Low-scale Supersymmetry from Inflation

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Abstract

We investigate an inflation model with the inflaton being identified with a Higgs boson responsible for the breaking of U(1)$_{B-L}$ symmetry. We show that supersymmetry must remain a good symmetry at scales one order of magnitude below the inflation scale, in order for the inflation model to solve the horizon and flatness problems, as well as to account for the observed density perturbation. The upper bound on the soft supersymmetry breaking mass lies between 1 TeV and 10$^3$ TeV. Interestingly, our finding opens up a possibility that universes with the low-scale supersymmetry are realized by the inflationary selection. Our inflation model has rich implications; non-thermal leptogenesis naturally works, and the gravitino and moduli problems as well as the moduli destabilization problem can be solved or ameliorated; the standard-model higgs boson receives a sizable radiative correction if the supersymmetry breaking takes a value on the high side $\sim 10^3$ TeV.
1 Introduction

The inflationary paradigm [1] has been well established so far. A number of theoretical difficulties of the standard big bang cosmology are naturally circumvented by the exponential expansion of the universe during inflation, and more important, the quantum fluctuation of the inflaton field can account for the observed density perturbation.

Despite the success of the inflationary paradigm, it has been considered extremely challenging to answer the question, what is the inflaton. If the inflaton is just a gauge singlet with extremely weak interactions with the standard-model particles, it would be almost impossible to identify the inflaton in a laboratory experiment. One way to avoid this conclusion is to build a successful inflation model in the framework of the standard model (SM) or its extensions. In the SM, the Higgs boson $\phi_{SM}$ is the only scalar field, and therefore a candidate for the inflaton. It is indeed possible to build an inflation model using $\phi_{SM}$, relying on a non-canonical kinetic term [2, 3] and/or a non-minimal coupling to gravity [4, 5].

Since the discovery of neutrino oscillations, the right-handed neutrinos, $\nu_R$, are usually incorporated in the minimal extension of the SM to explain the small, but non-vanishing neutrino masses. In particular, the extremely light neutrino mass scale can be naturally accounted for by the see-saw mechanism [6], which requires the heaviest right-handed neutrino at a scale of $10^{15}$ GeV close to the GUT scale. With the addition of the three right-handed neutrinos, it is then reasonable to introduce the $U(1)_{B-L}$ gauge symmetry which is required by the charge quantization condition and is also motivated by the GUT gauge group such as SO(10). Thus, we consider the framework, SM+$\nu_R$+$U(1)_{B-L}$, as the minimal extension of the SM. In this theoretical framework, we have another candidate for the inflaton, namely the Higgs boson, $\phi_{B-L}$, which is responsible for the breaking of the $U(1)_{B-L}$ symmetry. In this paper we explore a possibility that the Higgs boson $\phi_{B-L}$ plays a role of the inflaton and discuss its implications.

The SM has been successful in explaining numerous experimental data with a great accuracy, and there is no hard evidence for physics beyond the SM (with neutrino masses included). On the other hand, it has been known that there is a gauge hierarchy problem in the SM, which was the motivation to consider the physics beyond SM such as super-
symmetry (SUSY). However, after the LEP experiment, the supersymmetric extension of SM (SSM) turned out to be not free of fine-tunings. Indeed, typically a fine-tuning at the percent level is required for the correct electroweak breaking, which casts doubt on the conventional naturalness argument as the correct guiding principle for understanding the physics at and beyond the weak scale.

On the other hand, it is now widely accepted that SUSY should appear at a certain energy scale, which may be much higher than the weak scale, because the string theory, the most qualified candidate for the unified theory including gravity, requires supersymmetry for theoretical consistency, and it may remain in the effective 4D theory below the compactification scale. Further, the recent observation of a cosmological constant within the anthropic window strongly suggests the presence of the string landscape. Motivated by these considerations, we do not rely on the conventional naturalness argument for building an inflation model. For instance we do not care much about the fine-tuning needed to make the inflaton potential flat, because such a tuning may be easily compensated by the subsequent exponential expansion of the universe during inflation, and because clearly we cannot live in the universe which does not experience inflation. Instead, we take the existence of the inflationary phase (driven by $\phi_{B-L}$ in the model considered below) as a guiding principle. Also we assume the presence of SUSY, but we leave the SUSY breaking scale as a free parameter since it may be subject to the distribution of vacua in the landscape or anthropic selection. Indeed as we will see, the SUSY should remain a good symmetry at scales one order of magnitude below the inflation scale for the inflation model to be successful. Typically the soft SUSY breaking masses for the SSM particles lie between 1 TeV and $10^3$ TeV, whose precise value depends on the $B-L$ breaking scale and the inflaton potential. If there is a bias toward high-scale SUSY breaking in the landscape, the SUSY breaking masses may be close to $10^3$ TeV. It is intriguing that the low-scale SUSY emerges as a result of the inflationary selection, irrespective of the gauge hierarchy problem.

Before closing the introduction, let us here briefly mention the inflation scenario using

\[1\] Here and in what follows, the low-scale SUSY means that SUSY remains a good symmetry at scales much smaller than the Planck or string scale.
the GUT Higgs boson, since it is an old topic and was studied extensively in the past. The graceful exit problem of the original inflation relying on the first-order phase transition was avoided in the new inflationary universe scenario (new inflation) proposed by Linde [8]. The phase transition in the new inflation was of Coleman-Weinberg (CW) type [9], where the inflaton was the GUT Higgs boson with the mass at the origin being set to be zero. Although this scenario was very attractive, it was soon realized that the CW correction arising from the gauge boson loop makes the inflaton potential too steep to produce the density perturbation of the correct magnitude, $\delta \rho / \rho \approx 10^{-5}$ [11]. In fact, the required magnitude of the gauge coupling constant was many orders of magnitude smaller than the expected value of the unified gauge coupling constant, which clearly implied that some modification was needed. One solution was to consider a gauge singlet inflaton, which has extremely weak interactions with the SM particles. Although the inflation model may lose its connection to the GUT in this case [10], such gauge singlets are ubiquitous in the string theory, and so, one of them may be responsible for the inflation. There have been many works along this line [12]. Another way to resolve the problem is to introduce SUSY. Then the CW potential becomes suppressed because of the cancellation among bosonic and fermionic degrees of freedom running the loop [14]. In this paper we will explore the latter possibility in detail and estimate the required size of the SUSY breaking.

The rest of the paper is organized as follows. In Sec. 2 we discuss the inflationary dynamics of the $U(1)_{B-L}$ Higgs boson considering only the tree-level contributions, anticipating that the CW potential will be partially canceled by SUSY. We will consider the radiative corrections to the inflaton potential and derive the upper bound on the SUSY breaking scale in Sec. 3. We discuss implications of our scenario in Sec. 4. The last section is devoted for discussion and conclusions.

2 Set-up and inflaton dynamics

In the following we use $\phi$ to denote the Higgs boson responsible for the $U(1)_{B-L}$ breaking. Here we do not assume SUSY to allow a situation in which the inflation scale is lower.

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2 We note that the SUSY GUT provides a natural framework for hybrid inflation [12].
than the SUSY breaking scale. We will discuss the supersymmetric version in the next section.

Let us consider an inflaton potential given by
\[ V = V_0 - m_0^2 |\phi|^2 - \lambda_n \frac{|\phi|^{2n}}{M_{*}^{2n-4}} + \lambda_m \frac{|\phi|^{2m}}{M_{*}^{2m-4}}, \]
(1)
where \( m \) and \( n \) are integers satisfying \( m > n \geq 2 \) and \( M_\ast \) is a cut-off scale of the theory. We expect \( M_\ast \) to be not far from the GUT scale \( \sim 10^{15} \) GeV. The CW potential arising from the B–L gauge boson and the right-handed neutrino loops will be considered in the next section. For the moment we focus on the inflationary dynamics using the tree-level contributions.

After the \( \phi \) breaks the U(1)\(_{B-L}\) gauge symmetry, its phase component is absorbed into the massive B–L gauge boson. So we focus on its radial component:
\[ \phi \equiv \sqrt{2} \phi. \]
(2)
In terms of \( \phi \), we can write the scalar potential as
\[ V(\phi) = V_0 - \frac{m_0^2}{2} \phi^2 - \frac{\kappa}{2n} \frac{\phi^{2n}}{M_{\ast}^{2n-4}} + \frac{\lambda}{2m} \frac{\phi^{2m}}{M_{\ast}^{2m-4}}, \]
(3)
with
\[ \kappa \equiv \frac{n \lambda_n}{2n-1}, \]
(4)
\[ \lambda \equiv \frac{m \lambda_m}{2m-1}. \]
(5)
This potential has a global minimum at \( \phi = \varphi_{\text{min}} \) given by
\[ \varphi_{\text{min}} = \left( \frac{\kappa}{\lambda} \right)^{\frac{1}{(m-n)}} M_{\ast}, \]
(6)
which gives the U(1)\(_{B-L}\) symmetry breaking scale at low energy. For the above effective theory description to be valid, \( \varphi_{\text{min}} \lesssim M_\ast \) must be satisfied. Requiring that the cosmological constant vanishes at the minimum, we obtain
\[ V_0 = \left( \frac{m-n}{2mn} \right) \left( \frac{\kappa^m}{\lambda^n} \right)^{\frac{1}{m-n}} M_{\ast}^4. \]
(7)
Here we have assumed that the mass term is negligibly small compared to the higher order terms at the potential minimum.

The B−L breaking scale can be inferred from the neutrino oscillation data as follows. The Majorana mass $m_N$ for the right-handed neutrino $\nu_R$ is related to the B−L breaking scale through the following interaction,

$$L = -\frac{y_N}{2} \phi \bar{\nu}_R \nu_R + \text{h.c.}.$$  

and we obtain $m_N = y_N \varphi_{\text{min}} / \sqrt{2}$. The coupling constant $y_N$ is expected to be order unity for the heaviest $\nu_R$. Then the B−L breaking scale $\varphi_{\text{min}}$ is estimated to be about $10^{15}$ GeV, close to the GUT scale, using the seesaw formula [6].

For simplicity we drop the mass term, setting $m_0 = 0$, in the following analysis. All the results remain almost intact as long as the mass is much smaller than the Hubble parameter during inflation. Also, such a small mass may be favored since the total e-folding number of the inflation will be longer.

The inflation takes place if the initial position of $\varphi$ is sufficiently close to the origin. Since the inflaton has couplings to the B−L gauge boson, the right-handed neutrinos and the SM particles (through the U(1)B−L gauge interaction), we expect that the initial position of $\varphi$ before the inflation starts is naturally close to the origin, assuming the presence of thermal plasma in the universe. The inflation ends at the point where the slow-roll conditions are violated, namely, one of the slow-roll parameters $\eta$ becomes order unity,

$$\eta = M_p^2 \frac{V''}{V} \simeq -1,$$  

where $M_p \simeq 2.4 \times 10^{18}$ GeV is the reduced Planck mass. This occurs at

$$\varphi_{\text{end}} = \left[ \frac{m-n}{2mn} \frac{1}{2n-1} \left( \frac{\kappa}{\lambda} \right)^{\frac{n}{n-1}} M_s^{2n} M_p^{-2} \right]^{\frac{1}{2(n-1)}},$$

$$= \left[ \frac{m-n}{2mn} \frac{1}{2n-1} \right]^{\frac{1}{2(n-1)}} \varphi_{\text{min}}^{\frac{n}{n-1}} M_p^{-\frac{1}{n-1}}.$$  

\footnote{We assume that the universe had experienced another inflation before the last inflation by $\phi$ started. The radiation may have come from the decay of the inflaton responsible for the preceding inflation. Even without the radiation, the initial position may be naturally set to be at the origin by e.g. the Hubble-induced mass since it is the enhanced symmetry point.}
The field value of $\phi$ at $N$ e-foldings before the end of inflation, $\phi_N$, can be estimated as follows. The e-folding number $N$ is given by

$$N \simeq \int_{\phi_N}^{\phi_{\text{end}}} \frac{3H^2}{V'} d\phi \simeq \frac{V_0 M_s^{2n-4}}{2(n-1)M_p^2\kappa} \phi_N^{2-2n},$$

therefore we obtain

$$\phi_N \simeq \left(\frac{1}{N}\right)^{\frac{1}{2(n-1)}} \left(\frac{2n-1}{2(n-1)}\right)^{\frac{1}{2(n-1)}} \phi_{\text{end}},$$

where $\phi_N \ll \phi_{\text{end}}$ is assumed. The slow-roll parameter $\eta$ at $\phi = \phi_N$ is given by

$$\eta \simeq -\frac{2n-1}{2(n-1)N}.$$

The other slow-roll parameter, $\epsilon$, is much smaller than $|\eta|$, which is typically the case in the new inflation model. Thus the scalar spectral index, $n_s$, is then given by

$$1 - n_s = \frac{2n-1}{n-1} \frac{1}{N}.$$  

In the limit $n \gg 1$, it approaches to $n_s = 0.96$ for $N = 50$, which is close to the center value of the WMAP result [15]. For $n = 2(3)$, the spectral index is about 0.94(0.95), which is also consistent with observation. As we will see below, $n \geq 3$ is needed for non-thermal leptogenesis to work unless there is a degeneracy among the right-handed neutrino masses. Thus, the spectral index is predicted to be between 0.94 and 0.96 in the simple case where the inflaton potential is dominated by a single monomial term during the last 50 e-foldings. The precise value of $n_s$ actually depends on the details of the inflaton potential, and it is possible to slightly modify the prediction.

In order to account for the density perturbation by the quantum fluctuation of the inflaton, we impose the WMAP normalization condition [15],

$$\Delta_R^2 \simeq 2.42 \times 10^{-9},$$

where $\Delta_R$ denotes the power spectrum of the curvature perturbation $R$. In terms of the inflaton potential, it is given by

$$\frac{V_0^3}{M_p^6 V'^2} \simeq 2.9 \times 10^{-7}.$$
Using Eq. (7), we obtain
\[
\kappa \simeq 2.9 \times 10^{-7} \left( \frac{1}{2(n-1)N} \right)^{\frac{2n-2}{n-1}} \left( \frac{2mn}{m-n} \right)^{\frac{2n-2}{n-1}} \left( \frac{\varphi_{\text{min}}}{M_*} \right) - 2 \left( \frac{\varphi_{\text{min}}}{M_*} \right)^{\frac{2(n-2)}{n-1}}. \quad (17)
\]

In addition to \( n \) and \( m \), the inflaton potential of our interest has four parameters: \( V_0, \kappa, \lambda, \) and \( M_* \). \( V_0 \) is determined by requiring the vanishing cosmological constant at the true vacuum, see (7). Instead of \( \lambda \), we prefer to use the physically relevant quantity \( \varphi_{\text{min}} \), which is the B–L breaking scale, given by (6). The WMAP normalization condition (17) then fixes the value of \( \kappa \). As a result, we can parametrize the potential by two parameters: the cutoff scale \( M_* \) and the B–L breaking scale \( \varphi_{\text{min}} \).

Since the inflaton potential is dominated by \( V_0 \) during inflation, the Hubble parameter during inflation is estimated as
\[
H_{\text{inf}} = 3.1 \times 10^{-4} \left( \frac{m-n}{2mn} \right)^{\frac{1}{2(n-1)}} \left( \frac{1}{2(n-1)N} \right)^{\frac{1}{2(n-1)}} \left( \frac{\varphi_{\text{min}}}{M_p} \right)^{\frac{1}{2(n-1)}}. \quad (18)
\]

Note that the inflation scale is solely determined by the B–L breaking scale, independent of \( M_* \). The inflation scale is shown in Fig. 1 for \( \varphi_{\text{min}} = 10^{14} \text{GeV} \) (left) and \( 10^{15} \text{GeV} \) (right). One can see that, in the case of \( \varphi_{\text{min}} = 10^{14}(10^{15}) \text{GeV} \), the Hubble parameter lies in the range of \( 10^2(10^3) \text{GeV} \) and \( 10^7(10^8) \text{GeV} \). In the case of \( n = 2(3) \), the Hubble parameter lies in the range of \( 10^2(10^5) \text{GeV} \sim 10^4(10^6) \text{GeV} \).

The values of \( \kappa \) and \( \lambda \) are also plotted in Fig. 2 with \( \varphi_{\text{min}} = 10^{14} \text{GeV} \) and \( M_* = 10^{15} \text{GeV} \) (top left), \( \varphi_{\text{min}} = 10^{15} \text{GeV} \) and \( M_* = 10^{15} \text{GeV} \) (top right), and \( \varphi_{\text{min}} = 10^{15} \text{GeV} \) and \( M_* = 10^{16} \text{GeV} \) (bottom). In the case of \( \varphi_{\text{min}} < M_* \), the values of \( n \) and \( m \) are bounded above to avoid too large numerical coefficients. In the case of the top left and bottom panels, we obtain \( n \leq 4 \). However, this is sensitive to the relative magnitude of \( \varphi_{\text{min}} \) and \( M_* \). Indeed, for \( \varphi_{\text{min}} \approx M_* \), there is no such upper bound, and both \( \kappa \) and \( \lambda \) asymptote to \( \sim 10^{-5} \) as \( n \) becomes large (see the top right panel in Fig. 2).

Thus, if we allow fine-tuning of the parameters, \( m_0, \kappa \) and \( \lambda \), the successful inflation takes place using the B–L Higgs boson. The inflation scale is typically very low, and only negligible amount of the tensor mode is generated. Note however that we have only considered the tree-level potential. As we will see in the next section, the radiative corrections generically spoil the inflationary dynamics.
Figure 1: The Hubble parameter during inflation $H_{\text{inf}}$ with respect to $n$. Here we set $m = n + 1$. The left is for $\varphi_{\text{min}} = 10^{14}$ GeV and the right is for $10^{15}$ GeV. The relation between $H_{\text{inf}}$ and $n$ is almost same for $m = 2n$.

3 Radiative correction to the inflaton potential

Now let us turn to the issue of the radiative correction to the inflaton potential. The inflaton, the B–L Higgs boson, necessarily couples to the $\text{U}(1)_{\text{B–L}}$ gauge boson. Furthermore, it is expected to be coupled to the right-handed neutrinos to generate large Majorana masses. Due to these interactions, the inflaton potential receives corrections at the one-loop level. The general form of the CW effective potential is given by [9]

$$V_{\text{CW}} = \frac{1}{64\pi^2} \left[ \sum_{B = \text{boson}} m_B^4 \left( \frac{m_B^2}{\mu^2} - \frac{3}{2} \right) - \sum_{F = \text{fermion}} 2m_F^4 \left( \frac{m_F^2}{\mu^2} - \frac{3}{2} \right) \right], \quad (19)$$

where the sum over bosons counts a real scalar, and that over fermions counts a Weyl fermion. Here the subscript $B$ denotes bosons, and it includes the $\text{U}(1)_{\text{B–L}}$ gauge boson, while $F$ denotes fermions including the right-handed neutrinos. Since the masses of the $\text{U}(1)_{\text{B–L}}$ gauge boson as well as the right-handed neutrinos depend on the inflaton field $\varphi$, the inflaton potential receives the CW correction. For the moment we drop the contribution from the right-handed neutrinos and focus on the gauge boson contribution.

In fact, it is well known that the CW potential arising from the gauge boson loop makes the effective potential so steep that the resultant density perturbation becomes much larger than the observed one [11]. One way to solve the problem is to consider a gauge singlet inflaton. Here we stick to the inflation model using the B–L Higgs boson and
Figure 2: The $\kappa$ and $\lambda$ with respect to $n$. We set $m = n + 1$. The top left is for $\varphi_{\text{min}} = 10^{14}$ GeV and $M_* = 10^{15}$ GeV, the top right is for $\varphi_{\text{min}} = 10^{15}$ GeV and $M_* = 10^{15}$ GeV, and the bottom is for $\varphi_{\text{min}} = 10^{15}$ GeV and $M_* = 10^{16}$ GeV. $n$ is bounded above as $n \leq 4$ for the top left and bottom cases. If $m = 2n$, the bound becomes $n \leq 3$ for these cases.
explore other possibilities. Then we need to suppress the radiative correction somehow.\footnote{It is not possible to cancel the CW potential by tuning the tree-level potential, because of the logarithmic factor. Here we assume that the effective theory below $M_*$ is regular and that there are no additional light degrees of freedom other than SM+$\nu_R$+U(1)$_{B-L}$(+SUSY).}

One way to cancel or suppress the CW potential is to introduce SUSY. In the exact SUSY limit, contributions from boson loops and fermion loops are exactly canceled out. However, if SUSY is broken, the non-vanishing CW corrections remain.

In SUSY, two U(1)$_{B-L}$ Higgs bosons are required for anomaly cancellation. Let us denote the corresponding superfields as $\Phi(+2)$ and $\bar{\Phi}(-2)$ where the number in the parenthesis denotes their B–L charge. The $D$-term potential vanishes along the $D$-flat direction $\Phi\bar{\Phi}$, which is to be identified with the inflaton. Actually, a linear combination of the lowest components of $\Phi$ and $\bar{\Phi}$ corresponds to $\phi$.

The gauge boson as well as the scalar perpendicular to the $D$-flat direction have mass of $m_B^2 = g^2\varphi^2$. On the other hand, there are additional fermionic degrees of freedom, U(1)$_{B-L}$ gaugino and the B–L higgsino, whose mass eigenvalues are given by $m_F = g\varphi \pm \tilde{m}$, where $\tilde{m}$ denotes the SUSY breaking mass for the B–L gaugino. Because of the SUSY breaking mass $\tilde{m}$, the CW potential does not vanish and the inflaton receives a non-zero correction to its potential.\footnote{The possibility that the gauge non-singlet inflaton is protected from radiative corrections by SUSY was pointed out in Ref. [14], but the estimate of the CW potential during inflation is not correct and is different from ours by many orders of magnitude.} Inserting the field dependent masses into the CW potential (19), and expanding it by $\tilde{m}/(g\varphi)$, we find

$$V_{CW}(\varphi) \simeq \frac{g^2}{8\pi^2} \left(1 - 3 \ln \frac{g^2\varphi^2}{\mu^2}\right) \tilde{m}^2\varphi^2.$$  \hspace{1cm} (20)

Thus, in the presence of SUSY, the CW potential becomes partially canceled and the dependence of the inflaton field has changed from quartic to quadratic as long as $\tilde{m} \ll g\varphi$. Note that the correction still contains a logarithmic factor, which is not negligible if we consider the whole evolution of the inflaton from the origin.

If the mass of the CW potential exceeds the Hubble parameter, the inflaton does not slow-roll and inflation does not occur. In order not to disturb the inflationary dynamics studied in the previous section, therefore, the mass correction due to the CW potential should be sufficiently small for a certain range of the inflaton field. We require that the curvature of the inflaton potential is much smaller than the Hubble parameter everywhere
from $\varphi = H_{\text{inf}}/2\pi$ to $\varphi = \varphi_N$. This is a reasonable assumption because, in the new inflation scenario, the universe likely experiences the eternal inflation when the inflaton is near the origin where the quantum fluctuation dominates the dynamics. The presence of the eternal inflation may be favored in the landscape, since it can compensate the required fine-tuning of the parameters. For $n = 2 - 10$ with $N = 50$, we estimate the logarithmic factor as $\ln(\varphi_N^2/(H_{\text{inf}}/2\pi)^2) \sim 30$. Therefore, the following inequality must be satisfied,

$$ \frac{g^2}{8\pi^2} 90 \tilde{m}^2 < 0.01 H_{\text{inf}}^2, \quad (21) $$

therefore,

$$ \tilde{m} < 0.1 H_{\text{inf}}. \quad (22) $$

The reason why we put 0.01 in the right-handed side of Eq. (21) is that the observed spectral index, $1 - n_s$, is of $\mathcal{O}(0.01)$.

Thus, the SUSY breaking scale should be one order of magnitude smaller than the Hubble parameter during inflation, if one requires that the successful inflation take place. This is the main result of this paper.

In the gravity mediation, the soft SUSY breaking mass for $B-L$ gaugino is considered to be comparable to the soft SUSY masses for the SSM particles. For simplicity we assume the gravity mediation in the following. We will come back to this issue and consider other possibilities in Sec. 4.

We emphasize here that this novel bound on the soft SUSY breaking mass is derived from the requirement that the inflation should occur. Even if high-scale SUSY breaking scale is favored in the string landscape, the anthropic pressure by the inflation constrains the SUSY breaking scale below the inflation scale. Also, in this case we have a prediction that the SUSY breaking scale should be close to the inflation scale. For the choice of $n = 2$ and $\phi_{\min} = 10^{15}\text{GeV}$, the inflation scale is given by $H_{\text{inf}} = \mathcal{O}(10^4)\text{GeV}$. Thus, assuming that the soft SUSY mass is close to the upper bound, we obtain $\tilde{m} = \mathcal{O}(1)\text{TeV}$. That is to say, the SUSY breaking scale, $\tilde{m}$, happens to be close to the weak scale, independently of the gauge hierarchy problem. If this is the case, SUSY may be discovered in the TeV range at the LHC. As $n$ becomes larger, the SUSY breaking scale can be higher, but it is generically smaller than $\mathcal{O}(10^3)\text{TeV}$. The upper bound may be saturated in the landscape.

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6 Here we are interested in the inflation models which explain the observed data in our universe, since it is hard to estimate the likelihood of universes with observables taking different values.
if there is bias toward the longer duration of the inflationary phase. In this case, the SUSY particles are unlikely to be discovered at the LHC unless the SSM mass spectrum has some hierarchical structure, but we may be able to see the hint for the SUSY breaking scale of $O(10^3)$ TeV from the large radiative correction to the SM Higgs boson mass, which may fall in the range of $m_H \gtrsim 140$ GeV or so [37].

Since the SUSY remains a good symmetry below the inflation scale, it is possible to write down the inflation model in SUSY. Since the $\phi$ cannot have a large $F$-term when it is near the origin, there must be another superfield $S$ which has a non-vanishing $F$-term as in the usual SUSY inflation models. The model is similar to the two-field new inflation model in Ref. [16]. The superpotential is given by

$$W = S \left( v^2 - k \frac{\phi^{2\ell}}{M^2} \right),$$

(23)

where $v$ determines the inflation scale and $k$ is a coupling constant. It is possible to make the inflaton mass sufficiently small for a certain Kähler potential so that the inflation takes place. Since $S$ is a gauge singlet field, it does not modify our argument in the previous section. Note that the constant term in the superpotential, $W_0 = m_3/2M_p^2$, does not affect the inflation dynamics in this model, because $S$ is stabilized near the origin during and after inflation. This should be contrasted to the single-field new inflation [17] or hybrid inflation [12] (see also Refs. [18, 19]).

Lastly we note that it is actually possible to cancel the CW potential by tuning the coupling $y_N$, because the contributions from the right-handed neutrinos are accompanied with the minus sign in Eq. (19). This may be an interesting possibility, but it is not certain whether $y_N$ efficiently scans the desired range independently of the gauge coupling in the landscape. If this is the case, however, the inflation model studied in the previous section works, and most of the results concerning the inflationary dynamics (except for SUSY) in this paper remain valid.

Barring cancellation, a similar bound for a SUSY breaking mass for the right-handed sneutrino can be derived. Assuming the coupling $y_N$ of order unity for the heaviest right-handed neutrino, we have

$$\tilde{m}_N \lesssim 0.1 H_{\text{inf}},$$

(24)
where the right-handed sneutrino mass squared is given by $m_{N}^2 + \tilde{m}_{N}^2$, while the right-handed neutrino mass is $m_{N}$. For a generic Kähler potential, $\tilde{m}_{N}$ is expected to be of order the gravitino mass $m_{3/2}$.

4 Cosmological and phenomenological implications

Here we summarize features of our scenario and discuss its implications.

The inflation model based on the minimal set of particles

We have built an inflation model with the inflaton being identified with the Higgs boson responsible for the breaking of the $U(1)_{B-L}$ gauge symmetry. The particle content in our set-up is minimal in some sense: the SM particles, the right-handed neutrinos, $U(1)_{B-L}$ gauge symmetry and its associated Higgs boson $\phi_{B-L}$, and their superpartners at a certain energy scale.

The SUSY breaking scale

The CW potential spoils the inflaton dynamics if the soft SUSY breaking mass of the $B-L$ gaugino is higher than the Hubble scale during inflation. This sets an upper bound on the SUSY breaking scale as $\tilde{m}, \tilde{m}_{N} < 0.1H_{\text{inf}}$.

So far we have not specified how the SUSY breaking in another sector is transmitted to the $U(1)_{B-L}$ gaugino. In the gravity mediation, we expect $\tilde{m} \sim m_{3/2}$, and the soft SUSY breaking masses $m_{s}$ and $m_{\lambda}$ for the SSM particles will be the same order. Here $m_{s}$ and $m_{\lambda}$ collectively represent the soft SUSY breaking mass for scalars and gauginos, respectively. In the anomaly mediation, $\tilde{m}$ as well as $m_{\lambda}$ could be loop-suppressed with respect to the gravitino mass. In particular, since $\tilde{m}_{N}$ is expected to be of order $m_{3/2}$ for a generic (non-sequestered) Kähler potential, we obtain $\tilde{m}_{N} \sim m_{3/2} \gg \tilde{m}, \tilde{m}_{\lambda}$. Such a split mass spectrum is an interesting possibility when there is a bias toward high-scale SUSY. We then expect $m_{3/2} \sim 10^{3}$ TeV, and the SSM gauginos are in the TeV range. In particular, the Wino LSP of mass $\sim 3$ TeV would be a candidate for dark matter (DM) [20].

\footnote{The $B-L$ gaugino mass could be larger if $U(1)_{B-L}$ sector has Planck-suppressed couplings with the SUSY breaking sector, which may be realized in the extra dimensional framework.}
In the gauge mediation, the SUSY breaking mass for $U(1)_{B-L}$ gaugino will be suppressed as $\varphi$ becomes larger. If we require that the inflaton slow-rolls from around the origin to $\varphi_{\text{end}}$, the gravitino mass can be much lower than $\tilde{m}$. In this case, the SUSY mass spectrum is such that all the superpartners of the SM particles are in the TeV or higher, while the gravitino is much lighter and is the lightest SUSY particle (LSP). Therefore the gravitino is a candidate for DM in this case.

**Moduli problem**

In the present scenario the SUSY breaking scale is bounded by the Hubble scale during inflation. If there is a bias toward higher SUSY breaking scale, this bound may be saturated. Then, depending on the mediation mechanism of the SUSY breaking, the gravitino mass $m_{3/2}$ can be comparable to or even slightly larger than $H_{\text{inf}}$. If the modulus mass is of order $m_{3/2}$ or heavier, the modulus abundance is suppressed during inflation. Also the modulus mass is expected to be much heavier than the weak scale, it decays anyway before the big bang nucleosynthesis. Thus, the cosmological moduli problem can be solved, or at least relaxed considerably [21]. Moreover, it may avoid the modulus destabilization during inflation [22] in the KKLT setup [23], if $m_{3/2} \gtrsim H_{\text{inf}}$ is (marginally) satisfied.

**Gravitino problem**

Since the inflaton mass is light in our model, the reheating temperature is relatively low (see Fig. 3). Note that the reheating temperature is correlated with the gravitino mass, and that the cosmological bound on the gravitino is greatly relaxed as $m_{3/2}$ becomes large. The overclosure bound on the LSP abundance produced from the gravitino decay can be avoided if the R-parity is not conserved. Thus, the gravitinos produced by thermal particle scatterings are cosmologically harmless. Furthermore, the non-thermal gravitino production from the inflaton decay [24] is suppressed because the inflaton has a sizable coupling to the right-handed neutrino (see (8)) and the branching fraction into the gravitinos is small.

**Baryogenesis**

The $B-L$ Higgs couples to the right-handed neutrino $\nu_R$ as Eq. (8). The $B-L$ Higgs mass
around the potential minimum is given by

\[ m_\phi^2 = 2\lambda(m - n)\frac{\varphi_{\min}^{2m-2}}{M^{2m-4}}. \]  

(25)

If \( m_\phi > 2m_\nu \), the Higgs can decay into a pair of the right-handed neutrinos with the decay rate

\[ \Gamma_\phi \approx \frac{y_N^2}{8\pi} m_\phi. \]  

(26)

The reheating temperature is then estimated to be

\[ T_R = \left( \frac{10/\pi^2 g^*}{\sqrt{\Gamma_\phi M_P}} \right)^{1/4} \]  

where \( g^* \) counts the relativistic degrees of freedom at \( T = T_R \). We adopt in the following \( g^* = 228.75 \). In Fig. 3 the inflaton mass and the reheating temperature are plotted for \( \varphi_{\min} = 10^{14} \) GeV and \( M_\star = 10^{15} \) GeV (top left), \( \varphi_{\min} = 10^{15} \) GeV and \( M_\star = 10^{15} \) GeV (top right), and \( \varphi_{\min} = 10^{15} \) GeV and \( M_\star = 10^{16} \) GeV (bottom). Here we have assumed that \( m_\nu = 0.1m_\phi \) for simplicity and determined \( y_N \) accordingly. It is seen that the reheating temperature exceeds \( 10^6 \) GeV for \( n \geq 3 \). Thus the baryogenesis through non-thermal leptogenesis naturally works for \( n \geq 3 \) \[25\]. In the case of \( n = 2 \), the reheating temperature is about \( 10^4 \) GeV, and we need to assume degeneracy among the right-handed neutrinos to generate the right amount of the baryon asymmetry.

**Dark matter**

Since SSM particles are likely thermalized after reheating, the LSP can be DM if the R-parity is conserved. For instance, as we have seen, the Wino can be DM in a certain situation. However, there is an argument that the R-parity violation may be a common phenomenon in the string landscape \[26\]. If so, the dangerous operators leading to the proton decay must be absent due to some other reason(s) and the lifetime of LSP in the SSM may be too short to account for the DM\(^8\). Even in this case, the gravitino LSP may serve as DM. The decay of the gravitino DM may provide an observable signature in the the cosmic-ray spectrum \[28\].

\(^8\)It is possible that the lifetime of the Wino LSP is sufficiently long and its decay product contribute to the cosmic-ray spectrum \[27\].
Figure 3: The B–L Higgs mass $m_\phi$ and the reheating temperature $T_R$ with respect to $n$, with $m = n + 1$. The top left is for $\varphi_{\text{min}} = 10^{14}$ GeV and $M_* = 10^{15}$ GeV, the top right is for $\varphi_{\text{min}} = 10^{15}$ GeV and $M_* = 10^{15}$ GeV, and the bottom is for $\varphi_{\text{min}} = 10^{15}$ GeV and $M_* = 10^{16}$ GeV. We have set $m_N = 0.1m_\phi$. In the case of $m = 2n$, the relation between $T_R$ and $n$ is almost the same.
where $f_a$ denotes the PQ scale. This is very small for typical values of $H_{\text{inf}}$ and $f_a$, but if the inflation scale is on the high side, it is marginally consistent with the current observation and it may be detected by the Planck satellite.

The high quality of the PQ symmetry is often considered as a mystery, since any global symmetries are expected to be explicitly broken by Planck-suppressed operators according to the argument on the quantum gravity \[31, 32\]. One explanation is that the QCD axion arises from the string theory axions, and is subject to the other moduli stabilization mechanism \[33\]. There appears a small number in the moduli stabilization, namely the ratio of the gravitino mass to the Planck scale, which could be extremely small if the SUSY persists to low-energy scales. This hierarchy, the gravitino mass and the Planck scale may be responsible for the high-quality of the PQ symmetry. In our framework, therefore, the origin of the high-quality of PQ symmetry could be a result of the inflationary selection.

\textit{The SM Higgs boson mass}

The SM Higgs boson mass weakly depends on the soft masses of the SSM particles \[34, 35, 36\]. In our framework, the SUSY breaking scale can be as large as $10^6$ GeV, and if it is on the high side, the Higgs boson mass will receive sizable radiative corrections and becomes heavier than the case of the weak-scale SUSY. The Higgs boson mass lies within the range of 125 GeV and 155 GeV \[37\], which will be soon checked at the LHC. Although the SUSY particles are beyond the reach of LHC in this case, we may be able to obtain a hint for such a large SUSY breaking from the SM Higgs boson mass.

\textit{Spectral index and tensor mode}

As already calculated in \[14\], the spectral index $n_s$ varies from 0.94 to 0.96 depending on $n$, although it could be increased or decreased by further tuning the potential, which however is not needed. The amplitude of the tensor mode is negligibly small and it cannot be detected by future observations. Also no sizable non-Gaussianity is generated.
5 Discussion and Conclusions

In this paper we have built an inflation model with the inflaton being identified with the Higgs boson responsible for the $U(1)_{B-L}$ gauge symmetry breaking, in the minimal framework $\text{SM} + \nu_R + U(1)_{B-L}$. Our main conclusion is that the soft SUSY breaking masses for the SSM particles should be one order of magnitude smaller than the Hubble parameter during inflation. As a result, the SUSY breaking mass is bounded above as, $m_s \lesssim 1 \text{ TeV} - 10^3 \text{ TeV}$. It is intriguing that, requiring the inflation model based on the $B-L$ Higgs boson to work, the SUSY breaking scale is bounded above and it happens to be close to the weak scale, without relying on the conventional naturalness argument about the gauge hierarchy problem. There are several implications of our finding. First, the universes with the low-scale SUSY may be selected by the inflationary dynamics. Even if there is a bias toward high-scale SUSY in the string landscape, there is a hope that SUSY may be found in the collider experiments such as LHC. Also, the SUSY particles are not very far from the weak scale, even when they are beyond the reach of LHC. It should be emphasized here that this conclusion is derived not relying on the naturalness argument; the low-scale SUSY could emerge as a result of the inflationary selection. Second, even if SUSY is found at the LHC, the fine-tuning issue of obtaining the correct electroweak breaking (the little hierarchy problem) may not be a serious problem any more, because the driving force for the low-scale SUSY is not the fine-tuning issue, but the inflationary dynamics. Thus, the apparent fine-tuning could be a result of combination of the low-scale matter inflation and a bias toward high-scale SUSY.

Our inflation model has interesting cosmological implications. First, since the inflaton is charged under the $U(1)_{B-L}$ symmetry, it is reasonable to expect that the inflaton has a sizable coupling with the right-handed neutrinos, making the non-thermal leptogenesis scenario attractive and viable. Secondly, the inflaton may naturally sit at the origin before the inflation starts because of its gauge interactions with the high-temperature plasma. As the universe cools down, the energy density of the plasma decreases, and the inflation is considered to take place. Thirdly, since the inflationary scale is much lower than the GUT scale, the size of the tensor fluctuation is prohibitively small. The scalar spectral index is expected to be in the range of $n_s = 0.94 \sim 0.96$, but the precise value depends on
the detailed structure of the inflaton potential. The prime dark matter candidate is the QCD axion. Since the inflation scale is low, the isocurvature perturbation constraint on the QCD axion is not stringent; if the inflation scale is on the high side, it is marginally consistent and it may be detected by the Planck satellite.

We note that no topological defects are formed in the present model, since the gauge symmetry is already broken during inflation. This is contrasted to the case of GUT hybrid inflation model, in which the formation of topological defects is inevitable\(^\text{[38]}\).

Since the soft SUSY breaking masses should be in the range between 1 TeV and \(10^3\) TeV, the gauge coupling unification is improved compared to the case without low-scale SUSY. In fact, the unification looks reasonably good if \(m_s \sim 10^3\) TeV and \(m_\lambda \sim \text{TeV}\)\(^\text{[37]}\). Again, this is due to the inflationary selection.

Although we have focused on the B–L Higgs boson as the inflaton, it is straightforward to apply our mechanism to the other GUT Higgs bosons, and we expect similar conclusions about the SUSY breaking scale can be reached. Detailed discussion on this issue is left for future work.

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