Supergravity inflation with broken shift symmetry and large tensor-to-scalar ratio

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Abstract. We propose a class of inflation models with potential $V(\phi) = \alpha \phi^n \exp(-\beta^m \phi^m)$. We show that such kind of inflaton potentials can be realized in supergravity theory with a small shift symmetry breaking term in the Kähler potential. We find that the models with $(m = 1, n = 1)$, $(m = 1, n = 2)$, $(m = 2, n = 1)$, and $(m = 2, n = 3/2)$ can easily accommodate the Planck results. In particular, the tensor-to-scalar ratio is larger than 0.01 in the 1σ region. Thus, these models generate the typical large field inflation, and can be tested at the future Planck and QUBIC experiments.

Keywords: inflation, supersymmetry and cosmology

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1 Introduction

It is well-known that our Universe might experience an accelerated expansion, i.e., inflation \cite{1–4}, at a very early stage of evolution, as suggested by the observed temperature fluctuations in the cosmic microwave background radiation (CMB). Recently, the Planck satellite measured the CMB temperature anisotropy with an unprecedented accuracy. From its first-year observations, the scalar spectral index $n_S$, the running of the scalar spectral index $\alpha_s \equiv dn_s/d \ln k$, the tensor-to-scalar ratio $r$, and the scalar amplitude $A_s$ for the power spectrum of the curvature perturbations are respectively constrained to be \cite{5}

\begin{align}
    n_S &= 0.9603 \pm 0.0073, \\
    \alpha_s &= -0.0134 \pm 0.0090, \\
    r &\leq 0.11, \\
    A_s^{1/2} &= 4.6856^{+0.0566}_{-0.0628} \times 10^{-5}.
\end{align} \tag{1.1}

Furthermore, there is no sign of primordial non-Gaussianity in the CMB fluctuations. The generic predictions of the cosmological inflationary paradigm are qualitatively consistent with the Planck results. However, a lot of previously popular inflation models are challenged. For example, single field inflation models with a monomial potential $\phi^n$ for $n \geq 2$ are disfavoured due to large $r$. Interestingly, the Starobinsky $R+R^2$ model \cite{1, 6} predicts a value of $n_S \sim 0.96$, which is in perfect agreement with the CMB data, and a value of $r \sim 0.004$ that is comfortably consistent with the Planck upper bound \cite{5}.

From the particle physics point of view, supersymmetry is the most promising extension for the Standard Model (SM). In particular, the scalar masses can be stabilized, and the superpotential is non-renormalizable. Because gravity is also very important in the early Universe, it seems to us that supergravity theory is a natural framework for inflation model building \cite{7, 8}. However, supersymmetry breaking scalar masses in a generic supergravity theory are of the same order as the gravitino mass, giving rise to the so-called $\eta$ problem \cite{9–14}, where all the scalar masses are of the order of the Hubble parameter because of the large vacuum energy density during inflation \cite{15}. There are two elegant solutions: no-scale supergravity \cite{16–26}, and shift-symmetry in the Kähler potential \cite{27–35}.

The Planck satellite may measure the tensor-to-scalar ratio $r$ down to 0.03-0.05 in one or two years. The future QUBIC experiment targets to constrain at the 90% Confidence Level...
(C.L.) the tensor-to-scalar ratio of 0.01 with one year of data taking from the Concordia Station at Dôme C, Antarctica [36]. Interestingly, the well-known Lyth bound on tensor-to-scalar ratio is 0.01 [37]. Thus, the interesting question is how to construct inflation models which can be consistent with the Planck results and have large tensor-to-scalar ratio. For recent studies, see refs. [25, 33, 34, 38].

According to the Planck observations, the simple chaotic inflation models based on power law potentials \( \phi^n \) are out of 1\( \sigma \) region [5] due to the large tensor-to-scalar ratio. In this paper, we will modify the simple power law potentials by multiplying an exponential term as follows

\[
V(\phi) = \alpha \phi^n \exp(-\beta^m \phi^m),
\]

where the parameter \( \alpha \) relates to the energy scale of inflation, and \( \beta \) reveals the importance of exponential term. For simplicity, we take the Planck scale as unity \( M_{\text{Pl}} = 1 \) unless explicitly specified. These potentials with a constraint \( n = 4(1 - m) \) have been discussed in refs. [39–41] (For a review, see ref. [13].), where the expansion scale factor \( a(t) = \exp(\alpha \ln(t)\beta) \) is adopted. Besides, a concrete example, \( V(\phi) = \alpha \phi^2 \exp(-\beta^2 \phi^2) \) with a non-canonical kinetic term for \( \phi \) is obtained in ripple inflation from the modified no-scale supergravity [22, 25].

We show that the above inflaton potentials can be realized in the supergravity theory with a small shift symmetry breaking term in Kähler potential. Because we do not consider the non-renormalizable terms in Kähler potential, there are two cases for \( m \): \( m = 1 \) and \( m = 2 \). We find that the models with \( (m = 1, n = 1), (m = 1, n = 2), (m = 2, n = 1), \) and \( (m = 2, n = 3/2) \) can be highly consistent with the Planck observations. Especially, the tensor-to-scalar ratio is larger than 0.01 in the 1\( \sigma \) region, above the well-known Lyth bound [37]. Thus, these models produce the typical large field inflation, and can be tested at the future Planck and QUBIC experiments.

2 Supergravity inflation with shift symmetry breaking

The first step to realize inflation in supergravity is to get sufficient flat direction in the inflaton potential. The scalar potential in the supergravity theory with given Kähler potential \( K \) and superpotential \( W \) is

\[
V = e^K \left( (K^{-1})^{ij} D_i W D_j \overline{W} - 3 |W|^2 \right),
\]

where \( (K^{-1})^{ij} \) is the inverse of the Kähler metric \( K_{\overline{j} i} = \partial^2 K / \partial \Phi^i \partial \Phi_{\overline{j}} \), and \( D_i W = W_i + K_i W \). For the canonical Kähler potential \( K = \Phi \Phi^\dagger \), without miracle cancellation in eq. (2.1), the scalar potential \( V \) contains a term \( e^{\Phi^\dagger} \). Because it is very steep in the region with \( \Phi > 1 \), no slow-roll inflation can be realized for such potential. This is the so-called \( \eta \) problem, the major difficulty to achieve inflation in supergravity. The \( \eta \) problem can be naturally solved in the no-scale supergravity [16–20], in which the Kähler potential satisfies the flatness condition \( (K^{-1})^{ij} K_{ij} K_{\overline{k} i} K_{\overline{k} j} = 3 \), and the scalar potential can be arranged to be sufficient flat for inflation.

For general supergravity with a polynomial Kähler potential, normally the scalar potential varies significantly above the Planck scale. Interestingly, with a certain symmetry, one can decouple the inflaton from the Kähler potential. So the inflaton does not appear in the \( e^K \) term of the potential, and then no \( \eta \) problem in this direction. This fact, which provides another solution to \( \eta \) problem, realizes inflation in supergravity [27]. In such a case, the Kähler potential \( K \) is invariant under the shift symmetry \( \Phi \to \Phi + i C M_{\text{Pl}} \) with \( C \) a dimensionless real parameter, i.e., the Kähler potential \( K \) is a function of \( \Phi + \Phi^\dagger \) and
independent on the imaginary part of \( \Phi \). Consequently, it forms a flat direction along the imaginary part of \( \Phi \), and inflation may be triggered along this flat path.

Based on the fact that shift symmetry can keep the inflaton away from the \( \eta \) problem, several inflation models in supergravity were proposed [28–32]. As the inflation models with power law potentials are disfavored by the Planck observations, the supergravity inflation models with polynomial potentials have been studied recently [33–35]. In those works, the Kähler potentials are constrained by the shift symmetry in certain direction, while the superpotentials are generalized to be polynomial potentials.

In the following, we will obtain the potential \( \alpha \phi^n e^{-\beta \phi} \) in supergravity by introducing a new term in Kähler potential which slightly breaks the shift symmetry. Because the mass dimension of Kähler potential is two, we will not consider the models with \( m \geq 3 \) in this paper.

Before we study the general cases, a special potential with \( m = 1, n = 2 \) will be considered as a concrete example.

### 2.1 Supergravity realization of potential \( \alpha \phi^n e^{-\beta \phi} \)

The Kähler potential and superpotential used to get the potential \( \alpha \phi^n e^{-\beta \phi} \) are

\[
K = -b(\Phi + \Phi^\dagger) - \frac{1}{2}(\Phi - \Phi^\dagger)^2 + XX^\dagger, \\
W = a\Phi X ,
\]

in which the term \(-\frac{1}{2}(\Phi - \Phi^\dagger)^2\) has the shift symmetry \( \Phi \rightarrow \Phi + CM_{\text{Pl}} \). This symmetry is broken by the linear term \(-b(\Phi + \Phi^\dagger)\). As shown later, the shift symmetry breaking term introduces an exponential term \( e^{-\beta \phi} \) in the potential, which is crucial to fit the Planck results.

Actually, the model in eq. (2.2) is equivalent to the supergravity theory with canonical Kähler potential and superpotential as follows

\[
K = \Phi \Phi^\dagger + XX^\dagger, \\
W = ae^{-b\Phi - \frac{1}{2}\Phi^2} \Phi X .
\]

The new superpotential leads to the “miracle cancellation”, which plays the same role as the shift symmetry of Kähler potential in eq. (2.2) to solve the \( \eta \) problem. In addition, our model is equivalent to the supergravity model in refs. [31, 32] where the Kähler potential and superpotential are

\[
K = -\frac{1}{2}(\Phi - \Phi^\dagger)^2 + XX^\dagger, \\
W = ae^{-b\Phi} \Phi X .
\]

Especially, the shift symmetry is preserved in the Kähler potential. As we know, the superpotentials in eqs. (2.3) and (2.4) are not the traditional potentials in particle physics and string theory, and can not be generated by instanton effect or strong dynamics, etc. Thus, we do not consider these alternative explanations in this paper.

The scalar potential is

\[
V = a^2 \exp \left( -b(\Phi + \Phi^\dagger) - \frac{1}{2}(\Phi - \Phi^\dagger)^2 + XX^\dagger \right) \left( \Phi \Phi^\dagger + XX^\dagger - bXX^\dagger(\Phi + \Phi^\dagger) \right) + (1 + b^2)\Phi \Phi^\dagger XX^\dagger - XX^\dagger(\Phi^2 + \Phi^\dagger)^2 - \Phi \Phi^\dagger XX^\dagger(\Phi - \Phi^\dagger)^2 + \Phi \Phi^\dagger (XX^\dagger)^2 \right) .
\]
Decomposing the complex field $\Phi$ as follows

$$\Phi = \frac{1}{\sqrt{2}} (\phi + i\chi),$$

(2.6)

we obtain the scalar potential

$$V = a^2 e^{-\sqrt{2}b\phi + \chi^2 + XX^\dagger} \left( \frac{1}{2} (\phi^2 + \chi^2) + XX^\dagger - \sqrt{2}b\phi XX^\dagger + \frac{b^2 - 1}{2} \phi^2 XX^\dagger \right. + \left. \frac{b^2 + 3}{2} \chi^2 XX^\dagger + \chi^2 (\phi^2 + \chi^2) XX^\dagger + \frac{1}{2} (\phi^2 + \chi^2)(XX^\dagger)^2 \right).$$

(2.7)

So the potential $V$ depends on $\chi$ and $X$ through $e^{\chi\chi^\dagger}$ and $e^{XX^\dagger}$, as expected. While the shift symmetry breaking term $-b(\Phi + \Phi^\dagger)$ introduces a new term $e^{-\sqrt{2}b\phi}$ in the potential, which makes the potential flatter in the inflaton $\phi$ direction.

The masses of $\chi$ and $X$ with non-zero $\phi$ satisfy the following conditions

$$m^2 \chi = 2V + a^2 e^{-\sqrt{2}b\phi + \chi^2 + XX^\dagger} \left( 1 + (b^2 + 3)XX^\dagger + \cdots \right) > 6H^2,$$

$$m^2 X = V + a^2 e^{-\sqrt{2}b\phi + \chi^2 + XX^\dagger} \left( 1 - \sqrt{2}b\phi + \frac{b^2 - 1}{2} \phi^2 + 2\phi^2 XX^\dagger + \cdots \right),$$

(2.8)

where $H^2 = V/3$. Because $b$ is a small number for shift symmetry breaking, $m^2 \chi$ will be smaller than $3H^2$ if the magnitude of $X$ is smaller than $1/2$. To avoid having a very light scalar field $X$ during inflation, we can introduce a quartic term $-\xi(XX^\dagger)^2$ in the Kähler potential [31, 32]. Therefore, both $\chi$ and $X$ vanish fast and have little effects on the slow-roll inflation process. They can be safely fixed at the origin during inflation. The inflaton scalar potential in eq. (2.7) with $\chi = X = 0$ is simplified as follows

$$V = \frac{1}{2} a^2 \phi^2 e^{-\sqrt{2}b\phi},$$

(2.9)

which is the same as the potential proposed in eq. (1.2) with $(m = 1, n = 2)$ under the relations $\alpha = \frac{1}{2} a^2$ and $\beta = \sqrt{2}b$.

To get the general potential $\alpha\phi^n e^{-\beta\phi}$, we only need to use the same Kähler potential as above and a new superpotential

$$W = a\Phi^\frac{n}{2} X.$$  

(2.10)

It is easily shown that the extra fields $\chi$ and $X$ are still fixed at the origin during inflation, and then the inflaton scalar potential is

$$V = e^{-b(\Phi + \Phi^\dagger)} |W_X|^2 = 2^{-\frac{n}{2}} a^2 \phi^n e^{-\sqrt{2}b\phi}.$$  

(2.11)

2.2 Supergravity realization of potential $\alpha\phi^n e^{-\beta\phi^2}$

For the potentials with $m = 2$, we employ a different Kähler potential which has quadratic shift symmetry breaking term

$$K = -b^2(\Phi^2 + \Phi^\dagger 2) - \frac{1}{2} (\Phi - \Phi^\dagger)^2 + XX^\dagger,$$

(2.12)
Figure 1. The potentials $\alpha\phi^n\exp(-\beta\phi)$ and $\alpha\phi^n\exp(-\beta^2\phi^2)$ for various $n$. Here, the parameters $\alpha$ and $\beta$ are chosen to be consistent with the Planck observations, as shown later.

and the superpotential given by eq. (2.10). Similar to the Kähler potential in eq. (2.2), the above Kähler potential possesses the shift symmetry $\Phi \to \Phi + C M_{\text{Pl}}$ if $b = 0$, while this shift symmetry is slightly broken for a small but non-zero $b$.

For the Kähler potential in eq. (2.12), taking the complex field $\Phi = \frac{1}{\sqrt{2}}(\phi + i\chi)$, we obtain the exponential term in scalar potential

$$V \propto e^{-b^2\phi^2}(1+b^2)\chi^2+XX^\dagger. \quad (2.13)$$

All the other components in the potential are identical to the above scenario with $m = 1$. The mass of complex field $X$ is the same as that in eq. (2.8), while the mass of the real field $\chi$ is even larger

$$m_{\chi}^2 = 2(1+b^2)V + \cdots > 6(1+b^2)H^2. \quad (2.14)$$

Consequently, both $\chi$ and $X$ vanish during inflation. And the inflaton scalar potential is

$$V = e^{-b^2\phi^2}|W_X|^2 = 2^{-\frac{n}{2}}a^2\phi^n e^{-b^2\phi^2}. \quad (2.15)$$

Therefore, the inflation potential with $m = 2$ can also be realized in supergravity by introducing the shift symmetry breaking term in Kähler potential.

3 Slow-roll inflation

Based on the Planck results, inflation models with the pure power law potentials are disfavored: the $n_s - r$ relations with the e-folding number range $N \in [50, 60]$ are out of the $1\sigma$ region due to large $r$. With the exponential terms in the potentials, the $n_s - r$ relation will be significantly improved because of the following two points: (1) The potentials will become flatter due to exponential terms and more suitable for inflation; (2) The exponential term introduces one more free parameter $\beta$, which can be adjusted to fit the observation data.

The inflaton scalar potentials in eq. (1.2) with $m = 1$ and $m = 2$ are presented in figure 1 for various $n$. The inflation processes are assumed to be generated when the inflaton evolves from the local maxima to the origin.
3.1 Inflation models with $m = 1$

Taking $m = 1$, the inflaton potential turns into

$$V(\phi) = \alpha \phi^n \exp(-\beta \phi).$$  \hfill (3.1)

The inflation path is between the origin and the turning point $\phi_t = \frac{n}{\beta}$. The slow-roll parameters $\epsilon$ and $\eta$ are

$$\epsilon = \frac{(n - \beta \phi)^2}{2 \phi^2},$$

$$\eta = \frac{(n - \beta \phi)^2 - n}{\phi^2} = 2 \epsilon - \frac{n}{\phi^2}. \hfill (3.2)$$

During the inflation, the following slow-roll conditions should be satisfied

$$\epsilon \ll 1, \quad |\eta| \ll 1. \hfill (3.3)$$

At the end of inflation, we have either $\epsilon = 1$ or $|\eta| = 1$. In fact, the condition $\epsilon \ll 1$ will be violated first for small $n \leq 2$. From the condition $\epsilon = 1$, the exit point of inflation is determined as $\phi_f = n/(\sqrt{2} + \beta)$, and the corresponding $\eta$ is

$$\eta = 2 - \frac{(\sqrt{2} + \beta)^2}{n}. \hfill (3.4)$$

As we will see later, $\beta$ is very small to realize inflation processes that are consistent with the Planck observations. So for small $n \leq 2$, $\eta$ is always smaller than 1 before the inflation exit point. Its minimal value between $[\phi_f, \phi_t]$ is $-\beta^2 - 2\sqrt{2}\beta$ for ($n = 1$) or $-\frac{\beta^2}{n-1}$ for ($n > 1$), therefore, the slow-roll condition $|\eta| < 1$ is satisfied. While for $n = 3$, the condition $|\eta| \ll 1$ will be violated first, then the exit point of inflation $\phi_f$ is determined by the second slow-roll violation condition $|\eta| = 1$

$$\phi_f = \frac{1}{1 - \beta^2}(-3\beta + \sqrt{6 + 3\beta^2}) \cdot \hfill (3.5)$$

The e-folding number $N$ is

$$N = \int_{\phi_i}^{\phi_f} d\phi \frac{V}{V_{\phi} = -\frac{\phi}{\beta} - \frac{n}{\beta^2} \ln(n - \beta \phi)_{\phi_i}}. \hfill (3.6)$$

3.2 Inflation models with $m = 2$

The inflaton potential with $m = 2$ is

$$V(\phi) = \alpha \phi^n \exp(-\beta^2 \phi^2).$$  \hfill (3.7)

The turning point of this potential is $\phi_t = \sqrt{\frac{1}{2}} \beta$, and the slow-roll parameters $\epsilon$ and $\eta$ are

$$\epsilon = \frac{(n - 2\beta^2 \phi^2)^2}{2 \phi^2},$$

$$\eta = -2\beta^2 (1 + 2n) + \frac{n(n - 1)}{\phi^2} + 4\beta^4 \phi^2 = 2 \epsilon - 2 \beta^2 - \frac{n}{\phi^2}. \hfill (3.8)$$
The condition $\epsilon = 1$ or $|\eta| = 1$ determines the exit point of inflation. The first condition will be fulfilled at point $\phi_f = \frac{1}{2\sqrt{2}\beta} (-1 + \sqrt{1 + 4n\beta^2})$, and the corresponding $\eta$ is

$$\eta = 2 - \frac{1}{n} \left( 1 + \sqrt{1 + 4n\beta^2} \right).$$

(3.9)

So for $n \leq 2$, $\eta < 1$ at the exit point $\phi_f$. Besides, it is easy to check that in the range $\phi \in [\phi_f, \phi_i]$, the minimal value of $\eta$ is larger than $-1$. Thus, for the inflation models with $m = 2$ and $n \leq 2$, $\phi_f$ as given above is the exit point of inflation processes. The e-folding number $N$ is

$$N = \frac{1}{4\beta^2} \ln \frac{n - 2\beta^2 \phi_i^2}{n - 2\beta^2 \phi_f^2}.$$  

(3.10)

4 Numerical results

The $n_s - r$ relations for the inflation models with $m = 1$ and 2 are given in figure 2. In particular, for the models with $(m = 1, n = 1)$, $(m = 1, n = 2)$, $(m = 2, n = 1)$, and $(m = 2, n = 3/2)$, the predictions highly agree with the Planck observations. Notably, the tensor-to-scalar ratio is larger than 0.01 in the 1$\sigma$ region, above the well-known Lyth bound [37]. So these models generate the typical large field inflation. In the following years the Planck experiment may measure the tensor-to-scalar ratio down to 0.03-0.05, and the future QUBIC experiment may measure $r$ down to 0.01. Thus, it would be very interesting to compare the future observations with the predictions of our models.

For $m = 1$ and $m = 2$, the models with smaller $n$ are preferred. As can be seen from the two graphs in figure 2, the strips have a tendency to run out of the 1$\sigma$ region for larger $n$. For $m = 1$ and $m = 2$, the models respectively with $n > 4$ and $n > 3$ are out of the 1$\sigma$ region. So to be consistent with the observation data, $n$ cannot be too large. The strips for the number of e-folding $N$ from 50 to 60 end at certain lines, which correspond to the predictions of the chaotic inflation models with potential $\phi^n$ (See figure 1 in ref. [5]). In our models, when the parameter $\beta$ is close to zero, the exponential terms are ignorable. Thus, we return to the classical chaotic inflation models with potential $\phi^n$, which gives the upper limits for the predictions.
The relations between the parameter $\beta$ and inflation observables $n_s/r$ for $N = 55$ are given in figure 3. In general, to agree with the Planck observations in the 1σ region, the parameter $\beta$ should be small ($< 0.25$), but fortunately, $\beta$ is just close to the order of $O(10^{-1})$, no fine-tuning is needed here to fit the observation data. The $\beta - r$ relations in figure 3 are very interesting: to realize large $r$, $\beta$ should be very small, i.e., the exponential terms become less and less important when $r$ increases; in contrast, for the inflation with small $r$, the effects of the exponential terms becomes more important and non-ignorable.

Besides the $\phi^n$ models, the chaotic inflation can also be realized by the polynomial potentials [13]. Soon after the Planck results, the chaotic inflation based on the polynomial potential was studied [38], which can have large $r$ as well. The potential is $a\phi^2(b - \phi)^2$ and the predictions are consistent with the Planck observations. The results in ref. [38] are similar to our model with $\phi^2 e^{-\beta \phi}$ shown in figure 2. Actually, their potential can be considered as the approximation of our models with $(m = 1, n = 2)$

$$V(\phi) = a \phi^2 e^{-\beta \phi} \simeq a \phi^2 \left(1 - \beta \phi + \frac{1}{2} \beta^2 \phi^2 - \cdots\right),$$  

(4.1)

where $\beta$ is small and then all the non-renormalizable terms can be neglected. In other words, this approximation is valid for large $r$ and becomes worse for small $r$.

The second parameter $\alpha$ in the model determines the scale of inflation potential, which directly relates to the scalar amplitude $A_s$ for the power spectrum of the curvature perturbation

$$\alpha = \frac{3}{2} \pi^2 r A_s \phi_{i}^{-n} e^{\beta m \phi_{i}^m},$$  

(4.2)

where $\phi_{i}$ is the inflaton when the mode $k_{s}$ crosses the Hubble radius for the first time, and the parameter $A_s$, according to the Planck observation, is about $2.196 \times 10^{-9}$. From eq. (4.2), we can determine the regions for two parameters $\alpha$ and $\beta$ which agree with the Planck observations in the 1σ region. The results are given in figure 4. The parameter $\beta$ is of the order $10^{-1}$ or $10^{-2}$, which is natural as pointed out before. Another parameter $\alpha$ is of the order $10^{-10}$ for $n = 1$, and gets even smaller for the models with $n > 1$. In the supergravity inflation models, we have $\alpha = 2^{-n/2}a^2$. Thus, $a$ is of the order $10^{-5}$, and then it is not very fine-tuned for inflation.

5 Conclusion

We have studied a class of inflation models with potential $V(\phi) = \alpha \phi^n \exp(-\beta^m \phi^m)$, which can be realized in the supergravity theory with an small shift symmetry breaking term in the
Figure 4. Regions for $\alpha$ and $\beta$ in the inflation models $\alpha \phi^n e^{-\beta \phi}$ (left panel) and $\alpha \phi^n e^{-\beta^2 \phi^2}$ (right panel) that are consistent with the Planck observations within 1σ region. Here, $A_1^T$ is fixed to be $4.686 \times 10^{-5}$, and the effect of its uncertainty on the parameter $\alpha$ can be easily seen as they linearly depend on each other.

Kähler potential. We showed that the models with $(m = 1, n = 1)$, $(m = 1, n = 2)$, $(m = 2, n = 1)$, and $(m = 2, n = 3/2)$ have very good agreement with the Planck observations. Especially, the tensor-to-scalar ratio is larger than 0.01 in the 1σ region. Therefore, our models realize the typical large field inflation, and can be tested at the future Planck and QUBIC experiments.

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