Superconformal Subcritical Hybrid Inflation

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We consider D-term hybrid inflation in the framework of superconformal supergravity. In part of the parameter space, inflation continues for subcritical inflaton field value. Consequently, a new type of inflation emerges, which gives predictions for the scalar spectral index and the tensor-to-scalar ratio that are consistent with the Planck 2015 results. The potential in the subcritical regime is found to have a similar structure to one in the simplest class of superconformal $\alpha$ attractors.

I. INTRODUCTION

The observations of the cosmic microwave background (CMB) strongly support inflation as the paradigm of early universe. To discover the nature of inflation, intensive analysis of the CMB has been performed. The latest results by the Planck collaboration [11, 12] provide the bounds on the scalar spectral index $n_s$ and the tensor-to-scalar ratio $r$ of the primordial density fluctuations,

\[
s_s = 0.9655 \pm 0.0062 \ (68\% \ CL), \quad r < 0.10 \ (95\% \ CL).
\]

In fact, some inflation models, such as canonical chaotic inflation [3] and hybrid inflation [4], are already disfavored due to the bounds. Although they are not supported by the current observations, the models are simple and still attractive in theoretical point of view.

Recently Refs. [5, 6] studied the hybrid inflation in the framework of superconformal supergravity [7, 10]. It was found that the Starobinsky model [11] emerges in the supersymmetric D-term hybrid inflation [12, 14], to give a good accordance with the Planck observations. On the other hand, the D-term hybrid inflation was considered in a different context. In a shift symmetric Kähler potential [15], a ‘chaotic regime’ was found in the subcritical value of the inflaton field [16]. In the framework, inflation lasts even after the critical point of the hybrid inflation to give rise to different predictions from chaotic inflation. The following study [17] showed that the energy scale of inflation coincides with the Grand Unification (GUT) scale using the Planck 2013 data [18]. However, there is a tension between the predictions and the observations, especially the Planck 2015 data [11, 2].

In this letter we revisit D-term hybrid inflation in superconformal framework. It will be shown that there exists a single slow-rolling field in the subcritical value of the inflaton field. Since inflation continues for sufficiently long period, cosmic strings are unobservable as in Refs. [16, 17]. The potential in the subcritical region turns out to be in a general class of superconformal $\alpha$ attractors [19, 20], especially similar to the simplest version of the model. Consequently, non-trivial behavior and different predictions from the simplest ones are discovered.

II. SUBCRITICAL REGIME IN SUPERCONFORMAL D-TERM INFLATION

We consider D-term hybrid inflation in supergravity with superconformal matter [5, 6]. In the model three chiral superfields $S_\pm$ and $\Phi$, which have local $U(1)$ charge $\pm q \ (q > 0)$ and 0, respectively, are introduced. The superpotential and Kähler potential after fixing a gauge for the local conformal symmetry are respectively given by,

\[
W = \lambda S_+ S_- \Phi, \quad K = -3 \log \Omega^{-2},
\]

with,

\[
\Omega^{-2} = 1 - \frac{1}{3} (|S_+|^2 + |S_-|^2 + |\Phi|^2) - \frac{\chi}{6} (\Phi^2 + \bar{\Phi}^2),
\]

where $\lambda$ and $\chi$ are constants. The term proportional to $\chi$ in the Kähler potential breaks superconformal symmetry explicitly. In the model the Fayet-Iliopoulos (FI) term can be accommodated. Then, the D-term potential in the Einstein frame is [5],

\[
V_D = \frac{1}{2} g^2 \left( g \Omega^2 (|S_+|^2 - |S_-|^2) - \xi \right)^2,
\]

where $g$ is the gauge coupling and $\xi$ is the FI term, which is taken as a constant. (See Refs. [21, 25] for the subtleties of this issue in supergravity.) The F-term potential in the Einstein frame, on the other hand, is given in a simple form without exponentially growing terms [5, 22],

\[
V_F = \Omega^4 \lambda^2 \left[ |\Phi|^2 (|S_+|^2 + |S_-|^2) + |S_+ S_-|^2 - \frac{\chi^2 |S_+ S_- \Phi|^2}{3 + \frac{2}{3} (\Phi^2 + \bar{\Phi}^2) + \chi^2 |\Phi|^2} \right].
\]

As in the canonical hybrid inflation, $S_-$ is stabilized to its origin meanwhile $S_+$ suffers from the tachyonic instability depending on the field value of $\Phi$. The nature

\[1\] Throughout this letter we use the same notation for chiral superfields and scalar fields and take the reduced Planck mass $M_{pl} = 1$ unit.
of $\Phi$ depends on the value of $\chi$. In the Kähler potential there is a shift symmetry under $\mathrm{Re} \, \Phi (\mathrm{Im} \, \Phi) \rightarrow \mathrm{Re} \, \Phi (\mathrm{Im} \, \Phi) + \text{const.}$ for $\chi = -1 \pm 1$, and $\mathrm{Re} \, \Phi (\mathrm{Im} \, \Phi)$ can play a role of inflaton, as mentioned in Ref. \[\text{[5]}\].

We consider $\chi \leq -1$ in the later discussion without loss of generality. Then, the total potential is given by the waterfall field $s \equiv \sqrt{2} |S_+|$ and the inflaton field $\phi \equiv \sqrt{2} \mathrm{Re} \, \Phi$,

$$V_{\text{tot}}(\phi, s) = V_F + V_D$$

$$= \frac{\Omega^4(\phi, s)\lambda^2}{4} s^2 \phi^2 + \frac{g^2}{8} (q\Omega^2(\phi, s)s^2 - 2\xi)^2 ,$$

(7)

$$\Omega^{-2}(\phi, s) = 1 - \frac{1}{6} (s^2 + (1 + \chi)\phi^2) .$$

(8)

The waterfall field becomes tachyonic below the critical value $\phi_c$ of the inflaton field,

$$\phi^2_c = \frac{6qg^2\xi}{3\lambda^2 + (1 + \chi)qg^2\xi} .$$

(9)

After the tachyonic growth, the waterfall field is expected to reach its local minimum, which is obtained by $\partial V_{\text{tot}}(\phi, s)/\partial s = 0$,

$$s^2_{\text{min}} = \frac{2\xi\Omega^{-2}(\phi, 0)}{q(1 + \xi)} \frac{1 - \Psi^2}{q(1 + \xi)\Psi^2} ,$$

(10)

where $\xi \equiv \xi / 3q$ and,

$$\Psi = \frac{\Omega(\phi, 0)\phi}{\Omega(\phi_c, 0)\phi_c} = \frac{\Omega(\phi, 0)\phi}{\sqrt{2}qg^2\xi / \lambda^2} .$$

(11)

The expression for the local minimum given in Refs. \[\text{[16] \[\text{[17]}\] \text{[28]}}\ corresponds to the case for $\chi = -1$ (and $q = 1$) from the facts that $\Omega(\phi, 0)|_{\chi = -1} = 1$ and $\xi \sim O(10^{-4})$ in our targeted parameter space. Following Refs. \[\text{[16] \[\text{[28]}}\] (see also Appendix), we have confirmed numerically that the waterfall field reaches to the local minimum after $O(1/H_c)$ where $H_c = g\xi / \sqrt{6}$ is the Hubble parameter at the critical point, and then it becomes a single field inflation. Since the inflation lasts well over $O(10^2/H_c)$, cosmic strings, which are produced during the tachyonic growth, are unobservable. After the waterfall field relaxed to the local minimum, the dynamics of the inflaton is described by the potential,

$$V \equiv V_{\text{tot}}(\phi, s_{\text{min}})$$

$$= \frac{g^2\xi^2}{2\xi} (1 + \xi)\Psi^2 \frac{1 - \frac{g^2}{2(1 + \xi)}}{1 + 2\xi\Psi^2} .$$

(12)

As in Eq. \[\text{[10]}}\), it is easily to see that the potential $V$ with $\chi = -1$ agrees with one given in Refs. \[\text{[16] \[\text{[17]}\] \text{[28]}}\] up to $O(\xi)^2$.

\[\text{[2]}\] $q$ and $g$ can be absorbed by the redefinition of $\lambda$ and $\xi$, $\lambda \equiv \lambda / \sqrt{qg}$ and $\xi \equiv g\xi$ if we ignore terms proportional to $\xi$ that are irrelevant numerically. Although we will use $\lambda$ and $\xi$ in the following discussion, the results in terms of $\lambda$ and $\xi$ can be obtained by $q \rightarrow 1$, $g \rightarrow 1$, $\lambda \rightarrow \lambda$, and $\xi \rightarrow \xi$.\[\text{[2]}\]
The cosmological observables, i.e., the scalar amplitude $A_s$, the spectral index, and the tensor-to-scalar ratio, are then determined by,

$$A_s = \frac{V(\phi_s)}{24\pi^2(\phi_s)},$$

$$n_s = 1 + 2\eta(\phi_s) - 6\epsilon(\phi_s),$$

$$r = 16\epsilon(\phi_s).$$

We normalize the scalar amplitude by using the Planck 2015 data [2] $A_s = 2.198^{+0.076}_{-0.085} \times 10^{-9}$ and compute $n_s$ and $r$ for a given $N_\ast$.

As we have stated before, our target is the parameter space $\lambda \ll 1$. To search such a region, it is convenient to parametrize $\chi$ as,

$$\chi = -1 - \frac{3\lambda^2}{4g^2\xi} \delta \chi \quad (0 < \delta \chi < 1),$$

for the cosmological consequences. Fig. 1 shows the predictions of $n_s$ and $r$ in our model. Here $g = g_0 = 1$ is taken (see footnote 2), and $\lambda$ and $\xi$ are determined for a $\delta \chi$ and $N_\ast$ by using the scalar amplitude observed by the Planck collaboration. In Fig. 2, the allowed regions due to the bounds on $n_s$ and $r$ are shown for $N_\ast = 55–60$[4]. The upper and lower bounds on $\xi$ corresponds to the upper limit on $r$ and lower limit on $n_s$, respectively. In the $n_s$-$r$ plane, smaller values of $n_s$ and $r$ are obtained for larger $\lambda$ (and smaller $\xi$). In Fig. 2, the result in the previous work [17], i.e., the shift symmetric Kähler potential case, is also given as ‘shift sym.’ [4] We have checked that the result for $\delta \chi = 0$ agrees with it numerically and the similar behavior is seen around $\delta \chi \approx 0$. When $\delta \chi$ gets close to unity, on the contrary, a different behavior is observed. It is seen that $r$ gets smaller meanwhile $n_s$ tends to stay in the same value, which is within the Planck bounds. As a result, a wider allowed parameter space is obtained, which is seen in Fig. 2.

It is seen $\lambda \sim 10^{-4}–10^{-3}$ and $\sqrt{\xi} \sim 10^{16}$ GeV are consistent region with the Planck observation. Although the allowed region becomes larger, $\sqrt{\xi}$ tends to sit around the GUT scale even for $\delta \chi = 0.9$. As a consequence, the predicted $r$ is not extremely small. For example, $r > 0.0020$ ($0.075$) for $N_\ast = 60$ ($50$) for $\delta \chi = 0.9$. The value of $\chi$ in the allowed region, on the other hand, is found as $-1.41$ ($-1.016 < \chi < -1.0046$ ($-1.0092$) for $N_\ast = 60$ ($50$). Therefore, the parameter space $\lambda \ll 1$ and $\chi \sim -1$ is indeed favored by the observations.

In order to interpret the results, it is instructive to consider a canonically-normalized inflaton field $\phi$. Although the r.h.s of Eq. (14) is complicated, it can be approximated in the parameter space we are considering as,

$$\frac{d\phi}{d\hat{\phi}} \approx \sqrt{1 - \frac{1}{6}(1 + \chi)\phi^2}. \quad (19)$$

Then it is solved analytically,

$$\phi = \frac{1}{\sqrt{\beta}} \sinh \sqrt{\beta}(\hat{\phi} + C), \quad (20)$$

where $C$ is a constant and,

$$\beta = \frac{1 + \chi}{6} = \frac{\lambda^2}{2g^2\xi} \delta \chi = \frac{\delta \chi}{\phi^2\xi(1 - \delta \chi)}. \quad (21)$$

We have found that $C = 0$ is appropriate choice. Then $\Psi$ is simply given as,

$$\Psi \simeq \delta \chi^{-1/2} \tanh \sqrt{\beta} \hat{\phi}, \quad (22)$$

to express the potential in terms of $\hat{\phi}$,

$$V \simeq g^2\xi^2 \delta \chi^{-1}\tanh^2 \sqrt{\beta} \hat{\phi} \left[1 - \frac{\delta \chi^{-1}}{2} \tanh^2 \sqrt{\beta} \hat{\phi}\right]. \quad (23)$$

3 There is no allowed region for $N_\ast = 50$ except for $\delta \chi = 0.9$.
4 Do not confuse with the shift symmetric Kähler case with the present superconformal case where the shift symmetry is (weakly) broken in the Kähler potential.
This potential is valid in \( \hat{\phi} \leq \hat{\phi}_c = \frac{1}{\sqrt{2}} \sinh^{-1} \sqrt{3} \phi_c \). It is straightforward to check that the r.h.s is equal to \( g^2 \xi^2 / 2 \) for \( \hat{\phi} = \hat{\phi}_c \), and \( \hat{\phi}_c \to \infty \) for \( \delta \chi \to 1 \). We note that the potential coincides with a general class of superconformal \( \alpha \) attractors \[19\]. It especially resembles to the simplest class of the model,

\[
V_{\alpha\text{-attr.}} = \Lambda^4 \tanh^{2m} \frac{\hat{\phi}}{\sqrt{6} \alpha} .
\]

Due to the additional term, however, it has a different asymptotic behavior as we will see below.

In the small \( \lambda \) (and large \( \xi \)) region, \( \beta \) gets small, then the potential reduces to,

\[
V \simeq g^2 \xi^2 (1 - \delta \chi) \phi_c^2 \left[ 1 - \frac{1 + (4/3) \delta \chi}{2(1 - \delta \chi)} \phi_c^2 \right] .
\]

This is nothing but the potential for the shift symmetric Kähler case given in Refs.\[16\] \[17\] in the limit \( \delta \chi \to 0 \), which leads to \( \hat{\phi} \to \phi \). This feature is clearly seen in Fig.\[1\]. We note that the quadratic term is rewritten as \( (\lambda^2 \xi / 2g) \hat{\phi}^2 \), which is independent of \( \delta \chi \). Therefore, \( n_s \) and \( r \) approach to those in quadratic chaotic inflation in small \( \lambda \) limit (while \( \lambda^2 \xi \simeq \text{constant} \), independent of \( \delta \chi \)).

(Such a region is excluded, thus it is not shown in Fig.\[2\].) The potential \( V_{\alpha\text{-attr.}} \), on the other hand, has a similar structure,

\[
V_{\alpha\text{-attr.}} \simeq \Lambda^4 \hat{\phi}^2 \left[ 1 - \frac{m \phi_c^2}{9 \alpha} \right] .
\]

Although it coincides with \( V \) in the limit \( \hat{\phi} \to 0 \) for \( m = 1 \) and \( \Lambda^4 / 6 \alpha = \lambda^2 \xi / 2g \), it is not possible to get the same factor for the quartic term.

In large \( \lambda \) (and small \( \xi \)) region, on the contrary, \( \beta \) increases, which leads us to expand \( \Psi \) in large \( \sqrt{3} \phi \) limit to obtain,

\[
V \simeq \frac{1}{2} g^2 \xi^2 (2 - \delta \chi^{-1}) \left[ 1 + a_1 e^{-2 \sqrt{3} \phi} - a_2 e^{-4 \sqrt{3} \phi} \right] ,
\]

with \( a_1 = 8(1 - \delta \chi) / (2 \delta \chi - 1) \) and \( a_2 = 16(2 - \delta \chi) / (2 \delta \chi - 1) \). This expression should be compared with Eq.\[24\] in the \( \alpha \ll 1 \) limit. As shown in Ref.\[19\], it reduces to the potential in \( R^2 \) inflation \[29\] at large field value\[5\]

\[
V_{\alpha\text{-attr.}} \simeq \Lambda^4 \left[ 1 - 4me^{-\frac{2 \phi}{\sqrt{6} \alpha}} \right] .
\]

Now it is clear that the form of the potential with \( \delta \chi = 1 \) (in large \( \lambda \) region) reduces to \( R^2 \) inflation, or the simplest class of superconformal \( \alpha \) attractors in \( \alpha \ll 1 \) limit. To be specific, a choice of \( \Lambda^4 = g^2 \xi^2 / 2 \) and \( \alpha = 1 / 24 \beta \) leads to the same asymptotic form. Then, we get \( n_s \approx 1 - 2 / (N_\star + 1) - 3g^2 \xi / 8\lambda^2 (N_\star + 1)^2 \), \( r \approx g^2 \xi / \lambda^2 (N_\star + 1)^2 \), while satisfying \( \lambda^2 \xi \simeq \text{constant} \). Namely, when \( \lambda \) increases \( n_s \) approaches to \( 1 - 2 / (N_\star + 1) \) and \( r \) gets smaller and smaller. We have confirmed this behavior using Eq.\[23\] with \( \delta \chi = 1 \). Recall that, however, the critical value becomes infinity, which is unphysical.

Such a behavior, on the contrary, can not be seen for \( \delta \chi \neq 1 \) case shown in Fig.\[4\]. This arises from non-zero \( a_1 \) in Eq.\[27\]. This is why we have seen the different cosmological consequences.

In Fig.\[3\] the potential as function of canonically-normalized inflaton field \( \hat{\phi} \). \( \delta \chi = 0.9 \) and 1 cases (‘superconf.’) are shown, which are compared with superconformal \( \alpha \) attractors (‘\( \alpha \) attr.’), \( R^2 \) inflation (‘\( R^2 \)’), and the shift symmetric Kähler case (‘shift sym.’). The field values \( \hat{\phi}_T \) and \( \hat{\phi}_s \) at the end of inflation and the last 60 e-folds, respectively, are also indicated for \( \delta \chi = 0.9 \) and 1 cases.

\[5\] To be precise, \( \alpha = 1 \) gives the original \( R^2 \) inflation. The factor \( 4m(>0) \) is quantitatively irrelevant for the slow-roll predictions.
quantities are different, i.e., $n_s = 0.967$ and $r = 0.00044$. It is clear, on the other hand, that $\delta \chi = 0.9$ case shows a different behavior from the others. To summarize, the model has a nature of both the shift symmetric Kähler case and the simplest superconformal $\alpha$ attractors, and the slow-roll predictions change accordingly.

IV. CONCLUSION

We have revisited superconformal D-term hybrid inflation. After reaching its critical value, the inflaton field is slowly rolling thus inflation continues for a small coupling $\lambda$ of inflaton to the other fields. Because of a sufficiently long period of slow-roll regime, cosmic strings, which are formed during the tachyonic growth of the waterfall field, are unobservable. The potential which determines the dynamics of the canonically-normalized inflaton in the subcritical regime has been found to resemble to the simplest version of superconformal $\alpha$ attractors but with an additional term. Consequently, different predictions for the slow-roll parameters are obtained. For $\lambda \sim 10^{-4} - 10^{-3}$ and $\sqrt{\xi} \sim 10^{16}$ GeV, $n_s$ and $r$ are consistent with the Planck data.

The predictions depend on a parameter $\chi$ that explicitly breaks superconformal symmetry in the Kähler potential. In addition, the Kähler potential with $|\chi| = 1$ has a shift symmetry for the inflaton field, which is explicitly broken by non-zero $\lambda$ in the superpotential. On the other hand, $|\chi| \simeq 1$ is required from the consistency of the model setting, thus $\lambda \ll 1$ is parametrically natural. It has been found that the observational bounds indeed prefer such a parameter space.

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Appendix: Tachyonic growth of waterfall field

Around the critical value of the inflaton field, the dynamics of the waterfall field is governed by its tachyonic growth. For the evaluation, we define canonically-normalized waterfall field $\hat{s}$ near the critical point,

$$\frac{ds}{d\hat{s}} \simeq K_{S_+^+ S^-_+}^{-1/2} (\phi_c, 0). \quad (A.29)$$

The rest parts are parallel to Refs. [16, 17]. We expand the potential $V_{tot}(\phi, s)$ near the critical point, i.e., $\phi \simeq \phi_c + \hat{\phi} c t$ as

$$V_{tot}(\phi, s) = \frac{g^2 \xi^2}{2} - \frac{3}{2} t s^2 + O(t^2, s^4), \quad (A.30)$$

to leads the interaction term in the equation of motion of $\hat{s}$,

$$\frac{\partial V_{tot}(\hat{s}, s)}{\partial \hat{s}} = K_{S_+^+ S^-_+}^{-1/2} (\phi_c, 0) \frac{\partial V(\phi, s)}{\partial \phi} = -\hat{d}^3 \hat{s} + O(t^2, s^3), \quad (A.31)$$

where $\hat{d}^3 = (2 g^2 \xi^2) |\phi_c| / \lambda^2 \hat{s}$ and $|\phi_c| = -(\partial V_{11}(\phi) / \partial \phi) / (3 H_c) = \sqrt{\log 2} g^2 \gamma^2 \lambda^2 \xi / 4 \pi^2 \phi_c (3 \lambda^2 + g^2 \xi (1 + \chi^2))^2$. The equation of motion for $\hat{s}$ gives rise to that of momentum mode $\hat{s}_k$ of the quantum fluctuation [16, 17, 28].

$$\hat{s}_k + \left[ k^2 e^{-2 H_c t} - \frac{9}{4} H_c^2 - \hat{d}^3 t \right] \hat{s}_k = 0. \quad (A.32)$$

Here we have used $V_{11}$ given in Ref. [5]. Solving the equation numerically, we obtain the variance $\langle \hat{s}_k^2(t) \rangle$.

After the decoherence time $t_{dec}$ we match the variance with the classical motion of the waterfall field as $s(t_{dec}) = K_{S_+^+ S^-_+}^{-1/2} (\phi_c, 0) \sqrt{\langle \hat{s}_k^2(t_{dec}) \rangle}$, and solve the equations of motion for $\phi$ and $s$

$$\begin{align*}
3 H \dot{\phi} &= -K_{S_+^+ S^-_+}^{-1} (\phi, 0) \frac{\partial V(\phi, s)}{\partial \phi}, \quad (A.33) \\
3 H \dot{s} &= -K_{S_+^+ S^-_+}^{-1} (\phi, 0) \frac{\partial V(\phi, s)}{\partial s}. \quad (A.34)
\end{align*}$$

We have confirmed numerically that the obtained solutions coincide with the dynamics described by $V$ in Eq. (12).

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