Twinflation

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Abstract

We consider a supersymmetric hybrid inflation model with two inflaton fields. The superpotential during inflation is dominated by $W = (\kappa S + \kappa' S') M^2$, where $S, S'$ are the (twin) inflatons carrying the same $U(1)_R$ charge, $\kappa, \kappa'$ are dimensionless couplings, and $M \sim 10^{15-16}$ GeV is a dimensionful parameter associated with a symmetry breaking scale. One light mass eigenstate drives inflation, while the other heavier mass eigenstate is stuck to the origin. The smallness of the lighter inflaton mass for the scalar spectral index $n_s \approx 0.963$, which is the center value of WMAP5, can be controlled by the ratio $\kappa'/\kappa$ through the supergravity corrections.

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I. INTRODUCTION

Although inflation seems to be inevitable in cosmology to resolve the homogeneous and flatness problems, it is highly nontrivial to realize the idea in the scalar field theory framework \[1\]. It is because a small scalar mass is perturbatively unstable, and a heavy inflaton mass term destroys the slow-roll conditions needed for sufficient e-folds (\(\gtrsim 50–60\)). Supersymmetry (SUSY) is helpful for keeping the smallness of the inflaton mass against quantum corrections, but it is just up to the Hubble scale. Supergravity (SUGRA) correction in de Sitter space usually induces a Hubble scale mass term of the inflaton at tree level even with the minimal Kähler potential, unless the model is carefully constructed. It is called the “\(\eta\) problem.”

One of the promising models, which potentially avoids the \(\eta\) problem, is the SUSY hybrid inflation model \[2, 3\]. In that model, the superpotential is dominated by \(W = \kappa SM^2\) during the inflation era, where \(S\) denotes the inflaton superfield, and \(\kappa\) and \(M\) are dimensionless and dimensionful parameters, respectively. \(M\) turns out to be associated with a symmetry breaking scale. With this superpotential, the SUGRA correction does not induce the dangerous Hubble scale inflaton mass term \((3H^2|S|^2)\) in the scalar potential, if the Kähler potential is given by the minimal form \((K = |S|^2)\) \[2, 4\]: Such a mass term is accidentally cancelled out at tree level in this model. However, a quartic term with the dimensionless coefficient of order unity in the Kähler potential, \(K \supset c_1|S|^4/4M_P^2\), where \(M_P\) is the reduced Planck mass (\(\approx 2.4 \times 10^{18}\) GeV), generates the unwanted inflaton mass term in the scalar potential. Actually, only the quartic term in the Kähler potential is dangerous, while higher order terms with coefficients of order unity are all harmless, because \(|S| \lesssim M_P\). Therefore, only the coefficient \(c_1\) should be assumed to be adequately suppressed (\(\lesssim 10^{-2}\)). This naive assumption needs to be justified by a UV theory or a quantum gravity theory in the future.

A remarkable feature in the SUSY hybrid inflation model is that the CMB anisotropy \(\delta T/T\) is proportional to \((M/M_P)^2\) \[3\]. Thus, the observational data of \(\delta T/T \sim 10^{-5}\) \[5\] determine the spontaneous symmetry breaking scale: \(M \approx 10^{15–16}\) GeV, which is tantalizingly close to the scale of SUSY grand unified theory (GUT) \[3\]. As a result, the SUSY hybrid inflation model can be embedded in the models of SUSY GUT. Indeed, this idea has been combined with the particle physics models of SU(3)\(c\)×SU(2)\(L\)×SU(2)\(R\)×U(1)\(_{B-L}\) \[6\], SU(4)\(c\)×SU(2)\(L\)×SU(2)\(R\) \[7\], SU(5)×U(1)\(_X\) \[8\], and SO(10) \[9\]. In those models, \(M\) is
interpreted as the $U(1)_{B-L}$ breaking scale.

In the SUSY hybrid inflation model, the scalar spectral index is predicted:

$$n_s \approx 1 + 2\eta \approx 1 - \frac{1}{N_l} \approx 0.98,$$

(1)

where $\eta \equiv M_P^2 V'' / V$ is the slow-roll parameter, and $N_l$ denotes the e-folding number (= 50–60). On the other hand, the recent WMAP 5-year (WMAP5) observation result on the scalar spectral index is $n_s = 0.963^{+0.014}_{-0.015}$. Thus, the prediction of the SUSY hybrid inflation model, $n_s \approx 0.98$ is quite deviated from the center value of the WMAP5 result. Indeed, unless relatively larger SUGRA corrections are included or the model is much ameliorated, the deviation can not be easily overcome.

The literatures attempted to explain the deviation by considering a small (but relatively larger) quartic [10, 11] term and/or a more higher order term [12] in the Kähler potential, or a quite small soft SUSY breaking “A-term” [13] in the scalar potential. In this letter, we will propose another way, in which the superpotential plays an essential role in explaining $n_s \approx 0.963$.

II. MODIFICATION OF HYBRID INFLATION

The SUSY hybrid inflation model is defined as the following superpotential [2, 3];

$$W = \kappa S (M^2 - \phi \bar{\phi}),$$

(2)

where $\phi$ and $\bar{\phi}$ are a conjugate pair of superfields carrying gauge and/or global charges. At the SUSY minimum, $S = 0$ and $|\phi| = |\bar{\phi}| = M$ by including the D-term potential, breaking a symmetry by the non-zero vacuum expectation values (VEVs) of $\phi$ and $\bar{\phi}$. Inflation starts when the inflaton $S$ is much deviated from the minimum, $S \gtrsim M$. Then, the complex scalars, $\phi$ and $\bar{\phi}$ achieve heavy masses, by which $\phi = \bar{\phi} = 0$ during inflation. It is the quasi-stable point. Thus, the superpotential is dominated by $W = \kappa S M^2$ during inflation. It provides the positive constant vacuum energy density $\kappa^2 M^4$, which gives rise to inflation. As mentioned above, the superpotential $W = \kappa S M^2$ and the minimal Kähler potential do not raise the “$\eta$ problem.” Since the higher order terms of the singlet $S$ in the superpotential destroys the slow-roll conditions, they should be forbidden by introducing the $U(1)_R$ symmetry. The triggering condition for inflation, $S \gtrsim M$ would be possible, if the universe was hot enough before inflation was initiated.
Because of the positive vacuum energy, SUSY is broken and so the constant scalar potential is quantum mechanically corrected. Neglecting the SUGRA corrections, thus, the scalar potential is given by

$$V \approx \kappa^2 M^4 \left(1 + \frac{\kappa^2}{16\pi^2} \log \frac{\kappa^2 |S|^2}{\Lambda^2}\right),$$  \hspace{1cm} (3)

where the logarithmic term denotes the quantum correction when $S \gtrsim M$, and $\Lambda$ means the renormalization scale. It makes a small slope in the potential, leading the inflaton to the SUSY minimum. As shown in Eq. (1), however, the scalar potential Eq. (3) yields $n_s \approx 0.98$, unless it is somehow modified.

Let us consider the following form of the modified inflaton potential;

$$V = \mu^4 \left(1 + \alpha \log \varphi + \frac{\delta}{2} \varphi^2\right),$$  \hspace{1cm} (4)

where $\mu^4$ is the positive vacuum energy density leading to inflation. The dimensionless field $\varphi$ denotes an inflaton scalar defined as $S/M_P$. The logarithmic term arises from the quantum correction caused by SUSY breaking [3]. Comparison with Eq. (3) yields the relations, $\mu^4 = \kappa^2 M^4$ and $\alpha = \kappa^2/8\pi^2$. In Eq. (4), the inflaton’s mass term is introduced: $V \supset (\delta/2)\mu^4 \varphi^2 = (3\delta/2)H^2 S^2$, where $H (= \sqrt{\mu^4/3M_P^2})$ is the Hubble constant during inflation. For successful inflation, thus, the dimensionless coupling $\delta$ should be small enough, $|\delta| \ll 1$.

The slow-roll parameter $\epsilon$ is still much smaller than $|\eta|$. It is basically because $\varphi$ is assumed to be smaller than 1. In the presence of the mass term in Eq. (4), the expressions of the e-folding number and “$\eta$” are given by

$$N_i = \frac{1}{2\delta} \log \left(1 + \frac{\varphi^2}{\alpha}\right), \hspace{1cm} \text{and} \hspace{1cm} \eta = \delta \times \frac{e^{2\delta N_i} - 2}{e^{2\delta N_i} - 1}. \hspace{1cm} (5)$$

We note that in the limit of $\delta \to 0$, the expression for $N_i$ and $\eta$ become $\varphi^2/2\alpha$ and $-1/2N_i$, respectively, which are the expressions given in the original form of the SUSY hybrid inflation model. In Eq. (5), the limit $\alpha \to 0$ does not make sense. It means that the logarithmic quantum correction makes an important contribution to $N_i$.

With the help of the inflaton mass term, the scalar spectral index can be compatible with the center value of WMAP5:

$$n_s \approx 1 + 2\eta \approx 0.963 \hspace{1cm} \text{for} \hspace{1cm} \frac{\delta}{2} = -3.0 \times 10^{-3}. \hspace{1cm} (6)$$
Here we set $N_l = 55$, but $n_s$ is quite insensitive to large $N_l$s. Since the sign of the quadratic mass term in Eq. (4) is negative, the potential is convex-upward. If the inflaton starts at a point of $V' > 0$ or $\alpha + \delta \cdot \varphi^2 > 0$, the inflaton can roll down eventually to the origin. It is fulfilled for $\kappa \gtrsim 5 \times 10^{-2}$ $(5 \times 10^{-3})$ and $\varphi \sim 0.1$ $(0.01)$. Actually, inflation would take place near the local maximum, “hilltop” [10], unless $\varphi \ll 1$.

The curvature perturbation is estimated as

$$P_{R}^{1/2} = 1 \frac{V^{3/2}}{\sqrt{12 \pi M_{P}^{3} \bar{V}'}} \approx \left( \frac{M}{M_{P}} \right)^{2} \sqrt{\frac{2|1 - e^{2\delta N_l}|}{3|\delta|e^{4\delta N_l}}},$$

(7)

where we set $\mu^2 = \kappa M^2$. For $N_l = 55$, $\delta = -6.0 \times 10^{-3}$, and $P_{R}^{1/2} \approx 4.91 \times 10^{-5}$ [5], $M$ is approximately $4.5 \times 10^{15}$ GeV, which is slightly lower than that in the case of $\delta = 0$ $(5.7 \times 10^{15}$ GeV). If $10^{12}$ GeV $\lesssim M \lesssim 10^{15}$ GeV, the curvature perturbation should be supplemented by another scalar field, “curvaton” [14]. Then, inflation don’t have to occur near the local maximum, since the room between $M \lesssim S$ and $S \ll M_{P}$ (or $\varphi \ll 1$) can be much larger. If $M \ll 10^{15}$ GeV, however, the inflationary scenario can not be embedded in a SUSY GUT model any more.

With the scalar potential Eq. (4), the fraction of the tensor perturbation is unlikely to be detectable in the near future.

III. TWINFLATION

For explaining the negative small inflaton mass squared, let us introduce one more inflaton $S'$. It carries the same quantum number as $S$, but has a mass different from that of $S$. In the presence of the twin inflaton fields \{S, S'\}, and two pairs of the waterfall fields \{\phi_1, \overline{\phi}_1\}, \{\phi_2, \overline{\phi}_2\}, the general superpotential takes the following form;

$$W = S \left( \kappa_1 M^2 - \kappa_1 \phi_1 \overline{\phi}_1 - \kappa_2 \phi_2 \overline{\phi}_2 \right) + S' \left( \kappa'_1 M'^2 - \kappa'_1 \phi_1 \overline{\phi}_1 - \kappa'_2 \phi_2 \overline{\phi}_2 \right),$$

(8)

where we assign the U(1)$_R$ charges of 2 (0) to $S$ and $S'$ (\phi$_{1,2}$ and \overline{\phi}$_{1,2}$) such that the higher power terms of $S$ and $S'$ are forbidden [15]. The different coupling constants $\kappa_{1,2}$ and $\kappa'_{1,2}$ distinguish $S$ and $S'$. We assume that $\kappa_{1,2}^{(0)}$ and $M_{1,2}^{(0)}$ are real quantities for simplicity. At the SUSY minimum, $\phi_{1,2}$ and $\overline{\phi}_{1,2}$ achieve heavy masses as well as the VEVs, satisfying $\phi_{1}^{*} = \overline{\phi}_1$ $(\neq 0)$, $\phi_{2}^{*} = \overline{\phi}_2$ $(\neq 0)$, and $\kappa_1 M^2 - \kappa_1 \phi_1 \overline{\phi}_1 - \kappa_2 \phi_2 \overline{\phi}_2 = \kappa'_1 M'^2 - \kappa'_1 \phi_1 \overline{\phi}_1 - \kappa'_2 \phi_2 \overline{\phi}_2 = 0$. Hence, both $S$ and $S'$ also get the heavy masses, and so $S = S' = 0$. 

5
Similar to the original SUSY hybrid inflationary scenario, inflation in this model can be initiated at a quasi-stable point of

\[
\begin{align*}
(k_1 S + k'_1 S')^2 &\gtrsim |\kappa_1^2 M^2 + \kappa'_1 \kappa'_2 M^2|, \quad \text{and} \\
(k_2 S + k'_2 S')^2 &\gtrsim |\kappa_2^2 M^2 + \kappa_1 \kappa_2 M^2|,
\end{align*}
\]

for which the tree level scalar potential is minimized at \(\phi_1 = \bar{\phi}_1 = \phi_2 = \bar{\phi}_2 = 0\). The left (right) hand sides of Eq. (9) come from the (off-) diagonal components of the mass matrices for \((\phi_1, \bar{\phi}_1)\) and \((\phi_2, \bar{\phi}_2)\). Thus, inflation is described by the following effective superpotential;

\[
W = (\kappa S + \kappa' S') M^2, \tag{10}
\]

where we redefined \(\kappa\) and \(\kappa'\) as

\[
\kappa \equiv \kappa_1 \quad \text{and} \quad \kappa' \equiv \kappa'_2 \left(\frac{M^2}{M^2}\right). \tag{11}
\]

We will assume a mild hierarchy between \(\kappa\) and \(\kappa'\), i.e. \(\kappa'/\kappa = (\kappa'_2 M^2/\kappa_1 M^2) \lesssim O(1)\), and \(M \approx 4.5 \times 10^{15} \text{ GeV}\). With Eq. (10), we obtain again the constant vacuum energy at tree level, breaking SUSY. So the logarithmic quantum corrections will be generated in the scalar potential as in the single inflaton case.

The Kähler potential is expanded with the power of \(S^{(i)}/M_P (\lesssim 1)\) up to the quartic terms as

\[
K = |S|^2 + |S'|^2 + c_1 \frac{|S|^4}{4M_P^2} + c'_1 \frac{|S'|^4}{4M_P^2} + c_2 \frac{|S|^2|S'|^2}{M_P^2} + \frac{c_3 |S|^2 + c'_3 |S'|^2}{2M_P^2} (SS'^* + S'^* S'), \tag{12}
\]

where \(c_i^{(o)} \ (i = 1, 2, 3)\) are dimensionless coefficients. The quartic terms’ coefficients of order unity in the Kähler potential would destroy the slow-roll condition of the inflation. In the original version of the SUSY hybrid inflation model, as mentioned in Introduction, the quartic term coefficient \(c_1\) in the Kähler potential, \(K = |S|^2 + c_1 |S|^4/4M_P^2 + \cdots\), is assumed to be suppressed \((\lesssim 10^{-3})\) in order to satisfy Eq. (14) as well as the slow roll condition. Along the line of it, we also assume one fine-tuned relation among the parameters of the Kähler potential, \(c_1, c_2,\) and \(c_3:\)

\[
c_1 + \frac{c_3}{1 - c_2} \lesssim O(10^{-3}). \tag{13}
\]

It can be satisfied, e.g. if they all are of order unity or smaller, but related to each other by \(c_1 \approx -c_3^2/(1 - c_2) \ [\text{Case (A)}]\), or if they (particularly \(c_1\) and \(c_3\)) are sufficiently suppressed, \(c_1, c_3^2/(1 - c_2) \lesssim O(10^{-3}) \ [\text{Case (B)}]\).
With Eqs. (10) and (12), the corrections coming from the scalar potential in SUGRA,

\[ V_F = e^{K/M_P^2} \left[ K^{-1}_{ij} D_i W (D_j W)^* - \frac{3}{M_P^2} |W|^2 \right] \]  

(14)
can be estimated. In our case, \( i, j = \{ S, S' \} \). \( K^{-1}_{ij} \) and \( D_i W \) stand for the inverse Kähler metric and the covariant derivative of the superpotential, respectively. Upto the quadratic terms, their components are approximately given by

\[ K^{-1}_{SS} \approx 1 + \frac{c_2^2 |S|^2}{(1 - c_2) M_P^2} - c_2 |S|^2 M_P^2 - \frac{c_3}{M_P^2} (SS^* + S^* S') , \]  

(15)

\[ K^{-1}_{S'S'} \approx 1 - c_1' |S'|^2 M_P^2 - c_3' |S|^2 M_P^2 - c_3 |S|^2 M_P^2 - \frac{c_3'}{M_P^2} (SS^* + S^* S') , \]  

(16)

\[ K^{-1}_{SS'S'} \approx -c_2 S S^* M_P^2 - c_3 |S|^2 M_P^2 - c_3 |S|^2 M_P^2 - \frac{c_3'}{M_P^2} (SS^* + S^* S') , \]  

(17)

\[ D_S W \approx M^2 \left( \kappa + \kappa' \frac{|S|^2}{M_P^2} \right) , \quad D_{S'} W \approx M^2 \left( \kappa' + \kappa' \frac{|S'|^2}{M_P^2} + \kappa SS^* \right) . \]  

(18)

In Eq. (15), we inserted Eq. (13). The scalar potential Eq. (14) is, thus, estimated as

\[ V_F \approx \kappa^2 M^4 \left\{ 1 + \frac{c_2^2}{1 - c_2} |x|^2 + (1 - c_2) |y|^2 - c_3 (x^* y + x y^*) \right\} \]

\[ + \kappa'^2 M^4 \left\{ 1 + (1 - c_2) |x|^2 - c_1 |y|^2 - c_3' (x^* y + x y^*) \right\} \]

\[ - \kappa' \kappa M^4 \left\{ (1 + c_2) (x^* y + x y^*) + 2 c_3 |x|^2 + 2 c_3' |y|^2 \right\} \]

\[ = (\kappa^2 + \kappa'^2) M^4 + \kappa^2 M^4 (x^* y^*) \mathcal{M} (x y)^T , \]  

(19)

where \( x \equiv S/M_P, y \equiv S'/M_P, \) and the mass matrix \( \mathcal{M} \) is given by

\[ \mathcal{M} = \begin{bmatrix} \frac{c_2}{1 - c_2} - \frac{c_3}{\kappa} c_3 + \frac{c_3'}{\kappa'} (1 - c_2) & - c_3 - \frac{c_3'}{\kappa'} (1 + c_2) + \frac{c_3}{\kappa} c_3' \\ - c_3 - \frac{c_3'}{\kappa} (1 + c_2) + \frac{c_3}{\kappa} c_3' & (1 - c_2) - \frac{c_3'}{\kappa'} (2 c_3' + \frac{c_3}{\kappa} c_3') \end{bmatrix} . \]  

(20)

Note that in the absence of the second inflatons’s contribution to the superpotential Eq. (10), namely, \( \kappa'/\kappa \to 0 \), one of the mass eigenvalues of \( \mathcal{M} \) is zero. Of course, if the relation \( c_1 = -\frac{c_3^2}{1 - c_2} \) is just slightly relaxed, then the small negative mass term of Eq. (4) can be supported purely by the Kähler potential. In this letter, however, we intend to acquire the effect with the help of the superpotential.

Case (A): If \( c_1 = -\frac{c_3^2}{1 - c_2} \) and \( \kappa'/\kappa \lesssim \mathcal{O}(1) \), the mass eigenstates and eigenvalues during inflation are

\[ \begin{pmatrix} \phi_L \\ \phi_H \end{pmatrix} \approx \frac{1}{D^{1/2}} \begin{pmatrix} 1 - c_2 & c_3 \\ - c_3 & 1 - c_2 \end{pmatrix} \begin{pmatrix} x \\ y \end{pmatrix} \]  

for \( m_L^2 \approx -\frac{\kappa'}{\kappa} \left( \frac{2 c_3 (2 - 2 c_2 + c_3 c_3')}{D} \right) \)

\[ m_H^2 \approx 1 - c_2 + \frac{c_3^2}{1 - c_2} , \]  

(21)
where $D \equiv (1 - c_2)^2 + c_3^2$.

Case (B): If $c_1, c_3^2/(1 - c_2) \ll \mathcal{O}(1)$, fulfilling Eq. (13), and $\kappa'/\kappa \lesssim \mathcal{O}(1)$, the mass eigenstates and eigenvalues are given by

$$
\begin{pmatrix}
\phi_L \\
\phi_H
\end{pmatrix} \approx \begin{pmatrix}
1 & \frac{\kappa' + c_3}{\kappa} \\
-\frac{\kappa' + c_3}{\kappa} & 1
\end{pmatrix} \begin{pmatrix}
x \\
y
\end{pmatrix} \quad \text{for} \quad \begin{cases}
m_L^2 \approx -\left(\frac{\kappa'}{\kappa}\right)^2 \frac{4c_2}{1 - c_2}, \\
m_H^2 \approx 1 - c_2.
\end{cases}
$$

(22)

In both cases, the mass squared of the heavier component, $\phi_H \times M_P$ is of the Hubble scale [$\sim \mathcal{O}(\kappa^2 M^4/M_P^2)$]. Consequently, it is expected to be stuck to the origin during inflation, $\phi_H = 0$. On the other hand, the mass squared of the lighter component, $\phi_L \times M_P$ can be much lighter than the Hubble scale, if $(\kappa'/\kappa)c_3$ for Case (A), or $(\kappa'/\kappa)^2 c_2$ for Case (B) is small enough. Therefore, inflation can be driven only by lighter mass eigenstate. Moreover, the sign of $m_L^2$ can be negative, if $c_3$ for Case (A) [or $c_2$ for Case (B)] is positive. $\phi_L$ can be identified with the $\varphi$ of Eq. (1), and $\mu^4 = (\kappa^2 + \kappa'^2)M^4$. Quantum correction by the coupling between $\phi_L$ and $\{\phi_{1,2}, \phi_{1,2}^c\}$ would induce the logarithmic term in Eq. (1), which leads $\phi_L$ eventually into the origin. $\phi_L = \phi_H = 0$ implies $S = S' = 0$. As $S$ and $S'$ approach the origin, Eq. (2) becomes violated, and then $\phi_{1,2}$ and $\phi_{1,2}^c$ also roll down to the absolute minima, developing VEVs. Hence, SUSY is recovered after inflation terminates.

Identification of $m_L^2$ with $\delta/2$ of Eq. (9) yields

$$
\frac{\delta}{2} \approx -3.0 \times 10^{-3} \approx \begin{cases}
-\frac{2c_3(2 - 2c_2 + c_3^2)}{(1 - c_2)^2 + c_3^2} & \text{for Case (A)}, \\
-\left(\frac{\kappa'}{\kappa}\right)^2 \frac{4c_2}{1 - c_2} & \text{for Case (B)}.
\end{cases}
$$

(23)

In Case (A), hence, $\kappa'/\kappa = (\kappa'^2 M^2)/(\kappa_1 M^2) \sim \mathcal{O}(10^{-3} - 10^{-1})$ fulfills the constraint for $c_3 \sim \mathcal{O}(1 - 10^{-2})$. Particularly, if all the quartic terms in the Kähler potential are suppressed, i.e. $c_i^{(l)}$ ($i = 1, 2, 3$) including $c_3$ are of order $10^{-2}$ or smaller, $\kappa'/\kappa$ of order $10^{-1}$ is necessary. In Case (B), $\kappa'/\kappa \sim \mathcal{O}(10^{-2} - 10^{-1})$ satisfies the constraint for $c_2 \sim \mathcal{O}(1 - 10^{-1})$. Thus, the mildly hierarchical $\kappa'$ and $\kappa$ couplings (or $\kappa'^2 M^2$ and $\kappa_1 M^2$) can generate the small negative inflaton’s mass squared, explaining $n_s \approx 0.963$.

IV. CONCLUSION

In this letter, we proposed a SUSY hybrid inflation model, in which one more singlet field carrying the same quantum number with the inflaton is introduced. Inflation is dominated by the superpotential, $W = (\kappa S + \kappa' S')M^2$, but only one linear combination of $S$ and $S'$
drives inflation. The smallness of $\kappa'/\kappa \sim O(10^{-1}–10^{-3})$ is responsible for the small negative mass squared of the inflaton needed for explaining $n_s \approx 0.963$.

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