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Constraining $\mathcal{N} = 1$ supergravity inflationary framework with non-minimal Kähler operators

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Abstract: In this paper we will illustrate how to constrain unavoidable Kähler corrections for $\mathcal{N} = 1$ supergravity (SUGRA) inflation from the recent Planck data. We will show that the non-renormalizable Kähler operators will induce in general *non-minimal kinetic* term for the inflaton field, and two types of SUGRA corrections in the potential — the *Hubble-induced mass* (c_H), and the *Hubble-induced A-term* (a_H) correction. The entire SUGRA inflationary framework can now be constrained from (i) the *speed of sound*, c_s, and (ii) from the upper bound on the *tensor to scalar ratio*, $r_\star$. We will illustrate this by considering a heavy scalar degree of freedom at a scale, M_s, and a light inflationary field which is responsible for a slow-roll inflation. We will compute the corrections to the kinetic term and the potential for the light field explicitly. As an example, we will consider a visible sector inflationary model of inflation where inflation occurs at the point of *inflection*, which can match the density perturbations for the cosmic microwave background radiation, and also explain why the universe is filled with the Standard Model degrees of freedom. We will scan the parameter space of the non-renormalizable Kähler operators, which we find them to be order $\mathcal{O}(1)$, consistent with physical arguments. While the scale of heavy physics is found to be bounded by the tensor-to scalar ratio, and the speed of sound, $\mathcal{O}\left(10^{11} \leq M_s \leq 10^{16}\right)$ GeV, for $0.02 \leq c_s \leq 1$ and $10^{-22} \leq r_\star \leq 0.12$.

Keywords: Cosmology of Theories beyond the SM, Supersymmetric Effective Theories, Supergravity Models

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Contents

1 Introduction

The success of primordial inflation [1–4], for a review, see [5], can be gauged by the current observations arising from the comic microwave background (CMB) radiation [6–8]. The observations from Planck have put interestingly tight bounds on a number of unknown parameters of a generic inflationary model [6], in particular the *speed of sound*, c_s, of the perturbations, which also determines any departure from the Gaussian perturbations, the local type of non-Gaussianity, $f_{NL}^{\rm local}$, and the constraint on *tensor-to-scalar ratio*, $r_\star$, which can potentially unearth the scale of New Physics within any given *effective field theory* set-up.

The $\mathcal{N} = 1$ supergravity (SUGRA) [9, 10], for a review, see [11], is an excellent well-defined set-up where we can address some of the key questions about the physics of the new scale for instance. The Kähler metric determines the kinetic term for the inflaton potential, and one particular choice is the *minimal* Kinetic term for the inflaton field. However, quantum corrections to the Kähler potential is not very well-known. Generically corrections to the Kähler potential arise from integrating out the heavy physics, and due to lack of concrete knowledge on the details of the heavy physics, many times computing these corrections can be very challenging, see [12].

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The aim of this paper is to place a generic bound on the Planck suppressed corrections to the Kähler potential on top of the minimal Kähler potential. We will consider dimensional 3 and dimensional 4 *gauge invariant* non-renormalizable Kähler operators in this paper. Since these correction will lead to a departure from the minimal kinetic term for the inflationary potential, such corrections can now be bounded from the Planck data, especially from the speed of sound of the primordial perturbations. In fact the Kähler potential within $\mathcal{N} = 1$ SUGRA can induce corrections to the inflationary potential which can yield large Hubble-induced mass correction to the inflaton field [13–17], some times known as the SUGRA-η problem, and the Hubble-induced SUGRA A-term for the potential. In the context of $\mathcal{N} = 1$ SUGRA hybrid inflation [18–20], some of the Kähler corrections were constrained by the tensor to scalar ratio, $r_\star$, and the spectral tilt, n_S, of the power spectrum. In this paper we will be constraining these Kähler corrections systematically from the interaction of the heavy physics to the light inflaton field from the recent Planck data. Let us consider the scale of New Physics, M_s, to be within

$$M_p \geq M_s \geq H_{\rm inf}\,, \tag{1.1}$$

where $H_{\rm inf}$ is the Hubble parameter during inflation, and $M_p = 2.4 \times 10^{18}\,{\rm GeV}$. In order to constrain M_s, we would require to consider at least 2 fields, one which is heavy at the relevant scale, M_s, and the other which is light. We will assume that these two fields are coupled gravitationally. In past such a scenario has been considered by many authors, where the heavy field leaves interesting imprints in the dynamics of a low scale inflation [23–35]. Broadly speaking there are two possible scenarios which one can envisage:

- *The heavy field is dynamically frozen*: we can imagine that the heavy field is completely *frozen*, in which case it would be effectively a single light field with a canonical kinetic term for the inflaton field, with a speed of sound, $c_s = 1$. For a slow roll inflation, the perturbations will be primarily Gaussian. If the heavy field is settled down to its minimum VEV, i.e. zero, then there will be no effect from the heavy field at all. However, if the heavy field is settled with a finite non-zero VEV, and it remains dynamically inactive, means its velocity is strictly zero, then it can still contribute to the vacuum energy density of the inflaton, and this would be encoded in $H_{\rm inf}$. Also, the kinetic term for the light inflaton field will depart from being pure canonical. However the departure will depend on the scale of new physics. If $M_s \ll M_p$, then the departure from canonical kinetic term will be negligible for all practical purposes. Therefore, again the observational predictions for the CMB will be unaltered and will be similar to the previous case. Both of these scenarios were taken into account by various interesting papers, see for example [17, 18, 20–22, 36], and here we will not consider them in great details. We will analyse a slightly different scenario as mentioned below.

- *The heavy field is coherently oscillating during the initial phases of inflation*: in this case we will consider a very simple scenario, where we imagine that the heavy field is coherently oscillating at a VEV given by M_s with an amplitude M_s at the onset

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of inflation driven by the light field. The coherent oscillations of the heavy field will not last forever, its amplitude would be damped during inflation very rapidly within couple of e-foldings of inflation. However, just right at the onset of inflation, the relevant modes which are leaving the Hubble patch for the CMB can be constrainable. This will provide a window of opportunity for us to constrain such a scenario, see refs. [23–25] for probing the influence of heavy physics into the light inflaton field. In this paper we will consider a similar scenario, but in the context of $\mathcal{N} = 1$ SUGRA.

First of all the coherent oscillations of the heavy field around its non-zero vacuum would provide a non-zero vacuum energy density, i.e $\sim M_s^4$. Through its coupling to the light field in the Kähler potential, it would also yield non-canonical kinetic term contribution to the light field, and therefore $c_s \neq 1$ for a slow roll inflaton field. Eventually, the heavy field will be settled down to its minimum. We presume that the dominant contribution to the long wavelength fluctuations are still seeded by the light inflaton, but the fact that $c_s \neq 1$ for the inflaton, it would leave imprints which would be constrainable directly by the scale of heavy physics, M_s, and the non-minimal corrections to the kähler potential.

We will discuss this latter scenario in some details, and provide a full $\mathcal{N} = 1$ SUGRA potential for the light and the heavy field within a simple example. We will be using the following constraints from CMB, and also requirement for a guaranteed reheating of the Standard Model d.o.f for the success of big bang nucleosynthesis [37].

1. Successful single field inflation driven by ϕ field with the right amplitude and tilt of the power spectrum.

$$2.092 < 10^9 P_S < 2.297 \quad (\text{within } 2\sigma)\,, \tag{1.2}$$

$$0.958 < n_S < 0.963 \qquad (\text{within } 2\sigma)\,. \tag{1.3}$$

2. Speed of sound, c_s: the Planck analysis has constrained it to be [6, 8]:

$$0.02 \leq c_s \leq 1 \quad (\text{within } 2\sigma)\,. \tag{1.4}$$

3. Tensor-to-scalar ratio: the Planck constraint is $r_\star \leq 0.12$ [6]

$$r_\star \equiv \frac{P_T}{P_S} \leq 0.12. \tag{1.5}$$

4. Local type of non-Gaussianity, $f_{NL}^{\rm local}$: the Planck constraint on local non-Gaussianity is [6, 8]:

$$f_{NL}^{\rm local} = 2.7 \pm 5.8 \quad (\text{within } 1\sigma)\,. \tag{1.6}$$

In this paper we will not consider the constraints arising from various non-Gausisnaity bounds [8], but we will solely focus on the constraints arising from the speed of sound during perturbation, and the tensor to scalar ratio. In the companion paper we have discussed how non-Gaussianity can constrain the Kähler corrections.

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5. Particle physics constraint: we wish to ensure that the inflaton *solely* decays into the Standard Model (SM) d.o.f, therefore we embed the light fields within supersymmetric SM, such as minimal supersymmetric Standard Model (MSSM). In this case the inflaton carries *solely* the SM charges as in the case of the inflaton driven by *gauge invariant* combinations of squarks and sleptons [38–41]. This will naturally ensure that we obtain the right abundances for the dark matter and the baryons in the universe as required by the observations [42–45]. One can however follow a hidden sector or a SM gauge singlet inflaton, but it is not always straightforward to explain the universe with the right dark matter abundance [46–48], and baryon asymmetry, see [5].

The results of the first half of this paper will be very generic — applicable to any inflationary scenario. In section 2, we will discuss briefly the Planck constraints. In section 3, we will discuss the setup with one heavy and one light superfield which are coupled via gravitational interactions through Kähler potential. In section 4, we will describe the effective field theory potential for the light superfield Φ, and discuss the kinetic terms for various interesting scenarios. In section 5, we will discuss the role of non-canonical kinetic term and consider two possibilities, one where the heavy superfield is dynamically frozen, see section 5.1, and the more interesting scenario when the heavy field is oscillating at the onset of inflation, see section 5.2. We will scan the parameters for the Planck suppressed Kähler operators in subsection 5.3, we will discuss how tensor-to-scalar ratio, $r_\star$, can constrain the mass scale of the heavy physics.

2 Cosmological perturbations for $c_s \neq 1$

In this section we briefly recall some of the important formulae when $c_s \neq 1$, the scalar and tensor perturbations are given by [6, 49, 50]:

$$P_S(k) = \mathcal{P}_\mathcal{S}\left(\frac{k}{c_s k_\star}\right)^{n_S-1}\,,$$
$$P_T(k) = \mathcal{P}_\mathcal{T}\left(\frac{k}{c_s k_\star}\right)^{n_T}\,, \tag{2.1}$$

where the speed of sound at the Hubble patch is given by, $c_s k_\star = aH$ (where $k_\star \sim 0.002\,\mathrm{Mpc}^{-1}$). The amplitude of the scalar and tensor perturbations can be recast in terms of the potential, as [6]:

$$\mathcal{P}_\mathcal{S} = \frac{V_\star}{24\pi^2 M_p^4 c_s \epsilon_V}\,, \tag{2.2}$$

$$\mathcal{P}_\mathcal{T} = \frac{2V_\star}{3\pi^2 M_p^4} c_s^{\frac{2\epsilon_V}{1-\epsilon_V}}\,, \tag{2.3}$$

where running of the spectral tilt for the scalar and tensor modes can be expressed at $c_s k_\star = aH$, as:

$$n_S - 1 = 2\eta_V - 6\epsilon_V - s\,, \tag{2.4}$$

$$n_T = -2\epsilon_V\,, \tag{2.5}$$

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where running of the sound speed is defined by an additional slow-roll parameter, s, as:

$$s = \frac{\dot{c}_s}{H c_s} = \sqrt{\frac{3}{V}} \frac{\dot{c}_s}{c_s} M_p. \tag{2.6}$$

In all the above expressions, the standard slow-roll parameters (with $c_s = 1$) are defined by:

$$\epsilon_V = \frac{M_p^2}{2} \left(\frac{V'}{V} \right)^2, \qquad \eta_V = M_p^2 \left(\frac{V''}{V} \right). \tag{2.7}$$

Finally, the single field consistency relation between tensor-to-scalar ratio and tensor spectral tilt is modified by [6, 51]:

$$r_\star = 16 \epsilon_V c_s^{\frac{1+\epsilon_V}{1-\epsilon_V}} = -8 n_T c_s^{\frac{1-\frac{n_T}{2}}{1+\frac{n_T}{2}}}. \tag{2.8}$$

Using the results for $c_s \neq 1$ stated in eqs. (2.2)–(2.8), the upper bound on the numerical value of the Hubble parameter during inflation is given by:

$$H \leq 9.241 \times 10^{13} \sqrt{\frac{r_\star}{0.12}}\, c_s^{\frac{\epsilon_V}{\epsilon_V - 1}}\ \text{GeV} \tag{2.9}$$

where $r_\star$ is the tensor-to-scalar ratio at the pivot scale of momentum $k_\star \sim 0.002\,\text{Mpc}^{-1}$. An equivalent statement can be made in terms of the upper bound on the energy scale of inflation for $c_s \neq 1$ as:

$$V_\star \leq \left(1.96 \times 10^{16} \text{GeV}\right)^4 \frac{r_\star}{0.12}\, c_s^{\frac{2\epsilon_V}{\epsilon_V - 1}}. \tag{2.10}$$

Here in eqs. (2.9) and (2.10), the equalities will hold good for a high scale model of inflation.

Furthermore, for a sub-Plancikan slow-roll models of inflation, one can express the tensor-to-scalar ratio, $r_\star$, at the pivot scale, $k_\star \sim 0.002\,\text{Mpc}^{-1}$, in terms of the field displacement, $\Delta\phi$, during the observed $\Delta N \approx 17$ e-foldings of inflation, for $c_s \neq 1$ [52, 53]:

$$\frac{3}{25\sqrt{c_s}} \sqrt{\frac{r_\star}{0.12}} \left| \left\{ \frac{3}{400} \left(\frac{r_\star}{0.12} \right) - \frac{\eta_V(k_\star)}{2} - \frac{1}{2} \right\} \right| \approx \frac{|\Delta\phi|}{M_p}, \tag{2.11}$$

where $\Delta\phi = \phi_{cmb} - \phi_e \ll M_p$, where ϕ_{cmb} and ϕ_e are the values of the inflaton field at the horizon crossing and at the end of inflation.

3 Inflationary setup within $\mathcal{N} = 1$ SUGRA

Let us consider two sectors; heavy sector denoted by the superfield S, and the light sector denoted by Φ. Let us assume that the two sectors interact only gravitationally, S could denote the hidden sector, while Φ could denote the visible sector for example part of MSSM [54, 55]. Note that the origin of S superfield need not be always hidden sector, within MSSM it is possible to have a false vacuum at very high VEVs, see for instance [56].

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In the latter case both S and Φ could be embedded within MSSM for instance.[1] For the purpose of illustration, we will assume S to have a simple superfield potential given by:

$$W = W(\Phi) + W(S)\,, \quad \Phi = \text{Light}\,,\ S = \text{Heavy}\,, \tag{3.1}$$

$$= \frac{\lambda\Phi^n}{nM_p^{n-3}} + \frac{M_s}{2}S^2\,, \tag{3.2}$$

where $n \geq 3$ and $\lambda \sim \mathcal{O}(1)$, and Φ superfield is the D-flat direction of MSSM. The scale M_s governs the scale of heavy physics. Furthermore, we will assume $\langle s\rangle,\ \langle\phi\rangle \leq M_p$, where both s and ϕ are fields corresponding to the super field S and Φ. There are two flat directions which can drive inflation with $n = 6$, which are lifted by themselves [58–60],

$$\phi = \frac{\widetilde{u} + \widetilde{d} + \widetilde{d}}{\sqrt{3}}, \qquad \phi = \frac{\widetilde{L} + \widetilde{L} + \widetilde{e}}{\sqrt{3}} \tag{3.3}$$

where $\widetilde{u},\ \widetilde{d}$ denote the right handed squarks, and $\widetilde{L}$ denotes that left handed sleptons and $\widetilde{e}$ denotes the right handed slepton. In this case the inflaton mass for ϕ will be given by:

$$m_\phi^2 = \frac{m_{\widetilde{L}}^2 + m_{\widetilde{L}}^2 + m_{\widetilde{e}}^2}{3}\,, \qquad m_\phi^2 = \frac{m_{\widetilde{u}}^2 + m_{\widetilde{d}}^2 + m_{\widetilde{d}}^2}{3}\,, \tag{3.4}$$

for $\widetilde{L}\widetilde{L}\widetilde{e}$ and $\widetilde{u}\widetilde{d}\widetilde{d}$ directions respectively. Typically these masses are set by the scale of SUSY, which is typically of the order of $\geq \mathcal{O}(1)$ TeV, set by he ATLAS and CMS [76, 77].

Let us consider minimal Kähler potentials for both ϕ and s. For the purpose of illustration we will consider the simplest choice which produces minimal kinetic term, and the corrections are of the form:

$$K = s^\dagger s + \phi^\dagger\phi + \delta K\,, \tag{3.5}$$

where *gauge invariant* Kähler corrections:

$$\delta K = f\left(\phi^\dagger\phi, s^\dagger s\right),\ f\left(s^\dagger\phi\phi\right),\ f\left(s^\dagger s^\dagger\phi\phi\right),\ f\left(s\phi^\dagger\phi\right)\,. \tag{3.6}$$

The higher order corrections to the Kähler potentials are extremely hard to compute. In the following, we will assume that the leading order corrections are of the generic form — allowed by the gauge invariance.[2] For the purpose of illustration, let us consider the

[1] By visible sector we mean that the inflaton itself carries the SM charges, such as in the case of MSSM [38–41]. In all these examples the inflaton Φ is the D-flat direction made up of squarks and sleptons, see [57], which is lifted by the F-term of the non-renormalizable superpotential.

[2] For MSSM flat directions, some of these corrections were already considered before in the context of Affleck-Dine baryogenesis, see [58–61]. For MSSM inflation this paper is the first to deal with these corrections explicitly.

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following terms:

$$K^{(1)} = \phi^\dagger\phi + s^\dagger s + \frac{a}{M_p^2}\phi^\dagger\phi s^\dagger s + \cdots , \tag{3.7}$$

$$K^{(2)} = \phi^\dagger\phi + s^\dagger s + \frac{b}{2M_p} s^\dagger\phi\phi + \text{h.c.} + \cdots , \tag{3.8}$$

$$K^{(3)} = \phi^\dagger\phi + s^\dagger s + \frac{c}{4M_p^2} s^\dagger s^\dagger\phi\phi + \text{h.c.} + \cdots , \tag{3.9}$$

$$K^{(4)} = \phi^\dagger\phi + s^\dagger s + \frac{d}{M_p} s\phi^\dagger\phi + \text{h.c.} + \cdots , \tag{3.10}$$

where $a,\ b,\ c,\ d$ are dimensionless parameters.[3] These corrections will inevitably lead to a departure from the minimal kinetic energy for both the fields. Our aim will be to constrain these unknown parameters, i.e. $a,\ b,\ c,\ d$, and the scale of heavy physics, M_s, from the CMB constraints mentioned above in the introduction.

4 Effective field theory potential for inflaton from $\mathcal{N} = 1$ SUGRA

Typically, the scalar potential in $\mathcal{N} = 1$ SUGRA for the F-term can be written in terms of the superpotential, W, and the Kähler potential, K, see [11]:

$$V = e^{K\left(\Phi_i^\dagger,\Phi_i\right)/M_p^2}\left[(D_{\Phi_i}W(\Phi))K^{\Phi_i\bar{\Phi}_j}\left(D_{\bar{\Phi}_j}W^*\left(\Phi^\dagger\right)\right) - 3\frac{|W(\Phi)|^2}{M_p^2}\,,\right] \tag{4.1}$$

where $i = \Phi,\ S$ in our case, and $F_\Phi \equiv D_\Phi W = W_\Phi + K_\Phi/M_p^2$, and $K^{\Phi_i\bar{\Phi}_j}$ is the inverse matrix of $K_{\Phi_i\bar{\Phi}_j}$, and the subscript denotes derivative with respect to the field.

Typically at the leading order, the total potential will get contributions from [58–60]:

1. Interaction between flat direction and inflaton via exponential prefactor:

$$e^{K\left(\phi,\phi^\dagger\right)/M_p^2}V(s)\,.$$

2. Cross coupling terms between the flat direction induced Kähler derivative and the inflaton superpotential:

$$K_\phi K^{\phi\bar{\phi}}K_{\bar{\phi}}\frac{|W(s)|^2}{M_p^4}\,.$$

3. Interaction between the Kähler derivative and superpotential of the inflaton, supergravity Kähler metric and Kähler potential of the flat direction:

$$K_\phi K^{\phi\bar{s}}D_{\bar{s}}W^*\left(s^\dagger\right)\frac{W(s)}{M_p^2} + \text{h.c.}\,.$$

[3] The $\cdots$ contain higher order terms of type $(1/M_p)^2\left(\phi^\dagger\phi\right)^2 + (1/M_p)^2\left(s^\dagger s\right)^2 + \ldots$, and higher order terms, here we are ignoring them. These corrections have been taken into account in the context of SUGRA hybrid inflation in refs. [18–22]. Here we are mainly interested in considering the effects of heavy field s on the dynamics of a light field ϕ.

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4. Self coupling between inflaton via Kähler derivative interaction:

$$\left(D_S W(s)\right) K^{s\bar{s}} \left(D_{\bar{S}} W^* \left(s^\dagger\right)\right) .$$

Additionally, the Hubble-induced A terms arises from the following dominant contributions in the effective theory of supergravity [58–60]:

1. Cross coupling terms in the Kähler derivative between the derivative of the flat direction superpotential and inflaton superpotential:

$$W_\phi K^{\phi\bar{\phi}} K_{\bar{\phi}} \frac{W^*\left(s^\dagger\right)}{M_p^2} + \text{h.c.} .$$

2. Interaction terms between the flat direction superpotential and inflaton Kähler derivative:

$$K_s \frac{W(\phi)}{M_p^2} K^{s\bar{s}} \left(D_{\bar{S}} W^* \left(s^\dagger\right)\right) + \text{h.c.} .$$

3. Cross coupling terms between the flat direction and inflaton superpotential:

$$-\frac{3}{M_p^2} W^* \left(s^\dagger\right) W(\phi) + \text{h.c.} .$$

4. Couplings between the flat direction and inflaton:

$$W_\phi K^{\phi\bar{s}} \left(D_{\bar{S}} W^* \left(s^\dagger\right)\right) + \text{h.c.} .$$

The resulting *leading order* potential for the light field ϕ at low energies can be captured by the following terms [57–60]:

$$V(\phi) = V(s) + \left(m_\phi^2 + c_H H(t)^2\right) |\phi|^2 + \left(A \frac{\lambda \phi^n}{n M_p^{n-3}} + a_H H(t) \frac{\lambda \phi^n}{n M_p^{n-3}} + \text{h.c.}\right) + \lambda^2 \frac{|\phi|^{2(n-1)}}{M_p^{2(n-3)}} + \cdots , \tag{4.2}$$

where $A \sim m_\phi$ is the dimension full quantity, $\cdots$ contain terms of higher orders, c_H, a_H are numbers containing the information about the Kähler potential, we can infer them from table 1, and appendix B and C. Note that during inflation, $H(t) \sim H_{\text{inf}}$, is nearly constant. Note that there are two kinds of Hubble-induced terms; one proportional to the mass term, and the second of the order of the A-term.

5 Non-minimal Kähler potential and non-canonical kinetic terms

In this section we will consider two interesting possibilities, one which is the simplest and provides an excellent model for inflation with a complete decoupling of the heavy field. Inflation occurs via the slow roll of ϕ field within an MSSM vacuum, where inflation would end in a vacuum with an *enhanced gauge symmetry*, where the entire electroweak symmetry will be completely restored.

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Non-minimal Kähler potential	**Non-canonical kinetic term** $\mathcal{L}_{\rm Kin} = K_{ij^*}\left(\partial_\mu\Phi^i\right)\left(\partial^\mu\Phi^{j^*}\right)$	**Potential** $V(\phi)$ **for** $\lvert s\rvert \ll M_p$
$K^{(1)} = \phi^\dagger\phi + s^\dagger s + \frac{a}{M_p^2}\phi^\dagger\phi s^\dagger s$	$\mathcal{L}_{\rm Kin} = \left(1+\frac{a\lvert s\rvert^2}{M_p^2}\right)\left(\partial_\mu\phi\right)\left(\partial^\mu\phi^\dagger\right) + \frac{a}{M_p^2}\left\{\phi^\dagger s\left(\partial_\mu\phi\right)\left(\partial^\mu s^\dagger\right) + \phi s^\dagger\left(\partial_\mu s\right)\left(\partial^\mu\phi^\dagger\right)\right\} + \left(1+\frac{a\lvert\phi\rvert^2}{M_p^2}\right)\left(\partial_\mu s\right)\left(\partial^\mu s^\dagger\right)$	$V(s) + \left(m_\phi^2 + 3(1-a)H^2\right)\lvert\phi\rvert^2 - \frac{A\phi^n}{nM_p^{n-3}} - \left(1+a\frac{\lvert s\rvert^2}{M_p^2}\right)\left(1-\frac{3}{n}\right)\frac{s^2}{M_p^2}\frac{\lambda M_s\phi^n}{M_p^{n-3}} - \left(1-a\frac{\lvert s\rvert^2}{M_p^2}\right)\left(a-\frac{1}{n}\right)\frac{(s^\dagger)^2}{M_p^2}\frac{\lambda M_s\phi^n}{M_p^{n-3}} + \lambda^2\frac{\lvert\phi\rvert^{2(n-1)}}{M_p^{2(n-3)}} + \text{h.c.}$
$K^{(2)} = \phi^\dagger\phi + s^\dagger s + \frac{b}{2M_p}s^\dagger\phi\phi + \text{h.c.}$	$\mathcal{L}_{\rm Kin} = (\partial_\mu\phi)(\partial^\mu\phi^\dagger) + (\partial_\mu s)\left(\partial^\mu s^\dagger\right) + \frac{b\phi}{2M_p}\left(\partial_\mu\phi\right)\left(\partial^\mu s^\dagger\right) + \frac{b\phi^\dagger}{2M_p}\left(\partial_\mu s\right)\left(\partial^\mu\phi^\dagger\right)$	$V(s) + \left(m_\phi^2 + 3\left(1+b^2\right)H^2\right)\lvert\phi\rvert^2 - A\frac{\phi^n}{nM_p^{n-3}} - \left\{\left(1-\frac{3}{n}\right)\phi + \frac{b\phi^\dagger s}{nM_p}\right\}\frac{\lambda\phi^{n-1}M_s s^2}{M_p^{n-1}} - \left(\frac{s^\dagger\phi}{M_p} - bn\phi^\dagger\right)\frac{2M_s\lambda\phi^{n-1}s^\dagger}{nM_p^{n-2}} - \frac{bM_s s^2}{2M_p^2}\left(\frac{2M_s s}{M_p} - \frac{M_s s^2 s^\dagger}{M_p^3}\right)\phi\phi - \frac{4M_s^2 b\lvert s\rvert^2 s^\dagger}{M_p^3}\phi\phi + \lambda^2\frac{\lvert\phi\rvert^{2(n-1)}}{M_p^{2(n-3)}} + \text{h.c.}$
$K^{(3)} = \phi\phi^\dagger + ss^\dagger + \frac{c}{4M_p^2}s^\dagger s^\dagger\phi\phi + \text{h.c.}$	$\mathcal{L}_{\rm Kin} = (\partial_\mu\phi)\left(\partial^\mu\phi^\dagger\right) + (\partial_\mu s)\left(\partial^\mu s^\dagger\right) + \frac{cs^\dagger\phi}{4M_p^2}\left(\partial_\mu\phi\right)\left(\partial^\mu s^\dagger\right) + \frac{cs\phi^\dagger}{4M_p^2}\left(\partial_\mu s\right)\left(\partial^\mu\phi^\dagger\right)$	$V(s) + \left(m_\phi^2 + 3H^2\right)\lvert\phi\rvert^2 - A\frac{\phi^n}{nM_p^{n-3}} - \left\{\left(1-\frac{3}{n}\right)\phi + \frac{c\phi^\dagger ss}{2M_p^2}\right\}\frac{\lambda\phi^{n-1}M_s s^2}{M_p^{n-1}} + \frac{cM_s^2 s^2 s^\dagger s\phi\phi}{M_p^4} - \frac{M_s^2 c\lvert s\rvert^2 s^\dagger s^\dagger}{M_p^4}\phi\phi + \lambda^2\frac{\lvert\phi\rvert^{2(n-1)}}{M_p^{2(n-3)}} - \left(\frac{s^\dagger\phi}{M_p} - \frac{cn\phi^\dagger s}{M_p}\right)\frac{2M_s\lambda\phi^{n-1}s^\dagger}{nM_p^{n-2}} + \text{h.c.}$
$K^{(4)} = \phi\phi^\dagger + ss^\dagger + \frac{d}{M_p}s\phi^\dagger\phi + \text{h.c.}$	$\mathcal{L}_{\rm Kin} = \left(\frac{ds}{M_p} + \frac{ds^\dagger}{M_p} + 1\right)\left(\partial_\mu\phi\right)\left(\partial^\mu\phi^\dagger\right) + \left(\partial_\mu s\right)\left(\partial^\mu s^\dagger\right) + \frac{d\phi^\dagger}{M_p}\left(\partial_\mu\phi\right)\left(\partial^\mu s^\dagger\right) + \frac{d\phi}{M_p}\left(\partial_\mu s\right)\left(\partial^\mu\phi^\dagger\right)$	$V(s) + \left(m_\phi^2 + 3\left(1+d^2\right)H^2\right)\lvert\phi\rvert^2 - A\frac{\phi^n}{nM_p^{n-3}} - \left(1-\frac{3}{n}\right)\frac{\lambda\phi^n M_s s^2}{M_p^{n-1}} - \left(\frac{s^\dagger}{M_p} - d\right)\frac{2M_s\lambda\phi^n s^\dagger}{nM_p^{n-2}} + \lambda^2\frac{\lvert\phi\rvert^{2(n-1)}}{M_p^{2(n-3)}} + \text{h.c.}$

Table 1. Various supergravity effective potentials and non-canonical kinetic terms for $|s| \ll M_p$ in presence non-niminmal Kähler potential. Here both ϕ and s are complex fields, and so are the A-terms.

5.1 Heavy field is *dynamically* frozen

Let us first assume that the dynamics of the heavy field s is completely frozen during the onset and the rest of the course of slow roll inflation driven by ϕ. The full potential can be found in table 1. Note that the potential for s field, $V(s)$ contains soft term and the corresponding A-term:

$$V(s) \sim M_s^2|s|^2 + A' M_s s^2\,, \tag{5.1}$$

where A' is a dimensional quantity, and it is roughly proportional to $A' \sim M_s \gg$ TeV. In this case there are two possibilities which we briefly mention below:

- we can imagine that the heavy field, s, to have a global minimum at:

$$\langle s\rangle = 0, \qquad \langle\dot{s}\rangle = 0\,, \qquad V(s) = 0\,. \tag{5.2}$$

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In this particular setup, the kinetic terms for each cases, i.e. $1, 2, 3, 4$, become canonical for the ϕ field, therefore the heavy field is completely decoupled from the dynamics. One can check them from table 1. This is most ideal situation for a single field dominated model of inflation, where the overall potential for along ϕ direction simplifies to:

$$V(\phi) = m_\phi^2|\phi|^2 + \left(A\frac{\lambda\phi^n}{nM_p^{n-3}} + \text{h.c.}\right) + \lambda^2\frac{|\phi|^{2(n-1)}}{M_p^{2(n-3)}}\,. \tag{5.3}$$

The overall potential is solely dominated by the ϕ field, therefore Hubble expansion rate becomes, $H_{\rm inf} \propto V(\phi)/M_p^2$.

In this setup inflation can occur near a saddle point or an inflection point, where $\phi_0 \ll M_p$, and $m_\phi \gg H_{\rm inf}$, first discussed in refs. [38, 39]. During inflation the Hubble expansion rate is smaller than the soft SUSY breaking mass term and the A-term, i.e. $A \sim m_\phi \gg H(t)$ for $a_H \sim c_H \sim \mathcal{O}(1)$ in eq. (4.2), such that the SUGRA corrections are unimportant. This scenario has been discussed extensively, and has been extremely successful with the Planck data explaining the spectral tilt right on the observed central value, with Gaussian perturbations with the right amplitude [36, 62].

- On the other hand, if

$$\langle s\rangle \sim M_s \ll M_p\,, \qquad \langle\dot{s}\rangle = 0\,, \qquad V(s) = M_s^4\,, \tag{5.4}$$

then the kinetic term for ϕ field will be canonical for cases K^2 and K^3 by virtue $\dot{s} = 0$, see table 1. However for cases K^1 and K^4, the departure from canonical for the ϕ field will depend on M_s. If $\langle M_s\rangle \ll M_p$, and $a,\ d \sim \mathcal{O}(1)$, see table 1, then the kinetic term for ϕ will be virtually canonical, and as a consequence $c_s \approx 1$, while the potential will see a modification:

$$V_{\rm total} = M_s^4 + c_H H^2|\phi|^2 + \left(a_H H\frac{\lambda\phi^n}{nM_p^{n-3}} + \text{h.c.}\right) + \lambda^2\frac{|\phi|^{2(n-1)}}{M_p^{2(n-3)}}\,. \tag{5.5}$$

This large vacuum energy density, i.e. $M_s^4 \gg (\text{TeV})^4$, would yield a large Hubble expansion rate, i.e. $H_{\rm inf}^2 \sim M_s^4/M_p^2 \gg m_\phi^2 \sim \mathcal{O}(\text{TeV})^2$. Therefore, the Hubble induced mass and and the A-term would dominate the potential over the soft terms. Inspite of large mass, c_H, and a_H-term, there is *no* SUGRA-η problem, provided inflation occurs near the *saddle point* or the *inflection point* [17, 36]. We will not discuss this case any further, we will now focus on a slightly non-trivial scenario, where high scale physics can alter some of the key cosmological predictions.

5.2 Heavy field is oscillating during the onset of inflation

One dramatic way the heavy field can influence the dynamics of primordial perturbations is via coherent oscillations around its minimum, while ϕ still plays the role of a slow roll inflaton.[4] Furthermore, the heavy field would only influence the first few e-foldings of

[4]There could be other scenarios where the influence of heavy field is felt throughout the inflationary dynamics, see for instance in refs. [18–22, 29–31]. Here we will discuss a slightly simpler scenario where both heavy and light sectors are coupled gravitationally via the Kähler correction.

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inflation, once the heavy field is settled down its effect would be felt only via the vacuum energy density. Inspite of this short-lived phase, the heavy field can influence the dynamics and the perturbations for the light field as we shall discuss below.

Let us imagine the heavy field is coherently oscillating around a VEV, $\langle s \rangle \sim M_s$, during the initial phase of inflation, such that

$$V(s) \neq 0, \qquad \langle s \rangle \neq 0, \qquad \langle \dot{s} \rangle \neq 0\,. \tag{5.6}$$

The origin of coherent oscillations of s field need not be completely ad-hoc, such a scenario might arise quite naturally from the hidden sector moduli field which is coherently oscillating before being damped away by the initial phase of inflation, see for instance [80]. This is particularly plausible for high string scale moduli, where the moduli mass can be heavy and can be stabilised early on in the history of the universe. There could also be a possibility of a smooth second order phase transition from one vacuum to another during the intermittent phases of inflation [63, 64]. Such a possibility can arise within MSSM where there are multiple false vacua at high energies [56]. Irrespective of the origin of this heavy field, during this transient period, the heavy field with an effective mass, $M_s \gg H_{\rm inf}$, can coherently oscillate around its vacuum. We can set its initial amplitude of the oscillations to be of the order M_s.

$$s(t) = M_s + M_s \sin\left(M_s t\right)\,. \tag{5.7}$$

This also implies that at the lowest order approximation, $\langle s \rangle \sim M_s$ and $\langle \dot{s} \rangle \sim M_s^2$.[5] The contribution to the potential due to the time dependent oscillating heavy field, see eq. (5.7), is averaged over a full cycle $\left(0 < t_{\rm osc} < H_{\rm inf}^{-1}\right)$ is given by:

$$\langle V(s) \rangle \approx M_s^2 \langle s^2(t) \rangle \sim H_{\rm inf}^2 M_p^2\,. \tag{5.8}$$

The s field provides at the lowest order corrections to the kinetic term for the ϕ field, and to the overall potential, see table 1, for both kinetic and potential terms.

At this point one might worry, the coherent oscillations of the s field might trigger particle creation from the time dependent vacuum, see refs. [65–69], for a review see [70]. First of all, if we assume that the heavy field is coupled to other fields gravitationally, then the particle creation may not be sufficient to back react into the inflationary potential. Furthermore, inflation would also dilute the quanta created during this transient phase. We would not expect any imprint of this event on cosmological scales [71], except one interesting possibility could be to excite some non-Gaussianity [72–75]. In this paper we will not study the effects of non-Gaussianity, we shall leave this question for the companion paper.

Since the kinetic terms for the 4 cases tabulated in table 1 are now no longer canonical, they would contribute to the speed of sound, $c_s \neq 1$, which we can summarize case by case

[5] At this point one might say why we had taken the amplitude of oscillations for the heavy field to be M_s. In some scenarios, it is possible to envisage the amplitude of the oscillations to be M_p. This would not alter much of our discussion, therefore for the sake of simplicity we will consider the initial amplitude for the s field to be displaced by M_s, the same as that of the VEV.

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below:

$$c_s = \sqrt{\frac{\dot{p}}{\dot{\rho}}} \approx \begin{cases} \sqrt{\dfrac{\mathbf{X}_1(t) - \mathbf{X}_2(t) - \dot{\widehat{V}}}{\mathbf{X}_1(t) + \mathbf{X}_3(t) + \dot{\widehat{V}}}} & \text{for } \underline{\textbf{Case I}} \\ \sqrt{\dfrac{\mathbf{Y}_1(t) - \mathbf{Y}_2(t) - \dot{\widehat{V}}}{\mathbf{Y}_1(t) + \mathbf{Y}_3(t) + \dot{\widehat{V}}}} & \text{for } \underline{\textbf{Case II}} \\ \sqrt{\dfrac{\mathbf{Z}_1(t) - \mathbf{Z}_2(t) - \dot{\widehat{V}}}{\mathbf{Z}_1(t) + \mathbf{Z}_3(t) + \dot{\widehat{V}}}} & \text{for } \underline{\textbf{Case III}} \\ \sqrt{\dfrac{\mathbf{W}_1(t) - \mathbf{W}_2(t) - \dot{\widehat{V}}}{\mathbf{W}_1(t) + \mathbf{W}_3(t) + \dot{\widehat{V}}}} & \text{for } \underline{\textbf{Case IV}}\,, \end{cases} \tag{5.9}$$

where p is the effective pressure and ρ is the energy density. The dot denotes derivative w.r.t. physical time, t. All the symbols, i.e. $X_1,\ X_2,\ Y_1,\ Y_2,\ Z_1,\ Z_2,\ W_1,\ W_2$, appearing in eq. (5.9) are explicitly mentioned in the appendix. Additionally, here we have defined, $\widehat{V} = V(\phi) - V(s)$.[6]

5.3 Constraining non-renormalizable operators, i.e. a, b, c, d, and M_s

For the potential under consideration, we have $V(s) = 3H^2M_p^2 \sim M_s^2 s^2 \gg m_\phi^2|\phi|^2$, where $m_\phi \sim \mathcal{O}(\text{TeV})$ is the soft mass. In this case the contributions from the Hubble-induced terms are important compared to the soft SUSY breaking mass, m_ϕ, and the A term for all the four cases tabulated in table 1. The potential, eq. (4.2), after stabilizing the angular direction of the complex scalar field $\phi = |\phi|\exp[i\theta]$, see [17, 38–40], reduces to a simple form along the real direction, which is dominated by a single scale, i.e. $H \sim H_{\rm inf}$:

$$V(\phi) = V(s) + c_H H^2|\phi|^2 - \frac{a_H H\phi^n}{nM_p^{n-3}} + \frac{\lambda^2|\phi|^{2(n-1)}}{M_p^{2(n-3)}}, \tag{5.10}$$

where we take $\lambda = 1$, and,the Hubble-induced mass parameter c_H, for $s \ll M_p$, is defined as:[7]

$$c_H = \begin{cases} 3(1-a)\,, & \text{for } \underline{\textbf{Case I}} \\ 3\left(1+b^2\right)\,, & \text{for } \underline{\textbf{Case II}} \\ 3\,, & \text{for } \underline{\textbf{Case III}} \\ 3\left(1+d^2\right)\,, & \text{for } \underline{\textbf{Case IV}}\,. \end{cases} \tag{5.11}$$

Note that for only third case, i.e. K^3, the Hubble induced mass term does not contain any Kähler correction, i.e. δK. Similarly, we can express a_H, see appendix C for full

[6] As a side remark, our analysis will be very useful for the Affleck-Dine (AD) baryogenesis [58–60], especially when the minimum of the AD field is rotating in presence of the inflaton oscillations. Effectively, the AD field will have non-canonical kinetic terms, this has never been taken into account in the literature and one should take the non-canonical kinetic terms for the AD field in presence of the inflaton oscillations in order to correctly estimate the baryon asymmetry. The role of s field will be that of an inflaton and ϕ field will be that of an AD field.

[7] See appendix B and appendix C for details.

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expressions. Note that for all 4 cases, the kinetic terms are all non-minimal, and we have already listed in table 1. Fortunately for this class of potential given by eq. (5.10), inflection point inflation can be accommodated, when $a_H^2 \approx 8(n-1)c_H$. This can be characterized by a fine-tuning parameter, δ, which is defined as [38]:

$$\frac{a_H^2}{8(n-1)c_H} = 1 - \left(\frac{n-2}{2}\right)^2 \delta^2 \,. \tag{5.12}$$

When $|\delta|$ is small,[8] a point of inflection ϕ_0 exists, such that $V''(\phi_0) = 0$, with

$$\phi_0 = \left(\sqrt{\frac{c_H}{(n-1)}} H M_p^{n-3}\right)^{1/n-2} + \mathcal{O}\left(\delta^2\right). \tag{5.13}$$

For $\delta < 1$, we can Taylor-expand the inflaton potential around an inflection point, $\phi = \phi_0$, as [17, 78, 79]:

$$V(\phi) = \alpha + \beta(\phi - \phi_0) + \gamma(\phi - \phi_0)^3 + \kappa(\phi - \phi_0)^4 + \cdots, \tag{5.14}$$

where the expansion coefficients are now given by:

$$\alpha = V(\phi_0) = V(s) + \left(\frac{(n-2)^2}{n(n-1)} + \frac{(n-2)^2}{n}\delta^2\right) c_H H^2 \phi_0^2 + \mathcal{O}\left(\delta^4\right), \tag{5.15}$$

$$\beta = V'(\phi_0) = 2\left(\frac{n-2}{2}\right)^2 \delta^2 c_H H^2 \phi_0 + \mathcal{O}\left(\delta^4\right), \tag{5.16}$$

$$\gamma = \frac{V'''(\phi_0)}{3!} = \frac{c_H H^2}{\phi_0}\left(4(n-2)^2 - \frac{(n-1)(n-2)^3}{2}\delta^2\right) + \mathcal{O}\left(\delta^4\right), \tag{5.17}$$

$$\begin{aligned}\kappa &= \frac{V''''(\phi_0)}{4!} \\ &= \frac{c_H H^2}{\phi_0^2}\left(12(n-2)^3 - \frac{(n-1)(n-2)(n-3)\left(7n^2 - 27n + 26\right)}{2}\delta^2\right) + \mathcal{O}\left(\delta^4\right).\end{aligned} \tag{5.18}$$

Note that once we specify c_H and $H_{\rm inf}$, all the terms in the potential are determined. In this regard the potential indeed simplifies a lot to study the cosmological observables.

As an concrete example, we considered $n = 6$ case, where the flatness of the superfield Φ is lifted by the non-renormalizable operator. This is appropriate for both $\widetilde{u}\widetilde{d}\widetilde{d}$ and $\widetilde{L}\widetilde{L}\widetilde{e}$ flat directions. In our scans we allow the constraints from Planck observations [6, 7], see eqs. (1.2), (1.3), (1.4), (1.5).

Let us now scan the parameter space for c_H, a_H with the help of eqs. (2.2), (2.3), (2.4), (2.5), (2.8), by fixing $\lambda = \mathcal{O}(1)$ and $\delta \sim 10^{-4}$. In order to satisfy the Planck observational constraints on the amplitude of the power spectrum, $2.092 \times 10^{-9} < P_S < 2.297 \times 10^{-9}$ (within 2σ), spectral tilt $0.958 < n_S < 0.963$ (within 2σ), sound speed $0.02 \leq c_s \leq 1$ (within 2σ), and tensor-to-scalar ratio $r_\star \leq 0.12$, we obtain the following constraints on our parameters for $H_{\rm inf} \geq m_\phi \sim \mathcal{O}(\text{TeV})$, where successful inflation can occur via *inflection*

[8] We will consider a moderate tuning of order $\delta \sim 10^{-4}$ between c_H and a_H.

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point inflation:

$$\begin{aligned} c_H &\sim \mathcal{O}\left(10 - 10^{-6}\right) , && \text{for} \quad 10^{-22} < r_\star < 0.12 \\ a_H &\sim \mathcal{O}\left(30 - 10^{-3}\right) , && \text{for} \quad 10^{-22} < r_\star < 0.12 \\ M_s &\sim \mathcal{O}\left(9.50 \times 10^{10} - 1.77 \times 10^{16}\right) \text{ GeV} , && \text{for} \quad 10^{-22} < r_\star < 0.12 . \end{aligned} \tag{5.19}$$

Our motivation of doing such a scan is to generate feasible amplitude of power spectrum P_s , spectral tilt n_s, sound speed c_s and tensor to scalar ratio $r_\star$, which also satisfies the particle physics constraints in our prescribed inflationary setup. As these constraints are necessary to satisfy the inflation, we have to choose the parameter space in such a way that all of these constraints satisfy simultaneously. Inflation would not occur outside our scanning region since at least one of the constraints would be violated.

Note that for the above ranges, eq. (5.19), ϕ_0 gets automatically fixed by eq. (5.13),

$$\phi_0 \sim \mathcal{O}\left(10^{14} - 10^{17}\right) \text{ GeV} \quad \text{for} \quad 10^{-22} < r_\star < 0.12 . \tag{5.20}$$

Here the upper and lower bound appearing in eq. (5.19) and eq. (5.20) are obtained from large and small values of the tensor-to-scalar ratio varying within a wide range $10^{-22} < r_\star < 0.12$ for the pivot scale $k_\star \sim 0.002\,\text{Mpc}^{-1}$.

In figure 1, we have shown that the allowed ranges of the non-renormalizable coefficients of the operators mentioned in eqs. (3.7), (3.8), (3.9), (3.10). The solid blue and red curves are drawn for the sound speed $c_s = 0.02$ and $c_s = 1$ and the shaded regions are shown to point out the allowed region which satisfies the Planck 2σ constraints on the amplitude of power spectrum P_S and spectral tilt n_S as mentioned in eq. (1.2) and eq. (1.3) respectively. It is important to note that the non-minimal couplings "a", "b" and "d" directly controls both c_H, a_H in the inflaton potential. But the coupling "c" only affect a_H, while leaving c_H free from non-minimal correction, i.e. $c_H \sim 3$. For the consistency check see appendix where all the non-minimal couplings "a", "b", "c" and "d" are explicitly written in terms of the scale (VEV) of the heavy field M_s.

In figure 2(a) and figure 2(b) we have shown the constraints on the amplitude of the the power spectrum for scalar modes, P_S, and $\log(r)$, with respect to spectral tilt, n_S at the pivot scale $k_\star = 0.002$ Mpc^{-1} by *red* and *blue* curves for the sound speed, $c_s = 1$ and $c_s = 0.02$, respectively. If we consider the full parameter space as stated in eq. (5.19), there are solutions which have been shown in a *yellow* and *aqua* shaded regions. We have also shown the 2σ region allowed by the Planck data [6] for both the cases by *green* shaded region, i.e. P_S and n_S. It is important to note that if we consider the full parameter space then the low c_s fits the data well compared to the high value c_s.

However from figure 2(a) it is clearly observed that the high value of c_s also confronts the data well within a small patch for a specific choice of parameter space lying within the parameter scanning range mentioned in eq. (5.19). Consequently the full parameter space for low c_s and a tiny patch for high c_s fits the CMB power spectra well in the low l ($2 < l < 49$) and high l ($50 < l < 2500$) multipole region. But for the low l ($2 < l < 49$) region, the statistical error is too huge to differentiate between different c_s scenarios. Therefore, we will concentrate only on the high l ($50 < l < 2500$) region for the low and high c_s model discrimination with high statistical accuracy (2σ C.L.). See figure 3 for the details where we explicitly use this additional input.

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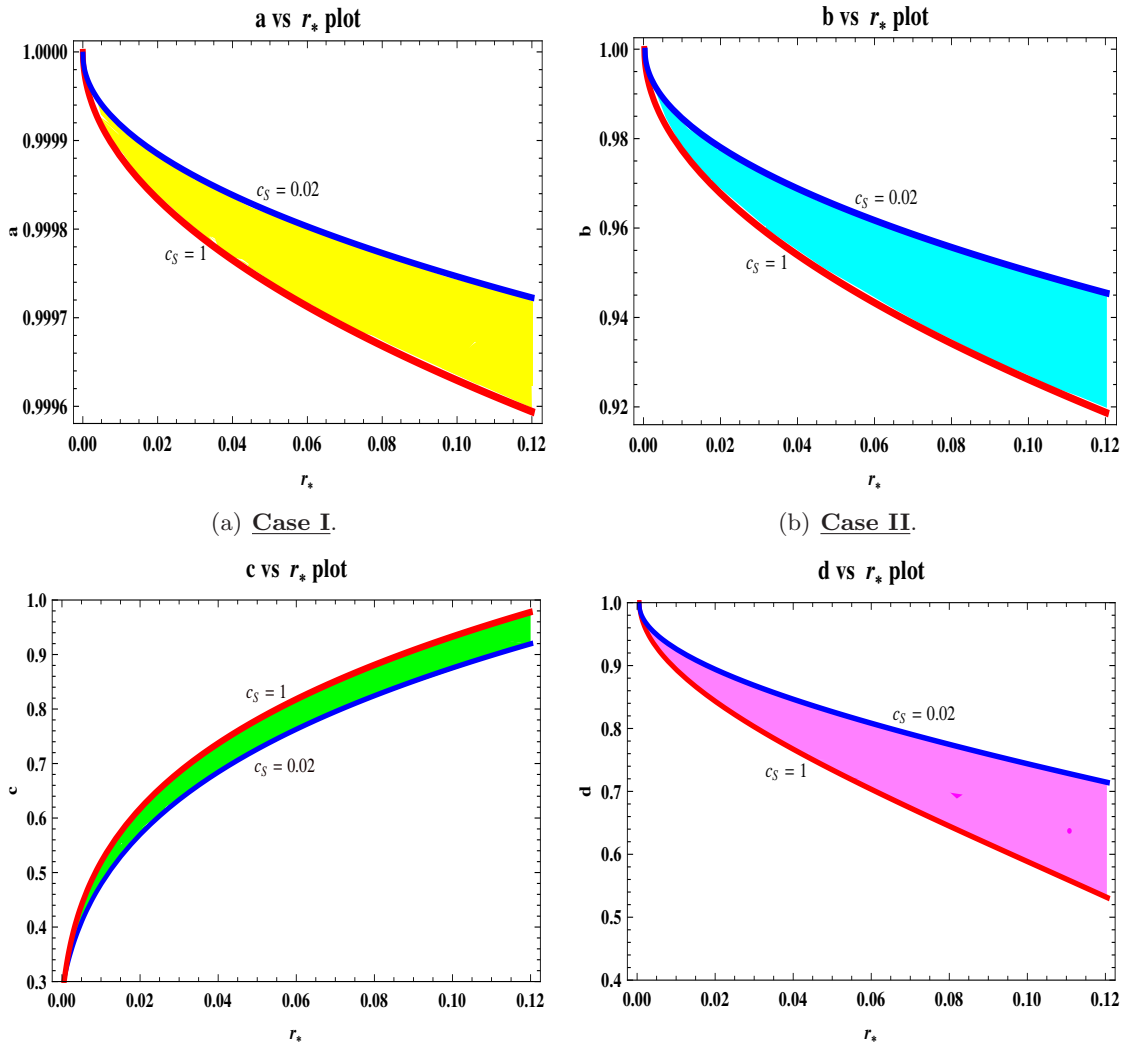

(a) **Case I**. (b) **Case II**.

(c) **Case III**. (d) **Case IV**.

Figure 1. We show the constraints on the non-renormalizable Kähler operators, "a","b","c" and "d" with respect to the tensor-to-scalar ratio $r_\star$ at the pivot scale $k_\star = 0.002\,\mathrm{Mpc}^{-1}$ when the heavy field s is oscillating during the initial phase of inflation, especially at the time when the interesting perturbations are leaving the Hubble patch for $H_{\rm inf} \gg m_\phi \sim \mathcal{O}(\mathrm{TeV})$. All the shaded regions represent the allowed parameter space for the Hubble induced inflation satisfying the Planck 2σ constraints on the amplitude of power spectrum $2.092 \times 10^{-9} < P_S < 2.297 \times 10^{-9}$ and spectral tilt $0.958 < n_S < 0.963$, as mentioned in eq. (1.2) and eq. (1.3) respectively. The dark coloured boundaries are obtained from the allowed range of the speed of sound c_s, within the window $0.02 \leq c_s \leq 1$.

Furthermore, by using *Planck*+WMAP-9 [6, 7], *Planck*+WMAP-9+high l [6, 7] and *Planck*+WMAP-9+BAO datasets [6, 7], we have shown r vs. n_s in the marginalized 1σ and 2σ CL. contours in figure 3. The *yellow* and *green* lines are drawn for the proposed model with $c_s = 1$ and $c_s = 0.02$ respectively. The region in between the *yellow* and *green* lines represent the allowed region obtained from the proposed model within the

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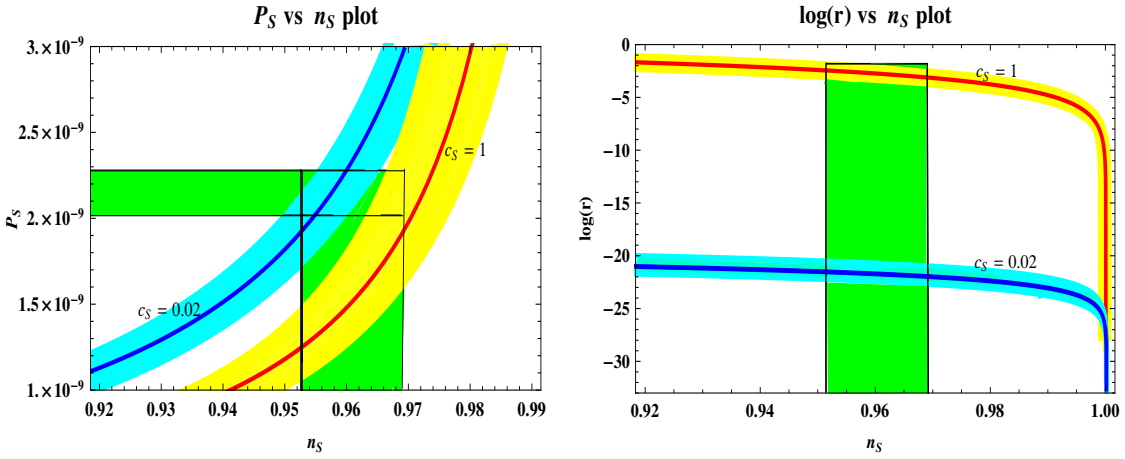

(a) P_S vs n_S. (b) $\log(r)$ vs n_S.

Figure 2. For $H_{\rm inf} \geq m_\phi \sim \mathcal{O}(\text{TeV})$, we have shown the variation of (a) P_S vs n_S, and (b) $\log(r)$ vs n_S, for MSSM flat direction (for $\widetilde{u}\widetilde{d}\widetilde{d}$ and $\widetilde{L}\widetilde{L}\widetilde{e}$ inflaton cadidates) in presence of non-minimal Kähler corrections. The *red* and *blue* curves are drawn for $c_s = 1$ and $c_s = 0.02$, which show the model parameters, $\delta \sim 10^{-4}$, $\lambda = 1$, $c_H = 2$, $a_H = 2.108$, $\phi_0 = 1.129 \times 10^{16}$ GeV, for the pivot scale $k_\star = 0.002\,\text{Mpc}^{-1}$. The *green* shaded region shows the 2σ CL. range allowed by the Planck data [6] for both the cases. Instead of getting a solid *red* and *blue* curves we obtain a *yellow* and *aqua* shaded regions if we would consider the full parameter space for ($H_{\rm inf} \geq m_\phi \sim \mathcal{O}(\text{TeV})$).

window $50 \leq N \leq 70$. In figure 3 we fix the number of e-foldings within the window, $50 \leq N \leq 70$, because at $N = 50$ and $N = 70$ the illustrated model satisfies the Planck 2σ combined constraints on the upper and lower bound of the amplitude of the power spectrum P_S, spectral tilt n_S, and the upper bound of tensor-to-scalar ratio $r_\star$ as mentioned in eq. (1.2), (1.3), (1.5) at the pivot scale $k_\star \sim 0.002$ Mpc^{-1} for both $c_s = 0.02$ and $c_s = 1$ branch.

Let us now derive an analytical expression for the scale of inflation, i.e. M_s. We consider a full cycle averaged within an interval $0 < t_{\rm osc} < H_{\rm inf}^{-1}$, and using eq. (5.13), eq. (5.15) and eq. (2.9) for n=6 flat directions, we can derive a following constraint on the scale of the heavy field, M_s for $k_\star \left(\sim 0.002\,\text{Mpc}^{-1}\right)$, by setting $\alpha \sim V(s) \approx M_s^4$ and the fine tuning parameter, $\delta \sim \mathcal{O}(10^{-4}) \ll 1$, the leading order contribution to the potential will be given by:

$$M_s \leq 1.77 \times 10^{16} \left(\frac{r_\star}{0.12}\right)^{1/4} c_s^{\frac{\epsilon_V}{2(\epsilon_V - 1)}}\ \text{GeV}\,. \tag{5.21}$$

The above eq. (5.21) will fix the upper bound on $M_s \sim \mathcal{O}\left(10^{16}\right)$ by setting $c_s = 1$.[9] Additionally, we also obtain a lower bound on M_s by considering the lower bound of the sound speed at, $c_s = 0.02$, which will generate very small value of the tensor-to-scalar ratio $r_\star \sim \mathcal{O}\left(10^{-22}\right)$ and satisfies the Planck observational constraints. Consequently we get the lower bound of the scale of the heavy field at, $M_s \sim \mathcal{O}\left(10^{11}\right)$ GeV.

[9] In the setup $\epsilon_V(k_* \approx k_{cmb}) \approx 0.0021$ which satisfies the WMAP+Planck constrain, as this combined data set puts an upper bound at $\epsilon_V < 0.008$at 95% CL. [7, 8] So for $0.02 \leq c_s \leq 1$, $(c_s)^{[\epsilon_V/2(\epsilon_V-1)]} \approx 1$ in the eq. (5.21). So the contribution in the scale of M_s comes from r_* and the prefactor sitting in the above eq. (5.21).

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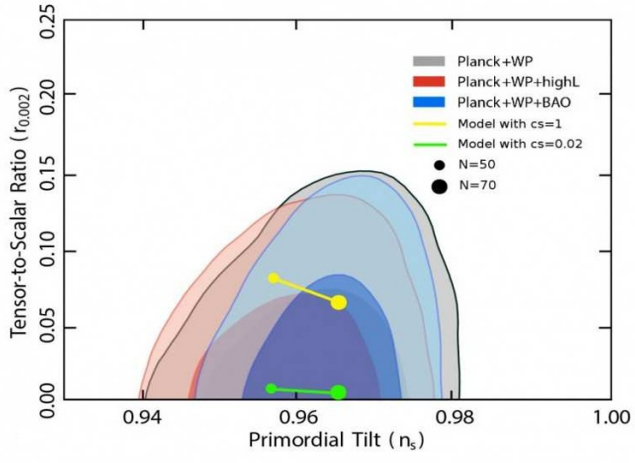

Figure 3. We show the joint 1σ and 2σ CL. contours using *Planck*+WMAP-9, *Planck*+WMAP-9+high l and *Planck*+WMAP-9+BAO data. The *yellow* and *green* lines are drawn for the proposed model with $c_s = 1$ and $c_s = 0.02$ where the model parameters are fixed at, $\delta \sim 10^{-4}$, $\lambda = 1$, $c_H = 2$, $a_H = 2.108$, $\phi_0 = 1.129 \times 10^{16}$ GeV, for the pivot scale $k_\star = 0.002\,\mathrm{Mpc}^{-1}$ respectively. The region in between the *yellow* and *green* lines represent the allowed region obtained from the model. The small circle on the left corresponds to $N = 50$, while the right big circle corresponds to $N = 70$.

At this point one might worry about the large vacuum energy density stored in the heavy field. This indeed helps inflation, in particular ameliorating the fine tuning parameter, we have taken $\delta \sim 10^{-4}$ in our scans [17, 78]. However such a large vacuum energy would need to be canceled after the end of slow roll inflation. In the string landscape [80], or in the MSSM landscape [56], it is plausible to have a bubble nucleation provided the rate of nucleation is large than the Hubble expansion rate. In the context of MSSM, these bubbles will naturally yield a low energy vacuum which is an *enhanced gauge symmetry point*, first suggested in ref. [56]. In the string vacua case, it is a challenge that the false vacuum governed by the heavy field s would nucleate to the MSSM vacuum [63]. Furthermore, the bubble nucleation could lead to an observational effects such as gravitational waves, etc. [81]. One may be able to constrain further the scale of heavy physics, M_s, from the high frequency gravitational waves, here we will not discuss these issues any further but we will leave this for future investigation. We can also envisage a smooth phase transition of the false vacuum as it can happen in the case of hybrid inflation [82, 83], possibly triggered by the MSSM inflaton itself as discussed in ref. [78]. In any of these scenarios we do not expect any modification on large scales, and therefore we do not expect these events to affect the primordial perturbations.

6 Conclusion

In this paper, we have shown that in any $\mathcal{N} = 1$ SUGRA inflation model when ever there are more degrees of freedom, non-minimal Kähler corrections would induce three distinct types of corrections: (i) non-minimal kinetic term for the inflaton, (ii) Hubble-induced mass correction to the inflaton, and (iii) Hubble-induced A-term in the potential.

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The exact nature of Kähler potential and Kähler corrections might not be known in all possible scenarios, but our aim has been to constrain the coefficients of the non-renormalizable Kähler higher dimensional operators phenomenologically, which are gauge invariant, from the recent Planck data. We assumed minimal Kähler potentials for all the fields to begin with. We first considered the heavy physics to be completely decoupled from the dynamics of the light inflaton field. We considered the light field to be embedded within MSSM, such that the reheating of the universe is guaranteed to be that of the SM dof. In the simplest setup when the heavy field is well settled down in its potential, it only affects via its vacuum energy density. The kinetic terms are mostly canonical, and therefore we do not obtain any constraint on the coefficients of the dimensional 3 and 4 non-reormalizable Kähler operators.

We further investigated an intriguing possibility, when the heavy field is coherently oscillating with a frequency larger than the Hubble parameter during the onset of inflation, while the light field is slowly rolling over the potential. In this particular scenario, we were able to constrain the coefficients of the Planck suppressed Kähler operators of dimensional 3 and 4. We scanned the four parameters, $a,\ b,\ , c\, , d$, and obtained a region of the parameter space where we can satisfy the current Planck observations, i.e. $P_S,\ n_S,\ c_s$ and $r_\star$ within 2σ CL, and we obtained all the coefficients to be of order $a,\ b,\ c,\ d \sim \mathcal{O}(1)$, as naturally expected in any non-rrenormalizable SUGRA theory. In fact, as we can see from figure 1 their magnitudes are always less than one.

In figure 2, we have shown for the range of non-renormalizable corrections, the parameter space for the allowed range of P_S versus n_S for the allowed range of $0.02 \leq c_s \leq 1$. In figure 3, we have plotted $r_\star$ vs. n_S, for $c_s = 1$ and $c_s = 0.02$ for the number of e-foldings, $N = 50,\ 70$. For the range of parameter space scanned, we were able to set an upper limit on the scale of new physics from the constraints arising from $r_\star$, which we obtained to be within $10^{11} \leq M_s \leq 10^{16}\,\mathrm{GeV}$. For the lower bound on M_s, we found $r_\star \sim \mathcal{O}\left(10^{-22}\right)$ and extremely negligible, and for the upper bound we saturated $r_\star = 0.12$. Note that the current Planck data mildly prefers lower value of the speed of sound, i.e. $c_s < 1$, this is visible from our scans and the plot on $r_\star$ versus n_S, see figure 3.

Finally, we would like to mention that all the above bounds have been obtained for a very particular kind of inflation model, which is fully embedded within MSSM, the inflaton is an MSSM flat direction and inflation happens at the point of inflection with a fine tuned parameter at the inflection point is roughly one part in 10^4. We chose MSSM inflation for its advantage that the dynamics can be well understood during inflation and after inflation. In particularly, we can ascertain that the universe after inflation would be filled with the SM degrees of freedom, and also the model is capable of explaining the Higgs mass constraint and the dark matter abundance, along with the constraints on the inflaton mass arising from the LHC [42–44]. Not every model of inflation enjoys such advantages, and therefore studying this model in some details along with SUGRA corrections yielded interesting constraints. Our methodology can be followed for other kinds of inflationary models too.

There is a further scope of improvement in our analysis. So far we have only used the Planck constraints from the power spectra, P_S, spectral tilt, n_S, tensor-to-scalar ratio, $r_\star$,

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and the constraint on the speed of sound, c_s. In principle we should be able to use the non-Gaussian parameters, $f_{NL}^{\rm local}$, $g_{NL}^{\rm local}$ and possibly $\tau_{NL}^{\rm local}$, to further constrain the non-renormalizable Kähler operators of dimension 3 and 4. In our companion paper, we would consider the non-Gaussian constraints in some details. All these cosmological constraints arising from Planck and future CMB missions can further improve our understanding of many different aspects of physics beyond the SM. With an improvement on tensor-to-scalar ratio, $r_\star$, we would be able to further constraint the scale of heavy physics, M_s.

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A XYZ

The symbols appearing in the eq. (5.9), in the definition of the sound speed c_s for $I \ll M_p$, after imposing the slow-roll approxiation are given by:

$$
\begin{aligned}
\mathbf{X}_1(t) &= \sqrt{\frac{2\epsilon_V(\phi)V(\phi)}{3}}\left\{\sqrt{\frac{2\epsilon_V(\phi)V(\phi)}{3}}\frac{aM_s^3}{M_p^2}\left[2\sin(2M_st)+4\cos(M_st)\right]\right.\\
&\qquad\left.-\frac{aM_s^4}{M_p^2}|\phi|\cos\mathbf{\Theta}\left[\cos(2M_st)-\sin(M_st)\right]\right\},\\
\mathbf{Y}_1(t) &= \sqrt{\frac{2\epsilon_V(\phi)V(\phi)}{3}}\left\{\sqrt{\frac{2\epsilon_V(\phi)V(\phi)}{3}}\frac{2bM_s^2}{M_p}\cos(M_st)+\frac{bM_s^3}{M_p}|\phi|\cos\mathbf{\Theta}\sin(M_st)\right\},\\
\mathbf{Z}_1(t) &= \sqrt{\frac{2\epsilon_V(\phi)V(\phi)}{3}}\left\{\sqrt{\frac{2\epsilon_V(\phi)V(\phi)}{3}}\frac{cM_s^3}{4M_p^2}\left[2\sin(2M_st)+4\cos(M_st)\right]\right.\\
&\qquad\left.-\frac{cM_s^4}{4M_p^2}|\phi|\cos\mathbf{\Theta}\left[\cos(2M_st)-\sin(M_st)\right]\right\},\\
\mathbf{W}_1(t) &= \sqrt{\frac{2\epsilon_V(\phi)V(\phi)}{3}}\left\{\sqrt{\frac{2\epsilon_V(\phi)V(\phi)}{3}}\frac{4dM_s^2}{M_p}\cos(M_st)+\frac{dM_s^3}{M_p}|\phi|\cos\mathbf{\Theta}\sin(M_st)\right\},\\
\mathbf{X}_2(t) &= \left(\mathbf{Y}_2(t)+\frac{a|\phi|^2M_s^5}{M_p^2}\sin(2M_st)\right),\\
\mathbf{Y}_2(t) &= \mathbf{Z}_2(t)=\mathbf{W}_2(t)=5M_s^5\sin(2M_st)+8M_s^5\cos(M_st),\\
\mathbf{X}_3(t) &= \left(\mathbf{Y}_3(t)-\frac{a|\phi|^2M_s^5}{M_p^2}\sin(2M_st)\right),\\
\mathbf{Y}_3(t) &= \mathbf{Z}_3(t)=\mathbf{W}_3(t)=3M_s^5\sin(2M_st)-8M_s^5\cos(M_st).
\end{aligned}
\tag{A.1}
$$

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Here the complex inflaton field ϕ is parameterized by, $\phi = |\phi| \exp(i\boldsymbol{\Theta})$. Here the new parameter $\boldsymbol{\Theta}$ characterizes the phase factor associated with the inflaton and it has a two dimensional rotational symmetry.

B Case -1, 2, 3, 4

Case–1: $\boldsymbol{K = \phi^\dagger\phi + s^\dagger s + \frac{a}{M_p^2}\phi^\dagger\phi s^\dagger s}$. For the above non-minimal Kähler interaction with $'a'$ being a dimensionless number. We have also computed the correction to the Hubble-induced mass term, for c_H for $|I| \ll M_p$:

$$c_H = \left\{3\left[(1-a)+(1+a)a\frac{|s|^2}{M_p^2}\right]+\left[(1+3a)+(1-3a)a\frac{|s|^2}{M_p^2}\right]\left(\frac{e^K|F_s|^2}{V(s)}-1\right)\right\} \approx 3(1-a)\,, \tag{B.1}$$

where we used the fact that: $V(s) = |W_s|^2 = 3H^2M_p^2 = 4M_s^2|s|^2$. Next we compute the correction to the Hubble-induced A term, $a_H H\frac{\phi^n}{nM_p^{n-3}}$, in presence of the non-minimal Kähler correction:

$$\begin{aligned}
&a_H H\frac{\phi^n}{nM_p^{n-3}} \\
&= \left(\left[1+a\frac{|s|^2}{M_p^2}\right]W_\phi\,\phi - 3W(\phi)\right)\frac{e^K W^*\left(I^\dagger\right)}{M_p^2} + \left[W(\phi)\frac{I^\dagger}{M_p} - aW(\phi)\frac{I^\dagger}{M_p}\frac{|I|^2}{M_p^2}\right. \\
&\quad \left. -aW_\phi\,\phi\,\frac{I^\dagger}{M_p}\left(1-a\frac{|s|^2}{M_p^2}\right)\right]\frac{e^K F_{\bar{s}}^*}{M_p} + \text{h.c.} \\
&\approx \left\{\left(1+a\frac{|s|^2}{M_p^2}\right)\left(1-\frac{3}{n}\right)\frac{s^2}{M_p^2} + \left(1-a\frac{|s|^2}{M_p^2}\right)\left(a-\frac{1}{n}\right)\frac{\left(s^\dagger\right)^2}{M_p^2}\right\}\frac{\lambda M_s\phi^n}{M_p^{n-3}} + \text{h.c.}\,,
\end{aligned} \tag{B.2}$$

which explicitly shows the Planck suppression for $|s| \ll M_p$ in the Hubble-induced A term.

Case–2: $\boldsymbol{K = \phi^\dagger\phi + s^\dagger s + \frac{b}{2M_p}s^\dagger\phi\phi + \text{h.c.}}$. Similarly, for the above non-minimal kähler correction where $'b'$ is a dimensionless number we can compute the correction to the Hubble-induced mass term, $c_H H^2|\phi|^2$, for $|s| \ll M_p$:

$$c_H = \left[3 + b^2\frac{e^K|W_s|^2}{H^2M_p^2} + \left(\frac{e^K|F_s|^2}{V(s)}-1\right)\right] \approx 3\left(1+b^2\right)\,, \tag{B.3}$$

where $V(s) = |W_s|^2 = 3H^2M_p^2 = 4M_s^2|s|^2$. And similarly the Hubble-induced A term, $a_H H\frac{\phi^n}{nM_p^{n-3}}$, in presence of snon-minimal Kähler correction read as:

$$
\begin{aligned}
& a_H H\frac{\phi^n}{nM_p^{n-3}} \\
&= \left(W_\phi\,\phi\, - 3W(\phi) + bW_\phi\,\phi^\dagger\frac{s}{M_p}\right)\frac{e^K W^*\left(I^\dagger\right)}{M_p^2} - \frac{b}{2}\frac{e^K W^*\left(I^\dagger\right)}{M_p^2}\left(\frac{W_s}{M_p} - \frac{s^\dagger}{M_p}\frac{W(I)}{M_p^2}\right)\phi\phi \\
&\quad + \left(W(\phi)\frac{s^\dagger}{M_p} - bW_\phi\,\phi^\dagger\right)\frac{e^K F_{\bar{s}}^*}{M_p} + 3bH^2\frac{s^\dagger}{M_p}\phi\phi + \text{h.c.} \\
&= \left\{\left(1-\frac{3}{n}\right)\phi + \frac{b\phi^\dagger s}{nM_p}\right\}\frac{\lambda\phi^{n-1}M_s s^2}{M_p^{n-1}} + \left(\frac{s^\dagger\phi}{M_p} - bn\phi^\dagger\right)\frac{2M_s\lambda\phi^{n-1}s^\dagger}{nM_p^{n-2}} + \frac{4M_s^2 b|s|^2 s^\dagger}{M_p^3}\phi\phi \\
&\quad - \frac{bM_s s^2}{2M_p^2}\left(\frac{2M_s s}{M_p} - \frac{M_s s^2 s^\dagger}{M_p^3}\right)\phi\phi + \text{h.c.}\,. \qquad \text{(B.4)}
\end{aligned}
$$

Case–3: $K = \phi\phi^\dagger + ss^\dagger + \frac{c}{4M_p^2}s^\dagger s^\dagger\phi\phi + \text{h.c.}$. In a similar way we can analyse the above non-minimal Kähler interaction, where c is the dimensionless number. We have computed the correction to the Hubble-induced mass term, $c_H H^2|\phi|^2$ for $|s| \ll M_p$ as:

$$
c_H = \left[3 + \frac{3c}{2}\frac{|s|^2}{M_P^2} + \left(1 + \frac{3c}{2}\frac{|s|^2}{M_P^2} - \frac{c^2}{4}\frac{|s|^4}{M_P^4}\right)\left(\frac{e^K|F_s|^2}{V(s)} - 1\right)\right] \approx 3 \qquad \text{(B.5)}
$$

where we have used $V(s) = |W_s|^2 = 3H^2M_p^2 = 4M_s^2|s|^2$. Next we compute the Hubble-induced A term, $a_H H\frac{\phi^n}{nM_p^{n-3}}$:

$$
\begin{aligned}
& a_H H\frac{\phi^n}{nM_p^{n-3}} \\
&= \left(W_\phi\,\phi - 3W(\phi) + \frac{c}{2}W_\phi\,\phi^\dagger\frac{ss}{M_p^2} - \frac{c}{2}\frac{s^\dagger}{M_p}\frac{W_s}{M_p}\phi\phi\right)\frac{e^K W^*\left(I^\dagger\right)}{M_p^2} + \left(W(\phi)\frac{s^\dagger}{M_p} - cW_\phi\,\phi^\dagger\frac{s}{M_p}\right)\frac{e^K F_{\bar{s}}^*}{M_p} \\
&\quad + \frac{3cH^2}{4}\frac{s^\dagger s^\dagger}{M_p^2}\phi\phi + \text{h.c.} \\
&= \left\{\left(1-\frac{3}{n}\right)\phi + \frac{c\phi^\dagger ss}{2M_p^2}\right\}\frac{\lambda\phi^{n-1}M_s s^2}{M_p^{n-1}} - \frac{cM_s^2 I^2 I^\dagger I\phi\phi}{M_p^4} + \left(\frac{s^\dagger\phi}{M_p} - \frac{cn\phi^\dagger s}{M_p}\right)\frac{2M_s\lambda\phi^{n-1}s^\dagger}{nM_p^{n-2}} \\
&\quad + \frac{M_s^2 c|s|^2 s^\dagger s^\dagger}{M_p^4}\phi\phi + \text{h.c.}\,. \qquad \text{(B.6)}
\end{aligned}
$$

Case–4: $K = \phi\phi^\dagger + ss^\dagger + \frac{d}{M_p}s\phi^\dagger\phi + \text{h.c.}$. For the above non-minimal Kähler potential, where d is the dimensionfull number, we can compute the Hubble-induced mass term, $c_H H^2|\phi|^2$, for $|s| \ll M_p$:

$$
\begin{aligned}
c_H = & \left[3\left[1 + d\frac{s+s^\dagger}{M_P} + d^2\left(1 + d\frac{s+s^\dagger}{M_P}\right)^{-1}\right] + \left[1 + d\frac{s+s^\dagger}{M_P}\right. \right. \\
& \left. \left. + 3d^2\left(1 + d\frac{s+s^\dagger}{M_P}\right)^{-1}\right]\left(\frac{e^K|F_s|^2}{V(s)} - 1\right)\right] \approx 3\left(1+d^2\right)\,, \qquad \text{(B.7)}
\end{aligned}
$$

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where we used $V(s) = |W_s|^2 = 3H^2M_p^2 = 4M_s^2|s|^2$. Next we compute the correction to the Hubble-induced A term, $a_H H \frac{\phi^n}{nM_p^{n-3}}$,

$$\begin{aligned} & a_H H \frac{\phi^n}{nM_p^{n-3}} \\ &= \left(W_\phi\,\phi - 3W(\phi)\right)\frac{e^K W^*\left(s^\dagger\right)}{M_p^2} + \left(W(\phi)\frac{I^\dagger}{M_p} - d\,W(\phi)\right)\frac{e^K F_{\bar{s}}^*}{M_p} + \text{h.c.} \\ &= \left(1-\frac{3}{n}\right)\frac{\lambda\phi^n M_s s^2}{M_p^{n-1}} + \left(\frac{s^\dagger}{M_p} - d\right)\frac{2M_s\lambda\phi^n s^\dagger}{nM_p^{n-2}} + \text{h.c.}\,. \end{aligned} \tag{B.8}$$

C Expression for a_H

Using these results in Hubble induced A-term, a_H can be computed from eqs. (B.2), eq. (B.4), eq. (B.6) and eq. (B.8) for the four physical situations, the simplified expressions turn out be:

$$a_H \sim \begin{cases} \frac{n}{2}\left(\frac{2}{3}\right)^{\frac{3}{4}}\sqrt{\frac{H_{\rm inf}}{M_p}}\left[1+a-\frac{4}{n}+\frac{35a}{4}\sqrt{\frac{2}{3}}\left(2-\frac{3}{n}\right)\frac{H_{\rm inf}}{M_p}-\frac{35a^2}{4}\sqrt{\frac{2}{3}}\frac{H_{\rm inf}}{M_p}\right] & \text{for } \underline{\textbf{Case I}} \\ \frac{1}{2}\left[3\left(1-\frac{1}{n}\right)+\frac{5b}{n}\sqrt[4]{\frac{2}{3}}\sqrt{\frac{H_{\rm inf}}{M_p}}\right]\left(\frac{2}{3}\right)^{\frac{3}{4}}\sqrt{\frac{H_{\rm inf}}{M_p}}+2\sqrt{\frac{2}{3}}\left(\left(\frac{3}{2}\right)^{\frac{3}{4}}\sqrt{\frac{H_{\rm inf}}{M_p}}-bn\right) & \\ +10b\left(\frac{3}{2}\right)^{\frac{5}{4}}\left(\frac{M_p}{\phi}\right)^{n-2}\left(\frac{H_{\rm inf}}{M_p}\right)^{\frac{3}{2}}-\frac{b}{2}\left(\frac{3}{2}\right)^{\frac{5}{4}}\left(\frac{M_p}{\phi}\right)^{n-2}\left(\frac{H_{\rm inf}}{M_p}\right)^{\frac{3}{2}}\left(5-\frac{67}{8}\sqrt{\frac{2}{3}}\frac{H_{\rm inf}}{M_p}\right) & \text{for } \underline{\textbf{Case II}} \\ \sqrt[4]{\frac{2}{3}}n\left[\sqrt{\frac{3}{2}}\left(1-\frac{3}{n}\right)+\frac{35c}{24}\frac{H_{\rm inf}}{M_p}\right]\sqrt{\frac{H_{\rm inf}}{M_p}}+(1-cn)\sqrt[4]{\frac{2}{3}}\sqrt{\frac{H_{\rm inf}}{M_p}} & \text{for } \underline{\textbf{Case III}} \\ \sqrt[4]{\frac{2}{3}}(n-3)\sqrt{\frac{H_{\rm inf}}{M_p}}+2\sqrt{\frac{2}{3}}\left[\left(\frac{3}{2}\right)^{\frac{3}{4}}\sqrt{\frac{H_{\rm inf}}{M_p}}-d\right] & \text{for } \underline{\textbf{Case IV}}. \end{cases} \tag{C.1}$$

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