R symmetry, with a class of logarithmic or semilogarithmic Kähler potentials which exhibit a prominent shift symmetry with a tiny violation, whose strengths are quantified by c_- and c_+. The inflationary observables provide an excellent match to the recent BICEP2/Keck Array and Planck results, setting $3.5\times10^{-3}\lesssim r_\pm = c_+/c_- \lesssim 1/N$, where $N=3$ or 2 is the prefactor of the logarithm. Inflation can be attained for sub-Planckian inflaton values, with the corresponding effective theories retaining the perturbative unitarity up to the Planck scale.

DOI: 10.1103/PhysRevD.92.121305 PACS numbers: 98.80.Cq, 04.50.Kd, 04.65.+e, 12.60.Jv

I. INTRODUCTION

Soon after the inflation's [1] introduction as a solution to a number of longstanding cosmological puzzles—such as the horizon and flatness problems—many efforts have been made so as to connect it with a *Grand Unified Theory* (GUT) phase transition in the early Universe—see e.g. Refs. [2–10]. According to this economical and highly appealing setup, the scalar field which drives inflation (called an inflaton) plays, at the end of its inflationary evolution, the role of a Higgs field [2–6] or destabilizes other fields which act as Higgs fields [7–11]. As a consequence, a GUT gauge group $G_{\rm GUT}$ can be spontaneously broken after the end of inflation. The first mechanism above can be also applied in the context of the *Standard Model* (SM) [12] or the next-to-*Minimal Supersymmetric SM* (MSSM) [13,14] and leads to the spontaneous breaking of the electroweak gauge group $G_{\rm SM}= SU(3)_{\rm C}\times SU(2)_{\rm L}\times U(1)_Y$ by the Higgs/inflaton field(s).

We here focus on the earlier version of this idea—i.e. the GUT-scale Higgs inflation—concentrating on its *supersymmetric* (SUSY) realization [3–11], where the notorious GUT hierarchy problem is elegantly addressed. The starting point of our approach is the simplest superpotential

$$W=\lambda S(\bar\Phi\Phi - M^2/4), \tag{1}$$

which leads to the spontaneous breaking of $G_{\rm GUT}$ and is uniquely determined, at renormalizable level, by a convenient [7] continuous R symmetry. Here, λ and M are two constants which can both be taken positive by field redefinitions; S is a left-handed superfield, singlet under $G_{\rm GUT}$; $\bar\Phi$ and Φ is a pair of left-handed superfields belonging to nontrivial conjugate representations of $G_{\rm GUT}$, and reducing its rank by their *vacuum expectation values* (VEVs)—see e.g. Refs. [8,10]. Just for definiteness we restrict ourselves to the $G_{\rm GUT}=G_{\rm SM}\times U(1)_{B-L}$ [5,8] gauge group which consists of the simplest GUT beyond the MSSM—where B and L denote the baryon and lepton number. With the specific choice of $G_{\rm GUT}$, Φ and $\bar\Phi$ carry $B-L$ charges 1 and -1, respectively.

Moreover, W combined with a judiciously selected Kähler potential, K, gives rise to two types of inflation, in the context of *supergravity* (SUGRA). In particular, we can obtain *F-term hybrid inflation* (FHI) driven by S with $\bar\Phi$ and Φ being confined to zero or *nonminimal Higgs inflation* (nMHI), interchanging the roles of S and $\bar\Phi-\Phi$. A canonical [8] or quasicanonical [9,10] K is convenient for implementing FHI, whereas a logarithmic K including a holomorphic function $c_{\mathcal{R}}\Phi\bar\Phi$ with large $c_{\mathcal{R}}>0$ [4] or tiny $c_{\mathcal{R}}<0$ [5] is dictated for nMHI. Although FHI can become compatible with data [15] at the cost of a mild tuning of one [8,9] (or more [10]) parameters beyond λ and M, it exhibits a serious drawback which can be eluded, by construction, in nMHI. Since $G_{\rm GUT}$ is broken only at the SUSY vacuum, after the end of FHI, topological defects are formed, if they are predicted by the $G_{\rm GUT}$ breaking. This does not occur within nMHI, since $G_{\rm GUT}$ is already spontaneously broken during it, through the nonzero $\bar\Phi$ and Φ values. By utilizing large enough $c_{\mathcal{R}}$'s [4] or adjusting three parameters (λ, $c_{\mathcal{R}}$ and M) [5], acceptable values for the (scalar) spectral index, $n_{\rm s}$, can be achieved with a low enough [4] or higher [5] tensor-to-scalar ratio, r. In the former case, though, the largeness of $c_{\mathcal{R}}$ violates the perturbative unitarity [16,17], whereas in the latter case, trans-Planckian values of the inflaton jeopardize the validity of the inflationary predictions.

In this paper, we show that the shortcomings above can be elegantly overcome, if we realize the recently proposed [18] idea of *kinetically modified nonminimal inflation* with a $G_{\rm GUT}$ nonsinglet inflaton. The crucial difference of this setting compared to the nMHI with large $c_{\mathcal{R}}$ [4] is that the slope of the inflationary potential and the canonical normalization of the Higgs inflaton do not depend exclusively on one parameter, $c_{\mathcal{R}}$, but separately on two parameters, c_+ and c_-, whose ratio $r_\pm=c_+/c_-\ll1$ determines $n_{\rm s}$ and r. In particular, restricting $r_\pm$ to natural values, motivated by an enhanced shift symmetry, the inflationary observables can be nicely cover the 1σ domain of the present data [15,19],

*cpallis@ucy.ac.cy

$$n_s = 0.968 \pm 0.0045 \quad \text{and} \quad r = 0.048^{+0.035}_{-0.032}, \qquad (2)$$

independently of M, which may be confined precisely at its value entailed by the gauge unification within MSSM. Contrary to our recent investigation [6], where we stick to quadratic terms for Φ and $\bar{\Phi}$ in the selected K's, we here parametrize the relevant terms with an exponent m. Moreover, we here insist on integer prefactors of the logarithms involved in K's, increasing thereby the naturalness of the model. As regards other simple and well-motivated inflationary models [20,21] which share similar inflationary potentials with the one obtained here, let us underline that the use of a gauge nonsinglet inflaton with sub-Planckian values together with the enhanced resulting r's, in accordance with an approximate shift symmetry constitute the main novelties of our approach.

Below we describe a class of Kähler potentials which lead to kinetically modified nMHI, outline the derivation of the inflationary potential, and restrict the free parameters of the models testing them against observations. Finally, we analyze the *ultraviolet* (UV) behavior of these models and summarize our conclusions.

II. KÄHLER POTENTIALS

The key ingredient of our proposal is the selection of a purely or partially logarithmic K including the real functions

$$F_{\pm} = |\Phi \pm \bar{\Phi}^*|^2 \quad \text{and} \quad F_S = |S|^2 - k_S|S|^4, \qquad (3)$$

which respect the symmetries of W; the star (*) denotes complex conjugation. As we show below, c_-F_- dominates the canonical normalization of inflaton, c_+F_+ plays the role of the nonminimal inflaton-curvature coupling, and F_S provides a typical kinetic term for S, considering the next-to-minimal term for stability/heaviness reasons [13]. Obviously, F_S is the same as that used in Ref. [18], apart from an overall normalization factor, whereas F_- and F_+ correspond to $F_{\rm K}$ and $F_{\mathcal{R}}$, respectively. However, F_+ is a real and not a holomorphic function such as $F_{\mathcal{R}}$. Actually, it remains invariant under the transformation $\Phi \to \Phi + c$ and $\bar{\Phi} \to \bar{\Phi} - c^*$ (where c is a complex number), whereas F_- respects the symmetry $\Phi \to \Phi + c$ and $\bar{\Phi} \to \bar{\Phi} + c^*$, which coincides with the former only for $c = 0$. Stability of the selected inflationary direction entails that the latter symmetry is to be the dominant one—see below. The particular importance of the shift symmetry in taming the so-called η problem of inflation in SUGRA is first recognized for gauge singlets in Ref. [22] and for gauge nonsinglets in Ref. [14].

In terms of the functions introduced in Eq. (3), we postulate the following form of K:

$$K_1 = -3\ln\left(1 + c_+F_+ - \frac{1}{3}(1 + c_+F_+)^m c_-F_- - \frac{1}{3}F_S + k_\Phi F_-^2 + \frac{1}{3}k_{S\Phi}F_-|S|^2\right), \qquad (4a)$$

where we take for consistency all the possible terms up to fourth order, whereas a term of the form $-k_{S+}F_+{}^m|S|^2/3$ is neglected for simplicity, given that F_+ is considered as a violation of the principal symmetry—we use throughout units with the reduced Planck scale $m_{\rm P} = 2.433 \times 10^{18}$ GeV being set equal to unity. Identical results can be achieved if we select $K = K_2$ with

$$K_2 = -3\ln\left(1 + c_+F_+ - \frac{1}{3}F_S\right) + \frac{c_-F_-}{(1 + c_+F_+)^{1-m}}. \qquad (4b)$$

If we place F_S outside the argument of the logarithm, we can obtain two other K's—not mentioned in Ref. [18]—which lead to similar results. Namely,

$$K_3 = -2\ln\left(1 + c_+F_+ - \frac{1}{2}(1 + c_+F_+)^m c_-F_-\right) + F_S \qquad (4c)$$

and

$$K_4 = -2\ln(1 + c_+F_+) + F_S + (1 + c_+F_+)^{m-1}c_-F_-. \qquad (4d)$$

To highlight the robustness of our setting, we use only integer prefactors for the logarithms, avoiding thereby any relevant tuning—cf. Refs. [6,23]. Note that for $m = 0$ [$m = 1$], F_- and F_+ in K_1 and K_3 [K_2 and K_4] are totally decoupled; i.e., no higher-order term is needed. If we allow for a continuous variation of the ln prefactor, too, we can obtain several variants of kinetically modified nMHI. For $m = 0$ this possibility is analyzed in Ref. [6].

Given that $M \ll 1$ does not affect the inflationary epoch, the free parameters of our models, for fixed m, are $r_\pm$ and λ/c_-, and not c_-, c_+ and λ as naively expected. Indeed, performing the rescalings $\Phi \to \Phi/\sqrt{c_-}$ and $\bar{\Phi} \to \bar{\Phi}/\sqrt{c_-}$ in Eqs. (1) and (4a)–(4d), we see that W and K depend exclusively on λ/c_- and $r_\pm$, respectively. Therefore, our models are equally economical as nMHI with $c_{\mathcal{R}} < 0$ [5], and they have just one more free parameter than nMHI with $c_{\mathcal{R}} > 0$ [4]—see also Ref. [21]. Unlike these models, however, where the largeness [4] or the smallness [5] of $c_{\mathcal{R}}$ cannot be justified by any symmetry, our models can be characterized as completely natural, in the 't Hooft sense, since in the limits $r_\pm = c_+/c_- \to 0$ and $\lambda \to 0$, they enjoy the following enhanced symmetries:

$$\Phi \to \Phi + c, \quad \bar{\Phi} \to \bar{\Phi} + c^* \quad \text{and} \quad S \to e^{i\varphi}S, \qquad (5)$$

where c and φ are a complex and a real number, respectively. The same argument guarantees the smallness of k_{S+} in a possible term $-k_{S+}F_+{}^m|S|^2/3$ inside the logarithms in Eq. (4a) or Eq. (4b). On the other hand, our models do not exhibit any no-scale-type symmetry like that postulated in Ref. [20].

III. INFLATIONARY POTENTIAL

The *Einstein frame* (EF) action within SUGRA for the complex scalar fields $z^\alpha = S, \Phi, \bar{\Phi}$—denoted by the same superfield symbol—can be written as [13]

$$\mathsf{S} = \int d^4x\sqrt{-\widehat{\mathfrak{g}}}\left(-\frac{1}{2}\widehat{\mathcal{R}} + K_{\alpha\bar{\beta}}\widehat{g}^{\mu\nu}D_\mu z^\alpha D_\nu z^{*\bar{\beta}} - \widehat{V}\right), \quad (6a)$$

where summation is taken over z^α; $\widehat{\mathcal{R}}$ is the EF Ricci scalar curvature; D_μ is the gauge covariant derivative, $K_{\alpha\bar{\beta}} = K_{,z^\alpha z^{*\bar{\beta}}}$; and $K^{\alpha\bar{\beta}}K_{\bar{\beta}\gamma} = \delta^\alpha_\gamma$—the symbol, z as subscript denotes derivation with respect to z. Also, $\widehat{V}$ is the EF SUGRA potential which can be found in terms of W in Eq. (1) and the K's in Eqs. (4a)–(4d) via the formula

$$\widehat{V} = e^K\left(K^{\alpha\bar{\beta}}D_\alpha W D^*_{\bar{\beta}}W^* - 3|W|^2\right) + \frac{g^2}{2}\sum_a \mathrm{D}_a\mathrm{D}_a, \quad (6b)$$

where $D_\alpha W = W_{,z^\alpha} + K_{,z^\alpha}W$, $\mathrm{D}_a = z_\alpha(T_a)^\alpha_\beta K^\beta$ and the summation is applied over the generators T_a of G_{GUT}. If we express $\Phi, \bar{\Phi}$ and S according to the parametrization

$$\Phi = \frac{\phi e^{i\theta}}{\sqrt{2}}\cos\theta_\Phi, \quad \bar{\Phi} = \frac{\phi e^{i\bar{\theta}}}{\sqrt{2}}\sin\theta_\Phi, \quad \text{and} \quad S = \frac{s + i\bar{s}}{\sqrt{2}}, \quad (7)$$

with $0 \le \theta_\Phi \le \pi/2$, we can easily deduce from Eq. (6b) that a D-flat direction occurs at

$$\bar{s} = s = \theta = \bar{\theta} = 0 \quad \text{and} \quad \theta_\Phi = \pi/4, \quad (8)$$

along which the only surviving term in Eq. (6b) is

$$\widehat{V}_{\mathrm{HI}} = e^K K^{SS^*}|W_{,S}|^2 = \frac{\lambda^2(\phi^2 - M^2)^2}{16 f_{\mathcal{R}}^2}, \quad (9a)$$

since we obtain

$$K^{SS^*} = \begin{cases} f_{\mathcal{R}} \\ 1 \end{cases} \quad \text{for } K = \begin{cases} K_i \\ K_{i+2} \end{cases} \quad \text{with} \quad i = 1, 2, \quad (9b)$$

where $f_{\mathcal{R}} = 1 + c_+\phi^2$ plays the role of a nonminimal coupling to the Ricci scalar in the *Jordan frame* (JF). Indeed, if we perform a conformal transformation [6,13,23] defining the frame function as $\Omega/N = -\exp(-K/N)$, where

$$N = 3 \quad \text{or} \quad N = 2 \quad \text{for } K = K_i \quad \text{or} \quad K = K_{i+2}, \quad (10)$$

respectively, we can easily show that $f_{\mathcal{R}} = -\Omega/N$ along the path in Eq. (8). It is remarkable that $\widehat{V}_{\mathrm{HI}}$ turns out to be independent of the coefficients c_-, k_Φ and $k_{S\Phi}$ in Eqs. (4a) and (4b). Had we introduced the term $-k_{S+}F_+{}^m|S|^2/3$ inside the logarithms in Eqs. (4a) and (4b), we would have obtained an extra factor $(1 + k_{S+}\phi^{2m})$ in the denominator of $\widehat{V}_{\mathrm{HI}}$. Our results remain intact from this factor provided that $k_{S+} \lesssim 0.001$. Note, finally, that the conventional Einstein gravity is recovered at the SUSY vacuum,

$$\langle S\rangle = 0 \quad \text{and} \quad \langle\phi\rangle = M \ll 1, \quad (11)$$

since $\langle f_{\mathcal{R}}\rangle \simeq 1$.

To specify the EF canonically normalized inflaton, we note that, for all choices of K in Eqs. (4a)—(4d), $K_{\alpha\bar{\beta}}$ along the configuration in Eq. (8) takes the form

$$(K_{\alpha\bar{\beta}}) = \mathrm{diag}(M_K, K_{SS^*}) \quad \text{with} \quad M_K = \frac{1}{f_{\mathcal{R}}^2}\begin{pmatrix}\kappa & \bar{\kappa} \\ \bar{\kappa} & \kappa\end{pmatrix}, \quad (12)$$

where $\kappa = c_- f_{\mathcal{R}}{}^{1+m} - Nc_+$ and $\bar{\kappa} = Nc_+^2\phi^2$. Upon diagonalization of M_K we find its eigenvalues, which are

$$\kappa_+ = c_-\left(f_{\mathcal{R}}^{1+m} + Nr_\pm(c_+\phi^2 - 1)\right)/f_{\mathcal{R}}^2, \quad (13a)$$

$$\kappa_- = c_-\left(f_{\mathcal{R}}^m - Nr_\pm\right)/f_{\mathcal{R}}, \quad (13b)$$

where the positivity of κ_- is assured during and after nMHI for $r_\pm \lesssim 1/N$ given that $\langle f_{\mathcal{R}}\rangle \simeq 1$. By inserting Eqs. (7) and (12) into the second term of the *right-hand side* (rhs) of Eq. (6a), we can define the EF canonically normalized fields (denoted below by a hat), which are found to be

$$\frac{d\widehat{\phi}}{d\phi} = J = \sqrt{\kappa_+}, \quad \widehat{\theta}_+ = \frac{J\phi\theta_+}{\sqrt{2}}, \quad \widehat{\theta}_- = \sqrt{\frac{\kappa_-}{2}}\phi\theta_-, \quad (14a)$$

$$\widehat{\theta}_\Phi = \phi\sqrt{\kappa_-}(\theta_\Phi - \pi/4), \qquad (\widehat{s}, \widehat{\bar{s}}) = \sqrt{K_{SS^*}}(s, \bar{s}), \quad (14b)$$

where $\theta_\pm = (\bar{\theta} \pm \theta)/\sqrt{2}$. Note, in passing, that the spinors ψ_S and $\psi_{\Phi\pm}$ associated with the superfields S and $\Phi - \bar{\Phi}$ are normalized similarly, i.e., $\widehat{\psi}_S = \sqrt{K_{SS^*}}\psi_S$ and $\widehat{\psi}_{\Phi\pm} = \sqrt{\kappa_\pm}\psi_{\Phi\pm}$, with $\psi_{\Phi\pm} = (\psi_\Phi \pm \psi_{\bar{\Phi}})/\sqrt{2}$.

Taking the limit $c_- \gg c_+$, we find the expressions of the masses squared $\widehat{m}^2_{\chi^\alpha}$ (with $\chi^\alpha = \theta_+, \theta_\Phi$ and S) arranged in Table I, which approach rather well the quite lengthy, exact expressions taken into account in our numerical computation. These expressions assist us to appreciate the role of $k_S > 0$ in retaining a positive $\widehat{m}_s^2$ for $K = K_i$ and one heavy enough for $K = K_{i+2}$. Indeed, $\widehat{m}^2_{\chi^\alpha} \gg \widehat{H}_{\mathrm{HI}}{}^2 = \widehat{V}_{\mathrm{HI0}}/3$ for $\phi_{\mathrm{f}} \le \phi \le \phi_\star$, where $\phi_\star$ and ϕ_{f} are the values of ϕ when $k_\star = 0.05/\mathrm{Mpc}$ crosses the horizon of nMHI and at its end, correspondingly. In Table I we display also the masses, M_{BL}, of the gauge boson A_{BL}—which signals the fact that G_{GUT} is broken during nMHI—and the masses of the corresponding fermions.

The derived mass spectrum can be employed in order to find the one-loop radiative corrections, $\Delta\widehat{V}_{\mathrm{HI}}$, to $\widehat{V}_{\mathrm{HI}}$. Considering SUGRA as an effective theory with a cutoff scale equal to m_{P}, the well-known Coleman-Weinberg formula can be employed, self-consistently taking into account only the masses which lie well below m_{P}, i.e., all the masses arranged in Table I besides M_{BL} and $\widehat{m}_{\theta_\Phi}$. The resulting $\Delta\widehat{V}_{\mathrm{HI}}$ leaves intact our inflationary outputs, provided that the renormalization-group mass scale Λ is determined by requiring $\Delta\widehat{V}_{\mathrm{HI}}(\phi_\star) = 0$ or $\Delta\widehat{V}_{\mathrm{HI}}(\phi_{\mathrm{f}}) = 0$. The possible dependence of our findings on the choice of Λ can be totally avoided if we confine ourselves to $k_{S\Phi} \sim 1$ and $k_S \sim 1$, resulting in $\Lambda \simeq 3.2 \times 10^{-5} - 1.4 \times 10^{-4}$. Under these circumstances, our inflationary predictions can be exclusively reproduced by using $\widehat{V}_{\mathrm{HI}}$ in Eq. (9a), cf. Ref. [6].

TABLE I. Mass-squared spectrum for $K=K_i$ and $K=K_{i+2}$ ($i=1$, 2) along the path in Eq. (8).

			Masses Squared		
Fields	Eigenstates	Symbol	$K=K_1$	$K=K_2$	$K=K_{i+2}$
2 real scalars	$\widehat\theta_+$	$\widehat m^2_{\theta+}$	$4\widehat{H}_{\rm HI}^2$	$6\widehat{H}_{\rm HI}^2$	
	$\widehat\theta_\Phi$	$\widehat m^2_{\theta_\Phi}$	$M_{BL}^2+4\widehat{H}_{\rm HI}^2$	$M_{BL}^2+6\widehat{H}_{\rm HI}^2$	
1 complex scalar	$\widehat s,\widehat{\bar s}$	$\widehat m^2_s$	$6(2k_S f_{\cal R}-1/3)\widehat{H}_{\rm HI}^2$		$12k_S\widehat{H}_{\rm HI}^2$
1 gauge boson	A_{BL}	M_{BL}^2		$g^2c_-({f_{\cal R}}^{m-1}-Nr_\pm/f_{\cal R})\phi^2$	
4 Weyl spinors	$\widehat\psi_\pm=\frac{1}{\sqrt2}(\widehat\psi_{\Phi+}\pm\widehat\psi_S)$	$\widehat m^2_{\psi\pm}$		$24\widehat{H}^2_{\rm HI}/c_-\phi^2 f_{\cal R}^{1+m}$	
	$\lambda_{BL},\widehat\psi_{\Phi-}$	M_{BL}^2		$g^2c_-(f_{\cal R}^{m-1}-Nr_\pm/f_{\cal R})\phi^2$	

IV. INFLATIONARY REQUIREMENTS

Applying the standard formulas quoted in Ref. [18] for $\widehat V_{\rm CI}=\widehat V_{\rm HI}$, we can compute a number of observational quantities, which assist us to qualify our inflationary setting. Namely, we extract the number, $\widehat N_\star$, of e-foldings that the scale $k_\star$ experiences during nMHI and the amplitude, $A_{\rm s}$, of the power spectrum of the curvature perturbations generated by ϕ for $\phi=\phi_\star$. These observables must be compatible with the requirements [15]

$$\widehat N_\star\simeq61.5+\ln\frac{\widehat V_{\rm HI}(\phi_\star)^{\frac12}}{\widehat V_{\rm HI}(\phi_{\rm f})^{\frac14}}\quad\text{and}\quad A_{\rm s}^{\frac12}\simeq4.627\times10^{-5},\qquad(15)$$

where we consider an equation-of-state parameter $w_{\rm int}=1/3$ corresponding to quartic potential, which is expected to approximate rather well $\widehat V_{\rm HI}$ for $\phi\ll1$. We can then compute the model predictions as regards $n_{\rm s}$, its running, $a_{\rm s}$ and r or $r_{0.002}$—see Ref. [18]. The analytic expressions displayed in Ref. [18] for these quantities are applicable to our present case too, for $m>-1$, performing the following replacements:

$$n=4,\qquad r_{{\cal R}K}=r_\pm,\quad\text{and}\quad c_{\rm K}=c_-,\qquad(16)$$

and multiplying by a factor of 2 the rhs of the equation which yields λ in terms of c_-. We here concentrate on $m>-1$, since for smaller m's, confining $n_{\rm s}$ to its allowed region in Eq. (2), the predicted r's, although acceptable, lie well below the sensitivity of the present experiments [24]. This happens because, when decreasing m below 0, the first term on the rhs of Eq. (13a) becomes progressively subdominant, and thus c_+ controls both the slope of $\widehat V_{\rm HI}$ and the value of J in Eq. (14a) as in the standard nMHI [4,5].

The inflationary observables are not affected by M, provided that it is confined to values much lower than $m_{\rm P}$. This can be done if we determine it by identifying the unification scale (as defined by the gauge-coupling unification within the MSSM) $M_{\rm GUT}\simeq2/2.433\times10^{-2}$ with the value of M_{BL}—see Table I—at the SUSY vacuum. Given that $\langle f_{\cal R}\rangle\simeq1$ and $\langle\kappa_+\rangle\simeq1-Nr_\pm$, we obtain, for $r_\pm\lesssim1/N$,

$$M\simeq M_{\rm GUT}/g\sqrt{c_-(1-Nr_\pm)},\qquad(17)$$

with $g\simeq0.7$ being the value of the GUT gauge coupling constant. This result influences the inflaton mass at the vacuum, which is estimated to be $\widehat m_{\delta\phi}\simeq\lambda M/\sqrt{2c_-(1-Nr_\pm)}$.

V. RESULTS

Imposing the conditions in Eq. (15), we restrict λ/c_- and $\widehat\phi_\star$, whereas Eq. (2) constrains mainly m and $r_\pm$. Focusing initially on $K=K_i$ with $i=1$, 2, we present our results in Figs. 1 and 2. Namely, in Fig. 1 we compare the allowed curves in the $n_{\rm s}-r_{0.002}$ plane with the observational data [15] for $m=-1/2,0,1$ and 10, shown by the double dot-dashed, dashed, solid, and dot-dashed lines, respectively. The variation of $r_\pm$ is shown along each line. Note that for $m=0$, the line essentially coincides with the corresponding one in Ref. [6]—cf. Refs. [18,21]—and declines from the central $n_{\rm s}$ value in Eq. (2). On the other hand, the compatibility of the $m=1$ line with the central values in Eq. (2) is certainly impressive. For low enough $r_\pm$'s—i.e. $r_\pm\leq10^{-4}$—the various lines converge to the $(n_{\rm s},r_{0.002})$'s obtained within quartic inflation; whereas, for larger $r_\pm$, they enter the observationally allowed regions and terminate for $r_\pm\simeq1/3$, beyond which κ_- in Eq. (13b) ceases to be well defined. Notably, this restriction provides a lower bound on $r_{0.002}$ which increases with m. Indeed, we obtain $r_{0.002}\gtrsim0.0017,0.0028,0.009$ and 0.025 for $m=-1/2,0,1$ and 10, correspondingly. Therefore, our results are testable in forthcoming experiments [24].

Repeating the same analysis for $(-1)\leq m\leq10$, we can identify the allowed range of $r_\pm$, as in Fig. 2. The allowed (shaded) region is bounded by the dashed line, which corresponds to $r_\pm\simeq1/3$, and the dot-dashed and thin lines, along which the lower and upper bounds on $n_{\rm s}$ and r in Eq. (2) are saturated, respectively. We remark that increasing $r_\pm$ with fixed m, $n_{\rm s}$ increases whereas r decreases, in accordance with our findings in Fig. 1. We also infer that $r_\pm$ takes more natural (lower than unity) values for larger m's. Fixing $n_{\rm s}$ to its central value in Eq. (2), we obtain the solid line along which we get clear predictions for r, $a_{\rm s}$ and $\widehat m_{\delta\phi}$. Namely,

$$0.18\lesssim m\lesssim10\quad\text{and}\quad1/3\gtrsim r_\pm\gtrsim3.5\times10^{-3},\qquad(18a)$$

$$0.4\lesssim r/0.01\lesssim7.6\quad\text{and}\quad5.4\lesssim-a_{\rm s}/10^{-4}\lesssim6,\qquad(18b)$$

with $2.4\times10^{-8}\lesssim\hat m_{\delta\phi}\lesssim8.7\times10^{-6}$. Since the resulting $|a_{\rm s}|$ remains sufficiently low, our models are consistent with the

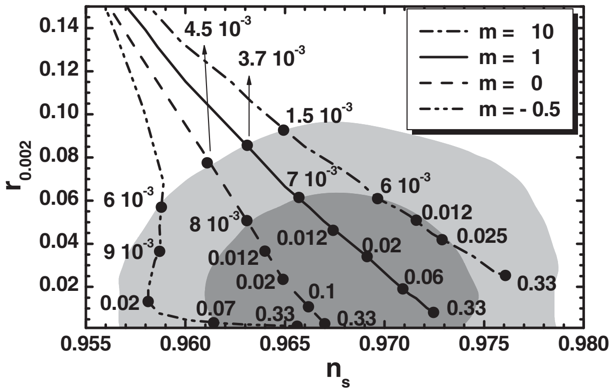

FIG. 1. Allowed curves in the $n_s - r_{0.002}$ plane for $K = K_i$ ($i = 1, 2$) and various m's (shown in the plot legend) and $r_\pm$'s indicated on the curves. The marginalized joint 68% [95%] regions from Planck, BICEP2/Keck Array, and Baryon Acoustic Oscillations (BAO) data are depicted by the dark [light] shaded contours.

fitting of data with the ΛCDM + r model [15]. Finally, the $\widehat{m}_{\delta\phi}$ range leaves open the possibility of nonthermal leptogenesis [25] if we introduce a suitable coupling between $\bar{\Phi}$ and the right-handed neutrinos—see e.g. Refs. [4,8].

Had we employed $K = K_{i+2}$, the various lines in Fig. 1 and the allowed regions in Fig. 2 would have been extended until $r_\pm \simeq 1/2$. This bound would have yielded $r_{0.002} \gtrsim 0.0012$, 0.002, 0.0066 and 0.023 for $m = -1/2$, 0, 1 and 10, correspondingly, which are a little lower than those designed in Fig. 1. The lower bounds of m, $r_\pm$ and r in Eqs. (18a) and (18b) become 0.19, 1/2, and 0.003; the upper bound on $\widehat{m}_{\delta\phi}$ moves on to 1.3×10^{-5}, whereas the bounds on $(-a_s)$ remain unaltered.

Although λ/c_- is constant in our setting for fixed $r_\pm$ and m, the amplitudes of λ and c_- can be bounded. This fact is illustrated in Fig. 3, where we display the allowed (shaded) area in the $\lambda - c_-$ plane, focusing on the $m = 1$ case. We observe that for any $r_\pm$ between its minimal (0.0037) and maximal (1/3) values—depicted by bold dot-dashed and dashed lines—there is a lower bound—represented by a faint dashed line—on c_-, above which $\phi_\star < 1$.

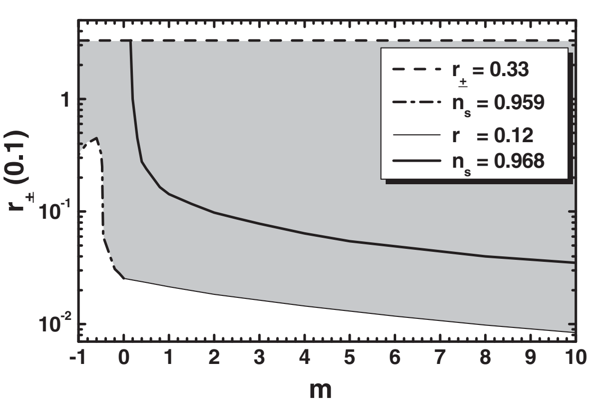

FIG. 2. Allowed (shaded) region in the $m - r_\pm$ plane for $K = K_i$. The conventions adopted for the various lines are also shown.

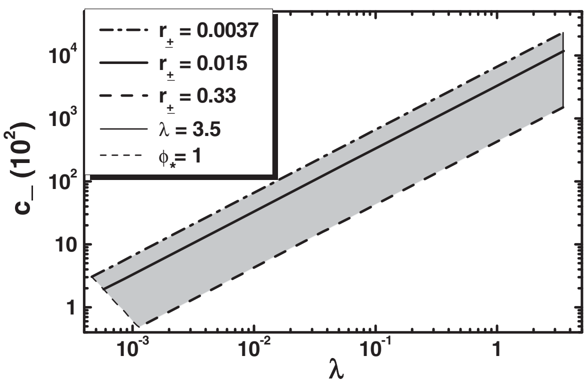

FIG. 3. Allowed (shaded) region in the $\lambda - c_-$ plane for $K = K_i$ (with $i = 1, 2$) and $m = 1$. The conventions adopted for the various lines are also shown.

Consequently, our proposal can be stabilized against corrections from higher order terms—e.g., $(\bar{\Phi}\Phi)^l$ with $l > 1$ in Eq. (1). The perturbative bound $\lambda = 3.5$ limits the region at the other end along the thin solid line. Plotted is also the solid line for $r_\pm = 0.015$, which yields $n_s = 0.968$. The corresponding $r = 0.043$ turns out to be impressively close to its central value in Eq. (2).

VI. THE EFFECTIVE CUTOFF SCALE

The fact that $\widehat{\phi}$ in Eq. (14a) does not coincide with ϕ at the vacuum of the theory—contrary to the pure nMHI [16,17]—assures that the corresponding effective theories respect perturbative unitarity up to $m_P = 1$, although c_- may take relatively large values for $\phi_\star < 1$—see Fig. 3. To clarify further this point, we analyze the small-field behavior of our models in the EF for $m = 1$. We focus on the second term on the rhs of Eq. (6a) for $\mu = \nu = 0$, and we expand it about $\langle\phi\rangle = M \ll 1$ in terms of $\widehat{\phi}$. Our result is written as

$$J^2\dot{\phi}^2 \simeq (1 + 3Nr_\pm{}^2\widehat{\phi}^2 - 5Nr_\pm{}^3\widehat{\phi}^4 + \cdots)\dot{\widehat{\phi}}^2. \qquad (19a)$$

Expanding similarly $\widehat{V}_{\rm HI}$—see Eq. (9a)—in terms of $\widehat{\phi}$ we have

$$\widehat{V}_{\rm HI} \simeq \frac{\lambda^2\widehat{\phi}^4}{16c_-^2}(1 - 2r_\pm\widehat{\phi}^2 + 3r_\pm{}^2\widehat{\phi}^4 - \cdots). \qquad (19b)$$

Similar expressions can be obtained for the other m's too. Given that the positivity of κ_- in Eq. (13a) entails $r_\pm \lesssim 1/N < 1$, we can conclude that our models do not face any problem with the perturbative unitarity up to m_P.

VII. CONCLUSIONS AND PERSPECTIVES

The feasibility of inflating with a superheavy Higgs field is certainly an archetypal open question. We here outlined a fresh look, identifying a class of Kähler potentials in Eqs. (4a)–(4d) which can cooperate with the superpotential in Eq. (1) and lead to the SUGRA potential $\widehat{V}_{\rm HI}$ collectively given by Eq. (9a). Prominent in the proposed Kähler potentials is the role of a shift-symmetric quadratic function F_- in Eq. (3) which remains invisible in $\widehat{V}_{\rm HI}$ while dominating the canonical normalization of the Higgs inflaton. Using $0.18[0.19] \leq m \leq 10$ and confining $r_\pm$ to the range $(3.5 \times 10^{-3} - 1/N)$ where $N = 3$ $[N = 2]$ for $K = K_i$ $[K = K_{i+2}]$, with $i = 1, 2$, we achieved observational predictions which may be tested in the near future and converge towards the "sweet spot" of the present data. These solutions can be attained even with sub-Planckian values of the inflaton requiring large c_-'s and without causing any problem with the perturbative unitarity. It is gratifying, finally, that our proposal remains intact from radiative corrections; the Higgs-inflaton may assume ultimately its VEV predicted by the gauge unification within MSSM, and the inflationary dynamics can be studied analytically and rather accurately.

As a last remark, we would like to point out that, although we have restricted our discussion to the $G_{\rm GUT} = G_{\rm SM} \times U(1)_{B-L}$ gauge group, kinetically modified nMHI has a much wider applicability. It can be realized, employing the same W and K's within other SUSY GUTs based on a variety of gauge groups, such as the left-right [10], the Pati-Salam [4], or the flipped $SU(5)$ group [10]—provided that Φ and $\bar{\Phi}$ consist of a conjugate pair of Higgs superfields so that they break $G_{\rm GUT}$ and compose the gauge-invariant quantities $F_\pm$. Moreover, given that the term $\lambda M^2 S/4$ of W in Eq. (1) plays no role during nMHI, our scenario can be implemented by replacing it with κS^3 and identifying Φ and $\bar{\Phi}$ with the electroweak Higgs doublets H_u and H_d of the next-to-MSSM [13]. In this case we have to modify the shift symmetry in Eq. (5), following the approach of Ref. [14]; consider the soft SUSY-breaking terms to obtain the radiative breaking of $G_{\rm SM}$; and take into account the renormalization group running of the various parameters from the inflationary up to the electroweak scale in order to connect convincingly the high- with the low-energy phenomenology. In all these cases, the inflationary predictions are expected to be quite similar to the ones obtained here, although the parameter space may be further restricted. The analysis of the stability of the inflationary trajectory may also be different, due to the different representations of Φ and $\bar{\Phi}$. Since our main aim here is the demonstration of the kinetical modification on the observables of nMHI, we opted to utilize the simplest GUT embedding.

[1] A. H. Guth, Phys. Rev. D **23**, 347 (1981); A. D. Linde, Phys. Lett. B **108**, 389 (1982); A. Albrecht and P. J. Steinhardt, Phys. Rev. Lett. **48**, 1220 (1982).
[2] D. S. Salopek, J. R. Bond., and J. M. Bardeen, Phys. Rev. D **40**, 1753 (1989); J. L. Cervantes-Cota and H. Dehnen, Phys. Rev. D **51**, 395 (1995).
[3] M. Arai, S. Kawai, and N. Okada, Phys. Rev. D **84**, 123515 (2011); K. Nakayama and F. Takahashi, J. Cosmol. Astropart. Phys. 05 (2012) 035; M. B. Einhorn and D. R. T. Jones, J. Cosmol. Astropart. Phys. 11 (2012) 049; L. Heurtier, S. Khalil, and A. Moursy, J. Cosmol. Astropart. Phys. 10 (2015) 045.
[4] C. Pallis and N. Toumbas, J. Cosmol. Astropart. Phys. 12 (2011) 002.
[5] M. Arai, S. Kawai, and N. Okada, Phys. Lett. B **734**, 100 (2014).
[6] G. Lazarides and C. Pallis, J. High Energy Phys. 11 (2015) 114.
[7] G. R. Dvali, Q. Shafi, and R. K. Schaefer, Phys. Rev. Lett. **73**, 1886 (1994).
[8] C. Pallis and Q. Shafi, Phys. Lett. B **725**, 327 (2013).
[9] M. Bastero-Gil, S. F. King, and Q. Shafi, Phys. Lett. B **651**, 345 (2007).
[10] M. Civiletti, C. Pallis, and Q. Shafi, Phys. Lett. B **733**, 276 (2014).
[11] S. Antusch, M. Bastero-Gil, J. P. Baumann, K. Dutta, S. F. King, and P. M. Kostka, J. High Energy Phys. 08 (2010) 100; J. Ellis, T. E. Gonzalo, J. Harz, and W. C. Huang, J. Cosmol. Astropart. Phys. 03 (2015) 039.
[12] J. L. Cervantes-Cota and H. Dehnen, Nucl. Phys. **442**, 391 (1995); F. L. Bezrukov and M. Shaposhnikov, Phys. Lett. B **659**, 703 (2008).
[13] M. B. Einhorn and D. R. T. Jones, J. High Energy Phys. 03 (2010) 026; H. M. Lee, J. Cosmol. Astropart. Phys. 08 (2010) 003; S. Ferrara, R. Kallosh, A. Linde, A. Marrani, and A. Van Proeyen, Phys. Rev. D **83**, 025008 (2011).
[14] I. Ben-Dayan and M. B. Einhorn, J. Cosmol. Astropart. Phys. 12 (2010) 002.
[15] Planck Collaboration, arXiv:1502.02114.
[16] J. L. F. Barbon and J. R. Espinosa, Phys. Rev. D **79**, 081302 (2009); C. P. Burgess, H. M. Lee, and M. Trott, J. High Energy Phys. 07 (2010) 007.
[17] A. Kehagias, A. M. Dizgah, and A. Riotto, Phys. Rev. D **89**, 043527 (2014).
[18] C. Pallis, Phys. Rev. D **91**, 123508 (2015); *Proc. Sci.*, PLANCK (2015) 095.
[19] P. A. Ade *et al.* (BICEP2/Keck Array and Planck Collaborations), Phys. Rev. Lett. **114**, 101301 (2015).
[20] J. Ellis, D. V. Nanopoulos, and K. A. Olive, Phys. Rev. Lett. **111**, 111301 (2013); **111**, 129902(E) (2013).
[21] R. Kallosh, A. Linde, and D. Roest, Phys. Rev. Lett. **112**, 011303 (2014).
[22] M. Kawasaki, M. Yamaguchi, and T. Yanagida, Phys. Rev. Lett. **85**, 3572 (2000).
[23] R. Kallosh, A. Linde, and D. Roest, J. High Energy Phys. 08 (2014) 052; C. Pallis, J. Cosmol. Astropart. Phys. 10 (2014) 058.
[24] P. Creminelli *et al.*, arXiv:1502.01983.
[25] G. Lazarides and Q. Shafi, Phys. Lett. B **258**, 305 (1991); K. Kumekawa, T. Moroi, and T. Yanagida, Prog. Theor. Phys. **92**, 437 (1994).