Double-Running Inflaton Mass for a Flat Potential and Assisted Hilltop Inflation

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We propose a scenario to flatten the inflaton potential for a small-field slow-roll inflation from the origin. In the scenario, the inflaton mass-square goes through a renormalization group running that depends on not only inflaton field but also an assisting field. Thanks to the assisting field participating in the running, the initial condition for a slow-roll inflation can be set naturally, and the inflaton potential can be flattened by the vacuum expectation value of the field. Applying this idea, we propose a scenario of inflation, dubbed here as assisted hilltop inflation, which is free from the initial condition and the flatness problems of inflation.

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I. INTRODUCTION

As a modern paradigm of cosmology, inflation [1–3] provides a very simple solution to the problems of hot Big-Bang cosmology while providing a compelling way to generate seeds of the large-scale structure in the present universe [4,5]. Such an attractive property is based on so-called slow-roll approximation that requires a flat inflaton potential having a curvature smaller than the expansion rate by about one order of magnitude during inflation. However, the slow-roll approximation has been recognized in practice to be quite non-trivial to realize in high-energy theories, as on general grounds, the inflaton receives Hubble scale mass contribution(s) caused by Planck-suppressed operator(s). In particular, in the framework of supergravity, which is believed to be the low-energy effective theory of string theories and the proper UV framework for inflation, the problem manifests itself in the form of a supergravity (F-term) scalar potential [6,7]. The slow-roll approximation (especially the smallness of the inflaton mass-square relative to the expansion rate during inflation) clearly does not hold in this case, and the nearly scale-invariant CMB power spectrum cannot be explained. This is the well-known flatness problem of the inflaton potential, which is also known as the ‘η-problem’ from the second slow-roll parameter.

A variety of ideas have been put forward to resolve or avoid the η-problem in constructions of inflation models (see, for example, Ref. 8). In regard to F-term inflation subject to the η-problem, most of the models use special patterns of the Kähler potential and the superpotential under an assumption of a certain symmetry. Introducing a new symmetry (for example, shift symmetry [9]) may be a convenient way to remove dangerous gravitational corrections to the inflaton mass. However, whether any symmetry other than gauge symmetry, which is free from gravitational breakings, exists is not clear. A pseudo-Goldstone boson might be considered instead of the modulus of a complex field such as in the case of natural inflation [10]. However, as a content of a global symmetry, it is again not free from gravitational symmetry-breakings. In addition, even if a small η is realized somehow, such a scenario requires trans-Planckian excursions of the inflaton, which is generically problematic in view of the validity of a given theory at hand, because higher order terms suppressed by the Planck scale are out of control. A possible way of circumventing this issue may lie in the fact that generically the inflaton lives in a multi-dimensional field space. In other words, the inflaton doesn’t need to be a fixed (straight) coordinate, but can be just a non-trivial trajectory in the field space. In principle, the trajectory can be compactified in a sub-Planckian region even though its length is much larger than the Planck scale [11–20] (see also Ref. 21). Notably, the resulting inflation can be effectively the same as a single-field slow-roll inflation. However, note that the elongation of the trajectory again does require a tuning of parameters, so it may not look like an improving of the original problem in a natural manner. However, it provides, at least, a way of obtaining large-field inflation.
scenarios that are free from the trans-Planckian issue.
Considering a non-minimal coupling to gravity is another approach [22], although its original motivation was not the problem. Depending on setups, the Jordan frame potential can be flattened or amplified in the Einstein frame, and the $\eta$-problem becomes irrelevant. However, even if the question on the origin of the large non-minimal coupling were to be put aside, the issues of fine-tuning higher-order terms and unitarity would still exist [23] (see, however, Ref. 24 for example).

In the conventional Einstein gravity, when blocking dangerous gravitational corrections to the inflation mass in the framework of SUGRA is difficult, considering the smallness of $\eta$ as a result of some dynamics might be more natural. In this regard, an interesting and seemingly natural approach to the $\eta$-problem is to use the renormalization group running of the inflaton mass, as discussed in Ref. 25. The idea was accommodated in Hybrid inflation to match observations [26], but a recent analysis showed that several well-known models suffered from either a blue spectrum or density perturbations caused by entropy perturbation that were too large [27]. Instead of inward inflaton, outward direction may possibly be used for inflation, but in such a case, the field point corresponding to the pivot scale of observations is too far away from the point ending inflation, causing too much $e$-foldings to match observations. Such a difficulty may be resolved if a rapidly varying potential is added [28,29]. However, in order for the scheme to work, even though this issue of initial condition may not be problematic, those contributions to the potential (i.e., a field-dependent mass term and a rapidly varying higher-order term) should be properly adjusted at a specific field value associated with observations. Moreover, the beta-function of the inflation mass should be in the vicinity of a specific value that may or may not be realized.

In this work, inspired by the idea of running mass inflation (RMI) [25], we propose a new way to flatten the inflaton potential, in which two orthogonal supersymmetric flat-directions ($\Phi$ and $\Psi$) sharing some gauge interactions participate in the renormalization group running of the mass-squared parameter ($m_\Phi^2$) of $\Phi$. The novel aspect of this idea is that, without resorting to any ad hoc symmetry in the K"ahler potential, $m_\Phi^2$ can be dynamically adjusted to be small by using a vacuum expectation value (VEV) of the other flat direction $\Psi$. The running of $m_\Phi^2$ along $\phi$, the radial direction of $\Phi$, appears only after crossing a cutoff scale set by the VEV of $\Psi$. When a hilltop potential is added to shorten $e$-foldings to a right amount, this idea allows a natural realization of a hilltop inflation [31] along $\phi$ with a natural initial condition. We may call this scenario as assisted hilltop inflation.

This work is organized as follows: We first introduce the original idea of RMI briefly and then describe the RG-running of the inflaton mass along two flat directions. Subsequently, we will consider a concrete model of inflation and show the result of a numerical analysis as an example. Finally, the conclusion will be drawn.

II. FIELD-DEPENDENT RUNNING MASS

In the framework of supergravity (SUGRA), the defining energy scale of a theory or model is assumed to be the reduced Planck scale ($M_P \simeq 2.4 \times 10^{18}$ GeV), the fundamental scale of 4-D SUGRA. Now, suppose we are interested in an inflation with a quadratic potential of a scalar field $\phi$ representing the canonically normalized radial component of a flat direction $\Phi$: \[
V = V_0 + \frac{1}{2} m^2 \phi^2 + \ldots \tag{1}
\]
where $m^2$ is the soft SUSY-breaking mass-squared parameter of $\Phi$, and ... represents possible higher-order terms including terms stabilizing $\phi$ in the case of $m^2 < 0$ all the way to the Planck scale. $V_0$ is set for vanishing cosmological constant at zero-temperature true vacuum. At the defining scale, the general form of the K"ahler potential with $O(1)$ numerical coefficients leads to \[m^2 = O(3H^2),\tag{2}\]
where $H$ is the expansion rate of the universe when $V_0$ dominates the universe. Hence, the second slow-roll parameter, $\eta \equiv M_P^2 V''/V$ with $V''$ the second derivative with respect to $\phi$, appears to be of the order of unity with either sign allowed. This is the heart of the $\eta$-problem, and a slow-roll inflation is not realized. However, if $\phi$ has interactions with light (relative to the energy scale of interest) particles, the mass-squared parameter will get through renormalization group (RG) running so that $m^2$ depends on the renormalization scale ($Q$) or the energy scale of interest. Now, suppose particles coupled to $\phi$ obtain masses proportional to $\phi$ which is much larger than the scale of soft SUSY-breaking mass parameters (say $m_{\text{soft}}$). Then, if the strength of the couplings is assumed to be of the order of unity for simplicity, the renormalization will be cutoff at a scale determined by $\phi$. Hence, setting $Q = \phi$, one finds a field-dependent mass-squared $m^2(\phi)$. It should be noted that such a field-dependence of the mass-squared parameter is actually the dominant contribution to the 1-loop Coleman-Weinberg potential [32] along a flat direction for a general form of the K"ahler potential [33].

From the renormalization group equations for $m^2$ and the mass parameters of the coupled fields, in principle, one can find the form of $m^2(\phi)$. Here, for simplicity, we consider 1-loop RG-running of $m^2$ and ignore the scale dependence of the beta-function of the RG-equation. Then, one can write $m^2(\phi)$ as \[
m^2(\phi) = m_0^2 + \beta_m \ln \left( \frac{\phi}{M_P} \right), \tag{3}\]
where $\beta_m \equiv dm^2/d\ln Q$, the beta-function of $m^2$. We assume both $m^2_0$ and $\beta_m$ to be negative in the following discussion. In the vicinity of a field value $\phi_\lambda (\gg m_{\text{soft}})$, one can rewrite $m^2(\phi)$ as

$$m^2(\phi) = m^2_\lambda + \beta_m \ln \frac{\phi}{\phi_\lambda},$$

(4)

where $m^2_\lambda \equiv m^2(\phi_\lambda)$. In this approximation, ignoring the running of $\beta_m$ and higher order terms of $V$ in Eq. (17), one finds derivatives of $V$ as

$$V' = \left[ m^2_\lambda + \beta_m \left( \frac{1}{2} + \ln \frac{\phi}{\phi_\lambda} \right) \right] \phi,$$

(5)

$$V'' = \beta_m + \left[ m^2_\lambda + \beta_m \left( \frac{1}{2} + \ln \frac{\phi}{\phi_\lambda} \right) \right].$$

(6)

Thus, if $V'(\phi_\lambda) = 0$,

$$m^2_\lambda = -\frac{1}{2} \beta_m$$

(7)

and

$$V''(\phi) = \beta_m \left( 1 + \ln \frac{\phi}{\phi_\lambda} \right).$$

(8)

Typically, $\beta_m/m^2_0 = \mathcal{O}(10^{-2})$. Hence, if the field value $\phi_\lambda$ associated with the Planck pivot scale is close to $\phi_\lambda$, i.e., $\phi_\lambda \sim \phi_\lambda$, a flat enough power spectrum matching Planck data can be achieved in a natural way. Inflation can take place along either side of $\phi_\lambda$, but a recent analysis showed that the case of $\beta_m < 0$ with inflation for $\phi_\lambda < \phi$ fits Planck data better [34]. However, two issues still exist in this case: (i) the initial condition, i.e., setting $\phi \sim \phi_\lambda$ from the beginning and (ii) too much $e$-folding to the end of inflation. The former may require eternal inflation [35–39] or tunneling to populate $\phi$ around the local maximum of the potential. The latter requires an additional feature to reduce $e$-foldings.

III. DOUBLE RUNNING OF THE MASS PARAMETER

The RG-running of the inflaton mass-squared parameter depends on the masses of particles having a sizable coupling to the inflaton. In a simple case, their masses depend on the field value of inflaton as used in the RMI scenario. However, those particles may also interact with other fields which develop large non-zero VEVs, obtaining additional masses. If such a mass generation takes place before the last inflation along the radial direction of $\Phi$ takes place, there will be threshold effects or cutoff in the RG-running of the inflaton mass-squared.

As an example, we consider two orthogonal flat directions $\Psi$ and $\Phi$ consisting of gauge-charged fundamental fields and assume that those fields consisting of $\Psi$ share some gauge-interactions with the fields consisting of $\Phi$. At the defining scale, for a generic form of the K"ahler potential in SUGRA, $\Psi$ and $\Phi$ are expected to have a Hubble-scale soft SUSY-breaking mass-squared with the possibility of either sign. Those mass parameters are subject to renormalization-group (RG) runnings in the presence of gauge and Yukawa couplings if coupled particles are light relative to the energy scale of interest. We assume that $\Psi$ has a negative mass-squared parameter for the whole range of the field and is stabilized by a higher-order term. On the other hand, for $\Phi$, we assume that the mass-squared parameter is negative at the input scale (i.e., Planck scale) and subject to a significant negative RG-running as in the case of the RMI scenario. Then, ignoring term(s) depending on phases of $\Psi$ and $\Phi$ for simplicity, and denoting their canonically normalized scalar radial components as $\psi$ and $\phi$, respectively, we can write the potential involving them during inflation as

$$V = V_0 + V_\psi(\psi) + V_\phi(\psi, \phi),$$

(9)

with

$$V_\psi = \frac{1}{2} m_\psi^2 \psi^2 + \left( \frac{\lambda_\psi \psi^q - 1}{M_P^{q-3}} \right)^2,$$

(10)

$$V_\phi = \frac{1}{2} m_\phi^2 \phi^2 + \ldots,$$

(11)

where $\lambda_\psi$ is a dimensionless coefficient of $\mathcal{O}(1)$ at most, and $q \geq 4$ as an integer. The mass parameters are taken to be

$$m_\psi = c_\psi m_0^2,$$

(12)

$$m_\phi^2 = -c_\phi m_0^2 + \frac{1}{2} \beta_m \ln \left( \frac{\psi^2 + \phi^2}{m_{\text{soft}}^2} \right),$$

(13)

where $c_\psi \sim c_\phi \sim \mathcal{O}(1) > 0$ are numerical coefficients, $m_0^2 (< 0)$ represents a negative Hubble-scale mass-squared parameter defined at the Planck scale, $\beta_m \equiv dm^2_\phi/d\ln Q$ is determined by loop-corrections in a given model, and $m_{\text{soft}}$ is a typical mass scale induced by the supersymmetry-breaking effect of the energy density for inflation. Note again that the second term on the right-hand side of Eq. (13) is actually the dominant contribution to the 1-loop Coleman-Weinberg potential. Even though $m_\phi^2$ may start with a negative value, it can run to a positive value (which we set as $-c_\phi m_0^2$) around the origin due to a rather strong running with $\beta_m < 0$. We assume this to be the case. This kind of negative RG-running of $m_\phi^2$ depending on both $\phi$ and $\psi$ can arise when the RG-running is dominated by gauge interactions shared by both $\phi$ and $\psi$.

A plausible thermal history in our scenario is as follows: At high temperature, the interactions with thermal bath can hold $\psi$ and $\phi$ around the origin by providing them large thermal masses. As time goes on, due to its negative mass-squared around the origin, $\psi$ rolls out. This causes particles (for example, gauge/gaugino fields interacting with $\psi$) to obtain $\psi$-dependent masses and
causes RG-running of $m_\phi^2$ along $\psi$. As a result, if $\psi_0$ (the VEV of $\psi$) is far away from the origin, the local minimum around the origin along $\phi$ may turn to a local maximum as $\psi$ crosses $\psi_\chi$ satisfying

$$\beta_m \frac{\ln (\psi_\chi)}{m_{\text{soft}}} = 1$$

in the 1-loop approximation\(^2\). Note that, although not automatic, $\psi_0$ can be determined by the interplay of the quadratic term and a higher order term so as to be $\psi_0 \sim \psi_\chi$ under a reasonable choice of parameters. Then, in the vicinity of $\psi_\chi$, one can rewrite $m_\phi^2$ as

$$m_\phi^2 = \frac{1}{2} \beta_m \ln \left( \frac{\psi_\chi^2 + \phi^2}{\psi_\chi^2} \right).$$

For $\psi_0 > \psi_\chi$, as $\psi$ across $\psi_\chi$, $\phi$ starts rolling out slowly. For $\phi \ll \psi_0$, the curvature along $\phi$ is found to be

$$\frac{d^2V}{d\phi^2} \simeq \beta_m \left[ \ln \frac{\psi_\chi}{\psi_\chi} + 3 \left( \frac{\phi^2}{\psi_0^2} \right) \right];$$

i.e., the curvature along $\phi$ is dominantly controlled by the ratio $\psi_0/\psi_\chi$. Hence, in addition to a non-slow-roll inflation along $\psi$, a slow-roll inflation along $\phi$ may be realized with the spectral index under control. Note that in this type of inflation, the initial condition for inflation can be set naturally by an earlier stage of thermal inflation realized at the symmetry-enhanced point (i.e., the origin in the current scenario) in the field space.

In regard to the $\eta$-problem, the smallness of $\eta$ is now from the closeness of $\psi_0$ to $\psi_\chi$. This may be still regarded as a tuning, but $\psi_0$ can fall in the vicinity of $\psi_\chi$ for a reasonable choice of parameters, making the inflaton mass-squared small even when it is of the Hubble scale at the defining scale of the theory.

### IV. A MODEL

In the mechanism discussed in the previous section, for the last inflation along $\phi$, the role of $\psi$ can be simply replaced by a cutoff scale in the RG-running of the inflaton mass-squared. Keeping this in mind, we consider an inflaton potential along $\phi$:

$$V(\phi) = \frac{1}{2} m_\phi^2(\phi) \phi^2 + M^4 \left[ 1 - \left( \frac{\phi}{\phi_0} \right)^n \right]^2,$$

where

$$m_\phi^2(\phi) = m_\phi^2 + \frac{1}{2} \beta_m \ln \left( 1 + \frac{\phi^2}{\psi_0^2} \right),$$

with $m_\phi^2 \equiv \beta_m \ln (\psi_0/\psi_\chi)$ and $n \geq 4$ being an integer. The second term in Eq. (17) was added to avoid having too much $e$-folding as explained in Sections I and II. The scales of $M$ and $\phi_0$ are treated as free parameters here, but they are assumed to satisfy $\psi_0 \ll \phi_0 \ll M_\phi$. For $\phi_\ast$ associated with a cosmological scale in observations, if $\phi_\ast \ll \psi_0$, the current scenario will produce essentially the same inflationary observables as the well-studied model with the quadratic term replaced by a tree-level mass term [28]. However, $\psi_0(\sim \psi_\chi)$ is not arbitrarily large for a reasonable choice of $\beta_m$. Actually, expecting for it to be of an intermediate scale may be reasonable. In this case, the effect of the RG-running of $m_\phi$ along $\phi$ may have a sizable impact on inflationary observables. When the $\beta_m$ term in Eq. (18) is non-negligible, even if slow-roll parameters and inflationary observables from the potential in Eq. (17) are easily found, an analytic expression for the $e$-foldings is non-trivial to obtain. Also, note that the number of $e$-foldings required

\(^2\) Equation (14) will be modified when higher-order corrections are taken into account, which will result in a different value of $\psi_\chi$. However, such an alteration does not affect the subsequent argument as long as $\psi_0$ can be in the vicinity of $\psi_\chi$.\n
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**Fig. 1.** (Color online) Parameter space matching observations for $n = 5$, $M = 2.3 \times 10^{12}$ GeV, and $\phi_0 = 10^{15}$ GeV. **Left:** $\psi_0 = M$ with $\beta_m = 1/\ln(\psi_0/M_\phi)$. **Right:** $\beta_m = 0$. In both panels, green and blue bands are the 1- and the 2-$\sigma$ allowance of the spectral index $n_s$; the red line is for $N_e = N_e^{\text{obs}}$, and the purple dashed lines are for $10^5 \times \Delta n_s / \Delta \ln k = 7, 8, 9, 10$ from left to right. The vertical dashed line stands for $\phi_\ast$ in the case of a pure hilltop potential without a quadratic term. For the estimate of $\Gamma_\ast$ in Eq. (19), we took $\Gamma_\ast = 200$.\n
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to match observations is constrained by the energy density of inflation and the reheating temperature [41]. In order to take this into account, we take the decay rate of inflaton for $M \ll \phi_0$ to be

$$\Gamma_\phi = \frac{\gamma_\phi m^3_{\phi,0}}{8\pi^2 \phi_0^2},$$

(19)

where $\gamma_\phi$ is a numerical constant taking into account of decay channels and couplings to light particles, and $m_{\phi,0} \approx \sqrt{2\hbar M^2/\phi_0}$ is the mass at the true vacuum.

In Fig. 1, as an example, we depicted a parameter-space matching observations for $n = 5$ with and without $\beta_m$ contribution in Eq. (18). Comparing the two panels in the figure, one can see how large the impact of $\beta_m$ on the spectral index can be. It depends on how large or small $\psi_0/M$ is. As expected from the parametric dependences of the power spectrum and the $\epsilon$-foldings on $M$ and $\phi_0$, we find that for a larger $M$, while the change of the red line is minor, the blue line is pushed downward significantly, going out of the 2-$\sigma$ allowed band. Such a pattern can be compensated for by taking a larger $\phi_0$, but then the redline is pushed down significantly. Hence, not much room exists for adjusting $M$ and $\phi_0$ (which are inter-related) for a given set of the other parameters. On the other hand, taking a large $\psi_0$ pushes both the red and the blue lines toward the best-fit region. This is because, as $\psi_0$ becomes larger, the impact of mass-running in the vicinity of $\phi_0$ becomes weaker. For a larger $n$, e.g., $n = 6$, the allowed parameter space can fall into to the 1-$\sigma$ best-fit region of the spectral index.

It should be noted that in a concrete realization of a hilltop potential, $\phi_0$ may not be an arbitrarily free parameter. For example, the potential we are considering here might be from a superpotential of the form

$$W = S \left( M^2 - \frac{\lambda_\Phi \Phi^n}{nM^2} \right) \equiv M^2 S \left[ 1 - \left( \frac{\Phi}{\langle \Phi \rangle} \right)^n \right],$$

(20)

where $S$ is a gauge singlet superfield, $\Phi$ is assumed to be a hidden sector flat-direction something like a MSSM $D$-flat-direction, $\lambda_\Phi$ is a numerical coupling constant of the order of unity or less, $M_\phi$ is a UV-cutoff scale, which may be either GUT or Planck scale, and $\langle \Phi \rangle^n \equiv nM^2M_\phi^{n-2}/\lambda_\Phi$. Hence, $\phi_0(= \sqrt{2}\langle \Phi \rangle)$ depends on mainly $M_s$, then $M$ for $n \geq 5$. Such a form of the superpotential may come from, for example, a proper choice of gauged $U(1)_R$ symmetry [42,43]. From numerical tests, we found that for $\psi_0 \sim O(10^{12-14})$ GeV, a set of working parameters may be found for $M_s = M_{\text{GUT}}$, but finding such a set for $M_s = M_P$ would be difficult.

V. CONCLUSIONS

In this work, we proposed a simple mechanism to flatten the inflaton potential when gravitational corrections to the Kähler potential add a Hubble scale mass to the inflaton. This mechanism is based on the renormalization group running of the inflaton mass-square, but the running depends on not only inflaton ($\phi$) but also an assisting field ($\psi$), which are, respectively, radial components of flat directions made of several gauge-charged fundamental fields. The curvature of the potential along the inflaton direction can be adjusted to be very small mainly by the vacuum expectation value of the assisting field.

As a concrete realization of inflation utilizing this mechanism of the double-running inflaton mass, we considered a hilltop inflation scenario, dubbed assisted hilltop inflation, where the inflaton potential consists of both a quadratic term, whose mass parameter is subject to RG-running, and a typical hilltop potential in addition to a potential for the assisting field only. Contrary to known hilltop inflation scenarios without the running effect of the mass parameter taken into account, assisted hilltop inflation does not suffer from the initial condition and flatness problems. It requires just an initial hot universe. Assisted hilltop inflation may be regarded as a UV-completion of some hilltop inflation scenarios.

The idea of a double-running inflaton mass presented here does not automatically make the inflaton mass small relative to the expansion rate during inflation. However, it provides a simple way of obtaining a small curvature of the inflaton potential for a reasonable choice of parameters.

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