Alchemical Inflation: inflaton turns into Higgs

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Abstract

We propose a new inflation model in which a gauge singlet inflaton turns into the Higgs condensate after inflation. The inflationary path is characterized by a moduli space of supersymmetric vacua spanned by the inflaton and Higgs field. The inflation energy scale is related to the soft supersymmetry breaking, and the Hubble parameter during inflation is smaller than the gravitino mass. The initial condition for the successful inflation is naturally realized by the pre-inflation in which the Higgs plays a role of the waterfall field.
1 Introduction

The existence of the inflationary era \cite{1} in the early Universe is strongly suggested by the observations \cite{2}. In particular, the density perturbations extending beyond the Hubble horizon at the last scattering surface can be interpreted as the evidence for the accelerated expansion of the Universe in the past.

While there are many inflation models, we still do not know which inflation model is realized in nature. The study of density perturbations such as isocurvature perturbations, non-gaussianity, tensor-mode, and their effects on the CMB power spectrum is a powerful diagnostic of the mechanism that laid down the primordial density fluctuations, but it is not enough at present to pin down the inflation model. This is partly because of our ignorance of thermal history of the Universe beyond the standard big bang theory, especially, how the inflaton reheated the Universe.

In the reheating process the inflaton transfers its energy to the visible sector. Usually it is assumed that it proceeds through perturbative or non-perturbative decay of the inflaton. After reheating, the visible sector is thermalized sooner or later, leading to the hot radiation dominated Universe.

The reheating process is also useful to constrain theory beyond the standard model (SM) because unwanted relics such as gravitinos \cite{3,4} are efficiently produced at the reheating. Recently there was progress in this regard; the inflaton with a finite vacuum expectation value (VEV) is necessarily coupled to the SM sector at both tree and one-loop levels \cite{5,6}, and so, the inflaton naturally decays and reheats the SM sector. However it turned out that such inflaton is also coupled to the supersymmetry (SUSY) breaking sector, and the inflaton decay generically produces gravitinos as well \cite{7,8,6,9}. The abundance of the non-thermally produced gravitinos is inversely proportional to the reheating temperature, and so, it tightly constrains the inflation models, especially if combined with the thermal production of the gravitinos.

In this letter, we propose a new inflation model, in which a SM gauge singlet inflaton turns into the Higgs field (or more precisely, $H_u H_d$) after inflation. The energy stored in the inflaton is directly transferred to the $H_u H_d$ condensate, where $H_u$ and $H_d$ are
up- and down-type Higgs fields, respectively. This is a new type of reheating; instead of producing the SM particles from the inflaton decay, the inflaton energy is directly transmuted into the $H_uH_d$ flat direction. As we will see shortly, the transmutation is due to the inflaton dynamics in the moduli space of supersymmetric vacua spanned by the inflaton and the Higgs field. Interestingly, because of this transmutation, the new inflation ends up in a symmetry vacuum of the Higgs boson, which enables rapid thermalization of the Higgs condensate. The inflation model has the following several cosmological and phenomenological virtues.

- The inflaton potential is induced by the SUSY breaking, and so, the SUSY breaking scale is directly related to the normalization of the density perturbations.
- Non-thermal gravitino production does not occur, both because the Higgs flat direction has a small VEV and because the soft mass for $H_uH_d$ is comparable to the gravitino mass.
- The inflation scale is lower than the gravitino mass by more than one order of magnitude:

$$H_{\text{inf}} \ll m_3^{3/2}. \quad (1)$$

This is advantageous from the point of view of the moduli stabilization [10].

- The initial condition of the inflation can be naturally realized by the primordial hybrid inflation [11] or its variant, the smooth hybrid inflation [12].

It is also possible to replace the role of the $H_uH_d$ with other D-flat directions of the SUSY SM (SSM). Then, the Affleck-Dine (AD) mechanism [13, 14] naturally works. What is interesting is that the AD field plays an important role in the inflationary dynamics, and its initial large VEV is a requisite for the inflation to take place.

Our model has similarity with the so-called MSSM inflation [15] in that the inflaton potential is induced by the SUSY breaking and the energy density after inflation is dominated by coherent oscillations of a D-flat direction in the SSM. However, there are great differences: (1) the inflation is mainly driven by a gauge singlet field, not the MSSM flat direction; (2) the energy transfer of the inflaton to the visible sector is due to non-trivial
dynamics in the moduli space; (3) the AD mechanism works; (4) the initial condition can be dynamically set by the pre-inflation.

The rest of this letter is organized as follows. In Sec. 2 we give the inflation model and study its dynamics in detail. The various implications of the model are discussed in Sec. 3. The last section is devoted for conclusions.

2 Inflation Model

We consider the following superpotential:

\[ W = S(\mu^2 - \lambda \chi^m - g \phi^n), \tag{2} \]

where \( S, \chi \) and \( \phi \) are chiral superfields, \( m \) and \( n \) are integers, and \( \mu, \lambda \) and \( g \) are taken to be real and positive by phase redefinition of the fields without loss of generality. We adopt the Planck units, in which \( M_p = 2.4 \times 10^{18} \text{ GeV} \) is set to be unity. The above form of the superpotential can be ensured by assigning R-charges as \( R(S) = 2 \), and \( R(\chi^m) = R(\phi^n) = 0 \) and appropriate discrete symmetry on \( \phi \) and \( \chi \). Later we will identify \( \chi^2 = H_uH_d \) or other flat directions.

The SUSY vacua is characterized by \( \chi \) and \( \phi \) satisfying

\[ \lambda \chi^m + g \phi^n = \mu^2. \tag{3} \]

The direction orthogonal to the moduli space has a large SUSY mass. In the low energy, we can integrate out those heavy degrees of freedom and focus on the dynamics of the moduli. There are two special symmetry-enhanced points in the moduli space, i.e., \( \chi = 0 \) and \( \phi = 0 \). As we will see shortly, the new inflation \[16] takes place as the field moves from one to the other.

Let us now introduce SUSY breaking which lifts the degeneracy of the moduli space. Assuming the gravity mediation for the scalar mass, we obtain

\[ V_{\text{soft}} = m_{\chi}^2 |\chi|^2 + m_{\phi}^2 |\phi|^2, \tag{4} \]

which is considered to be valid up to the Planck scale. The typical scale of \( m_{\chi} \) is given by the gravitino mass, \( m_{3/2} \), while \( m_{\phi} \) must be suppressed for successful inflation as we will
see later. We assume that the soft SUSY breaking scale is much lower than the SUSY mass scale in (2), namely,

\[
m_\chi \sim m_{3/2} \ll \text{Min} \left[ m\lambda \left( \frac{\mu^2}{\lambda} \right)^{\frac{n-1}{m}}, n g \left( \frac{\mu^2}{g} \right)^{\frac{n-1}{n}} \right].
\]

(5)

Then the scalar potential is obtained by substituting the SUSY condition (3) into \( V_{\text{soft}} \).

First let us study the scalar potential in the case of \( m_{\phi}^2 = 0 \). The moduli space is then lifted by the soft SUSY breaking mass of \( \chi \). The potential minimum and maximum are located at \( \chi = 0 \) and \( \phi = 0 \), respectively. Here and in what follows we focus on the real components of \( \phi \) and \( \chi \) and drop their imaginary components for simplicity.

The scalar potential around the maximum \( \phi = 0 \) can be expressed as

\[
V(\phi) \simeq m_\chi^2 \left( \frac{\mu^2}{\lambda} \right)^{\frac{1}{m}} \left( 1 - \frac{g}{\mu^2} \left( \frac{\phi}{\sqrt{2}} \right)^n \right)^{\frac{2}{m}}.
\]

(6)

where we have substituted (3) and defined \( \phi \equiv \sqrt{2} \text{Re}[\phi] \). The scalar potential around \( \phi = 0 \) is so flat that the inflation takes place. Note that the above expression (6) is not valid at \( \chi = 0 \), because \( \phi \) has a large SUSY mass there. The potential near \( \chi = 0 \) is simply given by the soft mass term of \( \chi \). See Fig. 1 for the typical shape of the inflaton potential. The inflation takes place near the top of the potential, and the inflaton \( \phi \) turns into \( \chi \) after inflation because of the SUSY condition given by (3). For sufficiently small \( |m_\phi^2| \) and \( m_{\phi}^2 < 0 \), the potential minimum and maximum are still located at \( \chi = 0 \) and \( \phi = 0 \) and the inflation is possible along the valley of the potential.

There are in general other contributions like

\[
\delta V = m_{3/2}^2 \lambda \chi^m + m_{3/2}^2 g \phi^n + \text{h.c.}
\]

(7)

which arise from the Kähler potential or inserting \( \langle S \rangle = O(m_{3/2}) \) in (2). However, these terms have negligible effects on the inflation dynamics as long as \( m_\chi \sim m_{3/2} \). (We will see that these terms are important for the AD mechanism later.)

Let us study the inflationary dynamics in detail. To simplify our analysis we focus on the inflaton dynamics around the top of the potential, \( \varphi \ll (\mu^2/g)^{1/n} \). The potential can

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1 This does not affect the inflation dynamics, but it may change the efficiency of the preheating, as will be discussed.

2 The \( S \) is known to develop a VEV of order \( m_{3/2} \) once the SUSY breaking is taken into account [17].
Figure 1: The shape of the inflaton potential. The right panel shows the contours of the inflaton potential in the left panel. The star and diamond denote the potential maximum and minimum in the moduli space, respectively. The inflation takes place along the valley of the potential near the maximum.

then be approximated as

\[ V(\varphi) \approx m_\chi^2 \left( \frac{\mu^2}{\lambda} \right)^\frac{1}{m} \left( 1 - \frac{2g}{m\mu^2} \left( \frac{\varphi}{\sqrt{2}} \right)^n \right) + \frac{1}{2} m_\phi^2 \varphi^2, \quad (8) \]

\[ \equiv V_0 - \kappa \varphi^n + \frac{1}{2} k V_0 \varphi^2. \quad (9) \]

For successful inflation, we require \( m_\phi^2 < 0 \) and \( |m_\phi^2| \ll V_0 \), or equivalently, \( k < 0 \) and \( |k| \ll 1 \). If this is satisfied, the inflation takes place at around \( \varphi = 0 \) and the inflaton \( \varphi \) turns into \( \chi \) after inflation and stabilized at \( \chi = 0 \). The fine-tuning of \( k \) is nothing but the so-called \( \eta \)-problem. We allow such fine-tuning if it is required for successful inflation. The inflation ends at \( \varphi = \varphi_{\text{end}} \) given by

\[ \varphi_{\text{end}} = \left( \frac{V_0 (1 + k)}{n(n-1)\kappa} \right)^{1 - n^{-2}}. \quad (10) \]

The position of the inflaton when the WMAP pivot scale exited the horizon is

\[ \varphi_N^{n-2} = \frac{k V_0}{n \kappa} \left[ 1 + \left( \frac{(n-1)k}{1+k} - 1 \right) e^{-N(n-2)k} \right]^{-1}, \quad (11) \]

where \( N \) denotes the e-folding number. The scalar spectral index is evaluated as

\[ n_s = 1 + 2k \left[ 1 - \frac{n-1}{1 + \left( \frac{(n-1)k}{1+k} - 1 \right) e^{-N(n-2)k}} \right]. \quad (12) \]
In the limit of $|k| \ll 1$, we obtain
\[ n_s \simeq 1 - \frac{2(n - 1)}{N(n - 2) + n - 1}. \] (13)

The spectral index for $n = 4, 5, 6$ and $7$ with $N = 50$ is shown in Fig. 2. We can see that the scalar spectral index is consistent with the WMAP result \cite{2}, $n_s = 0.968 \pm 0.012$, for $n \geq 5$, while the predicted value of $n_s$ for $n = 4$ has a slight tension with the WMAP result. This tension can be easily alleviated by the Coleman-Weinberg correction to the scalar potential if $\phi$ has a sizable Yukawa coupling \cite{18}.

The WMAP normalization is given by \cite{2}
\[ \frac{1}{12\pi^2} \frac{V(\phi)^3}{V'(\phi)^2} \simeq 2.43 \times 10^{-9}. \] (14)

For given parameters in the superpotential (2), the WMAP normalization condition \cite{14} fixes the soft SUSY breaking mass, $m_\chi$, and hence the Hubble scale during inflation, $H_{\text{inf}}$.

In Fig. 3 we show the contours of $\log_{10}[m_\chi/\text{GeV}]$ and $\log_{10}[H_{\text{inf}}/\text{GeV}]$. We can see that the soft SUSY breaking mass scale can be as low as about $10^7 \text{GeV}$. Here we have imposed a couple of constraints: the VEV of $\chi$ and $\phi$ should be smaller than $0.3M_p$; the mass orthogonal to the F-flat direction (3) is at least 3 times heavier than $m_\chi$. Similarly, we show the contours for the case of $m = n = 4$ in Fig. 4. The range of the soft mass
ranges from $10^7$ GeV to $10^9$ GeV for $\mu = 10^{12}$ GeV. If we take a large $g$, say, $g = (4\pi)^2$, the soft mass can be below $10^6$ GeV. For $n = 6$, the lower bound on the soft mass increases to $10^{11}$ GeV as one can see from Fig. 5. It is interesting that the range of the soft SUSY breaking mass agrees with those for which 125 GeV Higgs mass suggested by the recent ATLAS and CMS experiments [19, 20] can be realized [21].

Note that $H_{\text{inf}} < m_\chi (\sim m_{3/2})$ always holds in our model as long as $\chi_0 < M_P$. Therefore, for successful inflation, we need a tuning such that $|m_\phi| < H_{\text{inf}} < m_\chi$. As we can see from Figs. 3-5, the required amount of tuning is mild around the lower left region, and $m_\phi$ must be suppressed by about two orders of magnitude compared to its natural value.

So far we have not assumed the nature of $\chi$. An interesting possibility is that it consists of the D-flat direction of the SSM fields. As the simplest case, we can identify it with $\chi^2 = H_u H_d$. Then the inflaton transmutes into Higgs after inflation! This indicates that the efficient energy transfer to the visible sector takes place even if the $\phi$ is a gauge singlet and has no sizable interactions with SSM fields. The process of reheating will be discussed in the next section. It is also straightforward to identify $\chi^m$ with other flat directions such as $(\bar{u}d\bar{d})^2$ and $(LL\bar{e})^2$. In this case, the AD baryogenesis takes place naturally.

Figure 3: The contours of $\log_{10}[m_\chi/\text{GeV}]$ (left) and $\log_{10}[H_{\text{inf}}/\text{GeV}]$ (right) where we set $g = 1$, $m = 2$, $n = 4$ and $N = 50$. 

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3 Discussion

3.1 Reheating

An astonishing feature of the present inflation model is that the inflaton $\phi$ automatically transforms into another field $\chi$ after inflation. It is $\chi$ rather than $\phi$ that reheats the Universe. Therefore, even if $\phi$ is a singlet and its interaction to the SSM particles is suppressed, the efficient reheating can take place if $\chi$ has sizable interactions with the SSM particles, or $\chi$ itself can be the SSM field. Actually, $\chi$ can be identified with one of the D-flat directions in the SSM, like $\chi^2 = H_u H_d$ or $\chi^3 = L L \bar{e}$, etc. Here we show that the reheating is efficient in such a case because $\chi$ oscillates around the enhanced symmetry point, $\chi = 0$, and efficient particle production takes place in each oscillation. This should be contrasted to the conventional new inflation models, in which the inflaton oscillates around the large VEV and its perturbative decay rate can be suppressed.\(^3\)

Let us collectively denote the fields that couple to $\chi$ by $\psi$. The $\psi$ can be SSM (s)fermions and gauge bosons depending on the properties of $\chi$. Hereafter, we consider $\psi$

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\(^3\) The tachyonic preheating genetically takes place at the end of new inflation, however, the reheating is usually competed by the perturbative decay of the inflaton particle.
which has the strongest couplings to \( \chi \) and assume the couplings of the following form,

\[
\mathcal{L} = \begin{cases} 
  y^2|\chi|^2|\tilde{q}|^2 & \text{for a sfermion } (\psi = \tilde{q}) \\
  y\chi q\bar{q} + \text{h.c.} & \text{for a fermion } (\psi = q) \\
  y^2|\chi|^2A_\mu A^\mu & \text{for a gauge boson } (\psi = A_\mu) 
\end{cases}
\]

(15)

The mass of \( \psi \) is \( \chi \)-dependent, and is given by \( m_\psi \sim y \langle \chi \rangle \). Following discussion does not depend on whether \( \psi \) is a boson or fermion. For example, if \( \chi \) consists of the \( H_uH_d \), \( \psi \) represents top (s)quarks or weak gauge bosons, and the coupling \( y \) denotes either top Yukawa coupling or SU(2) gauge coupling. The exact form of the interactions does not matter in the following discussion.

After inflation ends, \( \chi \) oscillates around the minimum \( \chi = 0 \) as \( \chi(t) = \chi_0 \sin(m_\chi t) \) (here \( t = 0 \) is set to be the instant when \( \chi \) reaches to minimum \( \chi = 0 \)). We neglect the cosmic expansion because \( m_\chi \gg H_{\text{inf}} \) and because we are only interested in the dynamics within the time scale of dozens of \( \chi \) oscillations. The \( \psi \) field obtains a large mass through the coupling with \( \chi \) and the perturbative decay of \( \chi \) into \( \psi \) is prohibited during most of the time of oscillations except for the vicinity of \( \chi \sim 0 \), where the adiabaticity of the \( \psi \) is violated [22]. The frequency of the \( \psi \)-particle with a wavenumber \( k \) is given by \( \omega_k = \sqrt{y^2\chi^2 + k^2} \). (This \( k \) should not be confused with the one in the inflaton potential [9].) As \( \chi \) passes the origin, the adiabaticity is violated, i.e., \( |\dot{\omega}_k/\omega_k^2| > 1 \), and the \( \psi \)-particle is produced. The typical wavenumber \( k_* \) of the produced \( \psi \)-particle is estimated.
as \( k_* = \sqrt{y \chi_0 m_\chi} \). Therefore, each time the \( \chi \) passes through the origin, the number density of \( \psi \) increases by

\[
n_\psi \simeq \frac{k_*^3}{8\pi^3} \approx \frac{(y \chi_0 m_\chi)^{3/2}}{8\pi^3}.
\]

The produced \( \psi \) particle becomes massive as \( \chi \) goes far away from the origin, and then the \( \psi \) decays into light states. The decay rate of \( \psi \) is estimated as \( \Gamma_\psi \sim (f^2/8\pi)m_\psi \), where \( f \) is either a gauge or Yukawa coupling of \( \psi \), and \( m_\psi \approx y \chi \). Thus \( \psi \) decays at \( t = t_{\text{dec}} \sim \sqrt{8\pi/(f^2y \chi_0 m_\chi)} \) at which it has a mass of \( m_\psi \sim \sqrt{8\pi y \chi_0 m_\chi}/f \).

This happens much before the \( \chi \) returns again back to \( \chi = 0 \). Therefore, in each oscillation, the energy of \( \chi \) decreases by

\[
\frac{\Delta \rho_\chi}{\rho_\chi} \approx \frac{m_\psi(t_{\text{dec}})n_\psi}{m_\psi^2 \chi_0^2/2} \sim \frac{y^2}{\sqrt{2\pi^5/2}f}.
\]

This is \( O(0.1 - 0.01) \) depending on the couplings of \( \chi \) and \( \psi \). Therefore, the \( \chi \) energy density is efficiently transferred to the radiation within dozens of oscillations through the \( \psi \) particle production and its subsequent decay. The whole process is what is called the instant preheating [23]. (See also Ref. [24].) If the reheating is completed by the instant preheating, the reheating temperature is estimated as \( T_R \sim \sqrt{H_{\text{inf}}} \), and it is given by \( T_R \sim 10^{12} \text{GeV} \) for \( H_{\text{inf}} \sim 10^6 \text{GeV} \). Thermal leptogenesis works for such a high reheating temperature [25]. Note that, as the radiation energy increases through the instant preheating, the \( \psi \)-particle acquires a thermal mass, which can make the energy transfer inefficient and therefore can terminate the preheating process when the energy density of \( \chi \) becomes comparable to that of radiation. If such back reaction is relevant, the reheating is completed by the usual perturbative decay. In this case, the reheating temperature is considered to be about the soft SUSY breaking mass, \( T_R \sim m_\chi \).

So far we have assumed that \( \chi \) passes the origin. This may not be the case if \( \chi \) acquires an angular momentum in its complex plane. In fact, if we take account of the imaginary component of \( \phi \), it acquires a certain angular momentum during and after inflation, which is transferred to \( \chi \) through the SUSY condition (3). In this case, the preheating process may become inefficient and the perturbative decay will complete the reheating.

\[4\] Note that \( t_{\text{dec}} > 1/k_* \) and hence the implicit assumption that the \( \psi \) decays after the particle production ends is justified.
Note that SUSY particles can be produced both thermally and non-perturbatively. Since the SUSY breaking scale, which is characterized by $m_{3/2}$, is heavier than PeV in this model, the abundance of the lightest SUSY particle (LSP) likely exceeds the dark matter abundance if the R-parity is conserved. The LSP overabundance can be avoided if the R-parity is violated by a small amount so that the LSP decays before the big-bang nucleosynthesis (BBN). If the reheating temperature is higher than $m_{3/2}$, gravitinos are also efficiently produced by thermal scatterings, and they decay into LSPs much before the BBN and those LSPs quickly decay through the R-parity violating interactions. Alternatively, the cosmic density of the LSP can be suppressed if the LSP mass is sufficiently light as in the case of the axino LSP.

### 3.2 Initial condition

In order for our inflation scenario to work, the initial position of the inflaton must be near the maximum $\phi = 0$ along the flat direction $\chi$. This initial condition is dynamically realized by considering the pre-inflation which occurs in the same superpotential $S$. First note that the scalar potential is flat along $S$ for $\phi = \chi = 0$ where the potential energy is given by $\mu^4$. Thus the pre-inflation takes place for $\phi \sim \chi \sim 0$ with a sufficiently large value of $S$, and $S$ plays the role of inflaton. We assume $\phi$ is stabilized at the origin by a positive Hubble-induced mass during the pre-inflation. If $m = 2$, $\chi$ is stabilized at the origin until $\chi$ becomes tachyonic when $S \sim \mu/\sqrt{\lambda}$ and the pre-inflation ends. The inflationary dynamics is same as the hybrid inflation [11]. If $m > 2$, $\chi$ develops a small but non-zero $S$-dependent VEV, and the inflation ends when $S$ becomes of order $(\mu^2/\lambda)^{1/(2m-2)}$ as in the smooth hybrid inflation model [12]. In both cases, it is $\chi$ that develops large VEV at the end of pre-inflation while $\phi = 0$. Thus the desired initial condition for the new inflation is dynamically achieved.

### 3.3 Extensions to other D-flat directions, and AD baryogenesis

As already noticed, the $\chi$ field can be identified with other D-flat directions in the SSM, and if it has non-zero baryon and/or lepton numbers, it can create the baryon asymmetry

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5 Even if $\phi$ is not stabilized at the origin during pre-inflation, it may settle down at the origin after the pre-inflation through its interactions with the decay product of $S$ or $\chi$. 

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through the AD mechanism.

As we have mentioned in Footnote 1, $\chi$ can acquire a non-zero angular momentum from the inflaton dynamics. This however depends on the initial condition of the imaginary component of the inflaton $\phi$. Let us here concentrate on the dynamics of $\chi$ after inflation. For example, we consider $\chi^6 = (LL\bar{e})^2$. The AD field obtains a scalar potential of the form given by (17), which serves as the source of the baryon number violation and induces an angular motion on the complex plane of $\chi$. Although the AD field receives many angular kicks because of $m_\chi \gg H_{\text{inf}}$, the first one gives the dominant contribution. The resultant baryon asymmetry is estimated as

$$\frac{n_B}{s} \sim \frac{T_R}{H_{\text{inf}}^2} \frac{\lambda m_\chi^{2/2} \chi_0^m}{H_{\text{inf}}} \delta$$

where we have defined $\chi_0 \equiv (\mu^2/\lambda)^{1/m}$, and $\delta(<1)$ represents the CP phase. The particle production becomes inefficient if the $\chi$ has large angular momentum, because the $\chi$ does not pass the origin. If the non-perturbative decay of $\chi$ is suppressed, the reheating temperature is considered to be of order $m_\chi$. Therefore $T_R$ should be in the range between $m_\chi$ and $\sqrt{H_{\text{inf}}}$. We have found that in the case of $m = 6$, $n = 4$ and $\delta \sim O(0.1)$, the above baryon asymmetry varies from $O(10^{-9})$ to $O(10^{-3})$ in the allowed region of $\mu$ and $\lambda$. Therefore, it is possible to generate a right amount of the baryon asymmetry via the AD mechanism. The magnitude of the baryon isocurvature perturbation is given by $S_b \sim H_{\text{inf}}/(2\pi \chi_0)$ and it is much smaller than the observational constraint. We have confirmed it is smaller than $10^{-8}$ for the above case.

### 3.4 Other applications

Let us comment on other applications of our model. One extension of the model is consider an exponential dependence of $\phi$.

$$W = S(\mu^2 - \lambda \phi^m - g e^{-b\phi}).$$

The Kähler potential is considered to respect a shift symmetry of $\phi$: $\phi \rightarrow \phi + i\alpha$, where $\alpha$ is a real transformation parameter. The inflationary dynamics is similar to the case studied in Sec. 2 and some results can be obtained simply by taking a limit of $n \rightarrow \infty$. 
(For instance, $n_s$ can be obtained by taking this limit in Eq. (13).) Interestingly, if we take $m = 2$ and $\chi^2 = H_u H_d$, the resultant potential is quite similar to the Higgs inflation \cite{27}. The great difference is that, although the potential is similar, there is no large coupling in our model, avoiding the issue of the unitarity.

It is also possible to use the moduli space \cite{31} as a curvaton \cite{26}. In particular, the so-called hilltop curvaton can be easily realized because the potential maximum is a symmetry-enhanced point \cite{28, 29}.

So far we have focused on the case of $n \geq 3$. In the case of $n = 2$, the sufficient amount of inflation does not occur, but thermal inflation may take place, if $\phi$ acquires a thermal mass through its interactions with the ambient plasma.\footnote{Similar arguments have been made in \cite{30} in the context of anomaly-mediation with U(1) extension to solve the tachyonic slepton problem.} The reheating is induced by the decay of $\chi$ in this case, and the resultant reheating temperature will be so high that the leptogenesis will be possible. The reheating process is similar to that described before.

\section{Conclusions}

In this letter, we have proposed a new inflation model, in which the SM gauge singlet inflaton turns into the Higgs field (or more precisely, $H_u H_d$) after inflation. The energy stored in the inflaton is directly transferred to the $H_u H_d$ condensate. This is a novel type of reheating; instead of producing the SM particles from the inflaton decay, the inflaton energy is directly transmuted into the $H_u H_d$ flat direction. The transmutation is due to the inflaton dynamics in the moduli space of supersymmetric vacua spanned by the inflaton and the Higgs field. Interestingly, because of this transmutation, the new inflation ends up in a symmetry vacuum of the Higgs boson, which enables rapid thermalization of the Higgs condensate. The initial condition for the successful inflation can be naturally realized by the pre-inflation in which the Higgs field plays the role of a waterfall field. It is possible to replace the role of $H_u H_d$ with other D-flat directions in SSM. In particular, we have also pointed out that the AD baryogenesis naturally takes place in this case.
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