PeV-scale Supersymmetry from New Inflation

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Abstract

We show that heavy supersymmetric particles around $O(100)$ TeV to $O(1)$ PeV naturally appear in new inflation in which the Higgs boson responsible for the breaking of $U(1)_{B-L}$ plays the role of inflaton. Most important, the supersymmetric breaking scale is bounded above by the inflationary dynamics, in order to suppress the Coleman-Weinberg potential which would otherwise spoil the slow-roll inflation. Our scenario has rich phenomenological and cosmological implications: the Higgs boson mass at around 125 GeV can be easily explained, non-thermal leptogenesis works automatically, the gravitino production from inflaton decay is suppressed, the dark matter is either the lightest neutralino or the QCD axion, and the upper bound on the inflation scale for the modulus stabilization can be marginally satisfied.
1 Introduction

The concept of symmetry has been a guiding principle in modern physics. For instance, the structure of the standard model (SM) is dictated by the SM gauge symmetries, $SU(3)_C \times SU(2)_L \times U(1)_Y$. The central issue is then how to break symmetry, because clearly we are living in a broken phase: the observed rich structure in our Universe would be impossible in a completely symmetric vacuum. In the celebrated Higgs mechanism [1], gauge symmetry is spontaneously broken by a vacuum expectation value (VEV) of a Higgs field.

Recently the ATLAS and CMS collaborations have provided hints for the existence of a SM-like Higgs particle with mass of about 125 GeV [2]. The relatively light Higgs boson mass suggests the presence of new physics at scales below the Planck scale [3]. In a supersymmetric (SUSY) extension of the SM (SSM), the Higgs boson mass can be explained if the typical sparticle mass is at $O(10)$ TeV or heavier. This casts doubt on the conventional naturalness argument as the correct guiding principle for understanding the physics at and beyond the weak scale.

Once the existence of the SM-like Higgs boson is confirmed, it would immediately mean that the Higgs mechanism is indeed realized in nature, and some other gauge symmetries may be broken in a similar manner. Those symmetries may have been restored in the past because the Universe was much hotter and denser at early times. Thus probably our Universe experienced a series of phase transitions in course of its evolution.

The inflationary paradigm has been well established so far [4]. Despite its great success, it is not yet known what the inflaton is. It is natural to expect that one of the Higgs fields which trigger phase transitions in the early Universe is responsible for the inflation. Indeed, this possibility was extensively discussed in the early 80’s under the name of new inflation [5]. The phase transition in the new inflation was of Coleman-Weinberg (CW) type [6], where the inflaton was the Grand Unification Theory (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, $\delta \rho / \rho \sim 10^{-5}$ [7]. 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, such gauge singlets are ubiquitous in the string theory, and so, one of them may be responsible for the inflation. Another way to resolve the problem was to introduce SUSY [8]. Then the CW potential becomes suppressed because of the cancellation among bosonic and fermionic degrees of freedom running the loop.

Recently the present authors proposed a new inflation model in which a Higgs field responsible for the breaking of $U(1)_{B-L}$ symmetry plays the role of inflaton [9]. It was found that the SUSY must be a good symmetry at scales below the Hubble parameter during inflation\(^1\). Interestingly, we obtained an upper bound on the soft SUSY breaking mass, $\tilde{m} \lesssim O(10)$ TeV $- O(1)$ PeV for the $U(1)_{B-L}$ breaking scale of $10^{15}$ GeV inferred from the neutrino oscillation data [9]. Furthermore, the inflaton predominantly decays into a pair of right-handed neutrinos, and non-thermal leptogenesis [10, 11] works almost automatically. The implication for the SM-like Higgs boson mass in this framework was studied in Ref. [12].

In this paper, we study the inflationary dynamics of the $U(1)_{B-L}$ new inflation as well as its subsequent thermal history of the $U(1)_{B-L}$ new inflation in detail. The spectral index is calculated with a greater accuracy and found to be perfectly consistent with the current WMAP data, $n_s \simeq 0.968 \pm 0.012$ [13]. In particular we will see that the CW potential will play an important role to increase $n_s$ to provide a better fit to the WMAP data. We will consider the implication of the SUSY breaking mass from the inflationary dynamics for the SM-like Higgs boson mass. We also discuss various phenomenological and cosmological implications such as non-thermal leptogenesis, gravitino production from inflaton decay, dark matter (DM), the Polonyi problem, and the modulus destabilization problem [14]. It is noteworthy that in such a minimal extension of the SSM, the observed data such as the spectral index of the density perturbation and the SM-like Higgs boson mass can be explained while naturally creating the right amount of the baryon asymmetry

\(^1\) In Ref. [8], the soft SUSY breaking mass was (implicitly) assumed to scale in proportion to the GUT Higgs boson. In their Eq. (8), the dependence of the CW potential on the GUT Higgs boson was factored out, and then they substituted the mass splitting relation Eq. (21) into Eq. (8). In effect, this is equivalent to assuming that the soft SUSY breaking mass is proportional to the GUT Higgs boson VEV. Therefore the upper bound on the SUSY breaking was overestimated, and it was actually higher than the Hubble parameter during inflation, which clearly does not make sense, because it would mean that there is no SUSY at the inflation scale. To our knowledge, this error was not corrected until Ref. [9].
without the gravitino and Polonyi problems.

The rest of the paper is organized as follows. We will briefly review how the SUSY breaking is bounded from above in the new inflation in Sec. 2 and derive an important upper bound on the soft SUSY breaking mass. In Sec. 3 we discuss the dynamics of the U(1)$_{B-L}$ new inflation in detail. In Sec. 4 we discuss the reheating of the inflaton. The implications for the SM-like Higgs boson mass is discussed in Sec. 5. We discuss various implications of our scenario in Sec. 6. The last section is devoted for conclusions.

2 Upper bound on SUSY breaking

Let us briefly review how the SUSY breaking is bounded above for the successful new inflation using a gauge non-singlet inflaton, following Ref. [9]. The bound essentially comes from the requirement that the radiative correction to the inflation potential should be suppressed since otherwise the slow-roll inflation would not last long enough and the density perturbation would be too large.

Consider a Higgs boson $\phi$ responsible for the breaking of U(1)$_{B-L}$ symmetry. In the new inflationary scenario, the inflaton sits near the origin at the beginning of inflation. If the inflaton potential is sufficiently flat around the origin, the inflation takes place. As $\phi$ is charged under the U(1)$_{B-L}$ symmetry, the inflaton potential receives a radiative correction from the gauge boson loop. The general form of the CW effective potential is given by [6]

$$V_{CW} = \frac{1}{64\pi^2} \sum_i (2S + 1)(-1)^{2S}M_i^4(\phi) \ln \left( \frac{M_i^2(\phi)}{\mu^2} \right),$$

where $\mu$ is the renormalization scale, and the mass eigenvalues of the particles coupled to $\phi$ are represented by $M_i(\phi)$. Since the mass of the U(1)$_{B-L}$ gauge boson is given by $m_{GB} = \sqrt{2}g_{B-L}q_\phi \langle \phi \rangle$, the inflaton potential receives the CW correction as

$$V_{CW, gauge}(\sigma) = \frac{3}{64\pi^2}g_{B-L}^4q_\phi^4\sigma^4 \ln \left( \frac{g_{B-L}^2q_\phi^2\sigma^2}{\mu^2} \right),$$

where $g_{B-L}$ represents the gauge coupling of U(1)$_{B-L}$, $q_\phi$ is the U(1)$_{B-L}$ charge of $\phi$, and $\sigma$ denotes the radial component of $\phi$, $\sigma \equiv \sqrt{2} |\phi|$. 
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 [7]. One plausible way to solve the problem is to introduce SUSY [8]. In the exact SUSY limit, contributions from boson loops and fermion loops are exactly canceled out. However, if SUSY is broken, we are left with non-vanishing CW corrections, which are estimated below.

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 $\varphi$. We can simply relate $\Phi$ and $\bar{\Phi}$ to $\varphi$ as $|\Phi| = |\bar{\Phi}| = |\varphi|/\sqrt{2}$. The U(1)$_{B-L}$ charge of $\varphi$ is set to be $q_{\varphi} = 2$ in the following.

The gauge boson has mass of $m_S^2 = g_{B-L}^2 q_{\varphi}^2 \sigma^2$, where $\sigma \equiv \sqrt{2} |\varphi|$. On the other hand, there are additional fermionic degrees of freedom, the U(1)$_{B-L}$ gaugino and higgsino, whose mass eigenvalues are given by $m_F = g_{B-L} q_{\varphi} \sigma \pm \frac{1}{2} M_\lambda$, where $M_\lambda$ denotes the soft SUSY breaking mass for the U(1)$_{B-L}$ gaugino. Because of the SUSY breaking mass $M_\lambda$, the CW potential does not vanish and the inflaton receives a non-zero correction to its potential. Inserting the field dependent masses into the CW potential (1), and expanding it by $M_\lambda/(g q_{\varphi} \sigma)$, we find

$$V_{\text{CW, gauge}}^{\text{susy}}(\sigma) \simeq -\frac{3g_{B-L}^2}{8\pi^2} \left(\frac{q_{\varphi}}{2}\right)^2 M_\lambda^2 \sigma^2 \ln \left(\frac{g_{B-L}^2 q_{\varphi}^2 \sigma^2}{\mu^2}\right),$$

where we have also taken into account of the inflaton as well as the scalar perpendicular to the D-flat direction. 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 $M_\lambda \ll g_{B-L} q_{\varphi} \sigma$, in contrast to the result of Ref. [8]. Note that the correction still contains a logarithmic factor, which may not be negligible if we consider the whole evolution of the inflaton.

For successful inflation, we require the curvature of the CW potential (3) to be at least one order of magnitude smaller than $H_{\text{inf}}$ for $\sigma \lesssim \sigma_{\text{end}}$. Here $H_{\text{inf}}$ is the Hubble parameter during inflation, and $\sigma_{\text{end}}$ is the point where the slow-roll condition breaks down and the inflation ends. Therefore, we obtain the following constraint on the soft SUSY breaking
mass for the $U(1)_{B-L}$ gaugino:

$$g_{B-L}M_\lambda \lesssim O(0.1)H_{\text{inf}}.$$  

(4)

For the gauge coupling of order unity, this bound reads $M_\lambda \lesssim O(0.1)H_{\text{inf}}$.

The $U(1)_{B-L}$ Higgs boson is also coupled to the right-handed neutrinos to give a large Majorana mass. We consider the following interaction,

$$- \mathcal{L} = \sum_i \frac{y_{\phi,i}}{2} \bar{\nu}_{R,i}^c \nu_{R,i} + \text{h.c.},$$  

(5)

where the subscript $i$ represents the generation. The right-handed neutrino mass is given by $M_{N,i} = y_{\phi,i} \sigma/\sqrt{2}$. This interaction similarly contributes to the CW potential as

$$V_{\text{CW},N} = -\sum_i \frac{1}{8\pi^2} y_{\phi,i}^4 \sigma^4 \ln \left( \frac{2y_{\phi,i}^2 \sigma^2}{\mu^2} \right).$$  

(6)

This can similarly spoil the inflationary dynamics. In the presence of SUSY, there are right-handed sneutrinos. Let us write its mass as $M_{N,i}^2 = (\sqrt{2}y_{\phi,i} \sigma)^2 + m_{N,i}^2$, where $m_{N,i}$ represents the soft SUSY breaking mass for the right-handed sneutrinos. The CW potential is then

$$V_{\text{CW},N}^{\text{susy}}(\sigma) = \sum_i \frac{y_{\phi,i}^2}{8\pi^2} m_{N,i}^2 \sigma^2 \ln \left( \frac{2y_{\phi,i}^2 \sigma^2}{\mu^2} \right).$$  

(7)

For successful inflation, the soft mass is bounded above as before:

$$\sqrt{\sum_i y_{\phi,i}^2 m_{N,i}^2} \lesssim O(0.1)H_{\text{inf}}.$$  

(8)

If the Yukawa coupling for the heaviest right-handed neutrino $\nu_{R,3}$ is of order unity, the bound reads $m_{N,3} \lesssim O(0.1)H_{\text{inf}}$.

In the gravity mediation, $M_\lambda$, $m_{N,i}$ as well as the soft SUSY masses for the SSM particles are considered to be comparable to the gravitino mass $m_{3/2}$. On the other hand, in anomaly mediation [15], they may be suppressed compared to the gravitino mass, but for a generic form of the Kähler potential, $m_{\tilde{N}}$ and the sfermion masses are

\[\text{In principle it is possible to cancel} \ V_{\text{CW,gauge}} \text{with} \ V_{\text{CW,N}} \text{by fine-tuning the Yukawa coupling} \ y_{\phi,i}. \]

\[\text{In this case, the successful inflation takes place without SUSY, and the above upper bound on the soft} \]

\[\text{SUSY breaking mass does not hold. We do not consider this case further in this paper.}\]
comparable to the gravitino mass. We assume the latter case when we consider the case of anomaly mediation. On the other hand, in the gauge mediation, the relation between the soft masses and the gravitino mass is model-dependent, and we do not consider gauge mediation in this paper.

The inflation places a robust upper bound on the soft SUSY breaking parameter of the $U(1)_{B-L}$ gaugino and the right-handed sneutrino. In particular, for $g_{B-L}$ and $\sqrt{\sum_i g_{\tilde{\psi},i}^2}$ of order unity, both $M_\lambda$ and $m_{\tilde{N},i}$ should be smaller than $H_{\text{inf}}$. Furthermore, as long as $M_\lambda$ and $m_{\tilde{N},i}$ are comparable to the soft SUSY breaking mass for the SSM particles $\tilde{m}$ as in the gravity or anomaly mediation we obtain

$$\tilde{m} \lesssim O(0.1)H_{\text{inf}},$$

which relates the inflation scale to the SUSY breaking.\footnote{$\tilde{m}$ should be considered as representing the sfermion mass, if the gaugino mass is suppressed as in the anomaly mediation.} As we shall see later, the inflation scale varies from $10^6$ GeV ($n = 2$) to $10^{10}$ GeV or heavier ($n \geq 3$). (See Eq. (11) for the definition of the power $n$.) We will focus on the simplest case of $n = 2$, because it provides an interesting upper bound on $\tilde{m}$ and because the VEV of the inflaton is very close to the see-saw scale $\sim 10^{15}$ GeV suggested by the neutrino oscillation data. We shall see that in the case of $n \geq 3$ some of the nice features of the model are preserved, although the direct connection between the inflaton VEV and the see-saw scale is lost.

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 may constrain the SUSY breaking scale to be below the inflation scale. Also, in this case we have a prediction that the SUSY breaking scale should be close to the inflation scale. We assume that this is the case, because, if it is biased to lower SUSY breaking scale, we should have already seen SUSY particles at the collider experiments. Interestingly, as we will see, the observed value of the scalar spectral index even suggests that the upper bound is saturated. Even if the SUSY particles are too heavy to be discovered at the LHC, we may be able to see the hint for the SUSY breaking scale much higher than the electroweak breaking from the large radiative correction to the SM Higgs boson mass.\footnote{\cite{16}}
We will come back to this issue in Sec. 5.

Lastly let us mention the applicability of the inequality (9). As is clear from the derivation, the upper bound on SUSY breaking derived applies to any inflation models in which there are fields coupled to the inflaton with a coupling of order unity, and they have inflaton-dependent mass. In particular, this is the case if the inflaton is charged under gauge symmetry or if the inflaton has a Yukawa coupling with fermions, as we have seen above. Note that it is applicable to the gauge-singlet inflation models, if the inflaton has a sizable Yukawa coupling like (5).

3 \( U(1)_{B-L} \) New Inflation

In the previous section we have seen that the SUSY must be a good symmetry at the inflation scale. Therefore the inflation sector can be described in a supersymmetric Lagrangian.

The Kähler and super-potentials for the inflation are given by [17]

\[
K = |\Phi|^2 + |\bar{\Phi}|^2 + |\chi|^2 + \frac{k_1}{4}|\Phi|^4 + \frac{k_2}{4}|\bar{\Phi}|^4 + k_3|\Phi|^2|\chi|^2 + k_4|\bar{\Phi}|^2|\chi|^2 + \frac{1}{4}k_5|\chi|^4 + \cdots ,
\]

\[
W = \chi \left( v^2 - g(\Phi\bar{\Phi})^n \right),
\]

where \( k_i (i = 1 - 5) \) and \( g \) represent a coupling constant of order unity, \( \cdots \) denotes higher order terms and we adopt the Planck unit, \( M_p \approx 2.4 \times 10^{18} \text{GeV} = 1 \). The charge assignment of \( \Phi, \bar{\Phi} \) and \( \chi \) are shown in Table 1. Note that we have introduced a discrete \( Z_n \) symmetry under which only \( \bar{\Phi} \) is charged. Such a discrete symmetry is necessary to ensure a flat potential for the inflaton.

The \( U(1)_{B-L} \) and other symmetries may be restored in the early Universe, because of the thermal mass and/or the Hubble-induced mass. If so, the origin \( \Phi = \bar{\Phi} = 0 \) is chosen as the initial condition. As the Universe expands, the temperature and the Hubble parameter decrease, and finally the inflation takes place when the inflaton potential dominates the energy density of the Universe, if the inflaton potential is sufficiently flat.

The CW potential, which could spoil the slow-roll inflation, can be sufficiently suppressed if the typical SUSY breaking mass of the \( U(1)_{B-L} \) gaugino and the right-handed
Table 1: The charge assignment for $\Phi$, $\bar{\Phi}$ and $\chi$.

|           | $\Phi$ | $\bar{\Phi}$ | $\chi$ |
|-----------|--------|---------------|--------|
| $U(1)_{B-L}$ | -2     | 2             | 0      |
| $U(1)_R$     | 0      | 0             | 2      |
| $Z_n$        | 0      | 1             | 0      |

The neutrino is (much) smaller than the inflation scale. We assume that this is the case for the moment and consider the supersymmetric part of the inflaton potential. We shall discuss the effect of the CW potential ((3) and (7)) on the inflation dynamics, especially on the spectral index $n_s$, later in this section. The effect of a constant term in the superpotential was studied in Ref. [9]; assuming $|k_5| = O(1)$, it was found that the inflaton dynamics is not affected as long as $m_{3/2} \lesssim O(0.1)H_{\text{inf}}$, which is similar to Eq. (9). We will come back to this issue in Sec. 6.

For a field value greater than the Hubble parameter during inflation, the D-term potential forces $\Phi$ and $\bar{\Phi}$ to be along the D-flat direction, $|\Phi| = |\bar{\Phi}| = \frac{1}{\sqrt{2}}|\varphi|$, where $\varphi$ is a complex scalar field. Focusing on the radial component, $\sigma \equiv \sqrt{2}|\varphi|$, the scalar potential is approximately given by

$$V(\sigma, \chi) \simeq v^4 - \frac{1}{2} \left( \frac{k_3 + k_4 - 2}{4} \right) v^4 \sigma^2 - \frac{g}{2^{2m-1}} v^2 \sigma^{2n} + \frac{g^2}{2^{4m}} \sigma^{4n} - k_5 v^4 |\chi|^2.$$  \hspace{1cm} (12)

We assume $k_5 < -\frac{3}{4}$ so that $\chi$ is stabilized at the origin during inflation. In order for the slow-roll inflation to take place, we also require the inflaton mass term is much smaller than the Hubble parameter,

$$k \equiv \frac{k_3 + k_4 - 2}{4} \lesssim O(0.01).$$ \hspace{1cm} (13)

The tuning of the inflaton mass is known as the $\eta$-problem. We do not care about this fine-tuning at the level of 1%, because it can be easily compensated by the subsequent exponential expansion and because perhaps we cannot live in an Universe which has not experienced inflation. We note that, in general, $k_3$ and $k_4$ do not have to be small, and we expect them to be of order unity. \footnote{Since either $k_3$ or $k_4$ is likely greater than unity, either $\Phi$ or $\bar{\Phi}$ acquires a tachyonic mass in the vicinity of the origin, developing a local minimum. This may make the eternal inflation more likely.} The inflation dynamics in this model is same as in the single-field new inflation model [18], which was studied in detail in Ref. [19].
Let us rewrite the inflaton potential, assuming $\chi$ is stabilized at the origin:

$$V(\sigma) \simeq v^4 - \frac{1}{2} kv^4 \sigma^2 - \frac{g}{2^{n-1}} v^2 \sigma^{2n} + \frac{g^2}{2^{4n}} \sigma^{4n}.$$  \hspace{1cm} (14)

After inflation, the inflaton $\sigma$ is stabilized at the potential minimum given by

$$\sigma_{\text{min}} \simeq 2 \left( \frac{v^2}{g} \right)^{\frac{1}{2n}}.$$  \hspace{1cm} (15)

The $U(1)_{B-L}$ symmetry is spontaneously broken by the inflaton vev. We define the breaking scale as

$$v_{B-L} \equiv \sigma_{\text{min}} \sqrt{2} = \sqrt{2} \left( \frac{v^2}{g} \right)^{\frac{1}{2n}}.$$  \hspace{1cm} (16)

Note that $v_{B-L}$ cannot take an arbitrary value because the coupling $g$ should not be much larger than $O(1)$ for the Kähler potential Eq. (10) to be valid.

In order to estimate the Hubble parameter during inflation, we need to solve the inflation dynamics and estimate the density perturbation. When the inflaton sits near the origin, the slow-roll inflation takes place. As the inflaton rolls down on the potential, the curvature of the potential becomes gradually non-negligible, and finally the slow-roll inflation ends when one of the slow-roll parameters, $\eta$, becomes order unity. The $\eta$ is given by

$$\eta \equiv \frac{V''(\sigma)}{V(\sigma)} \simeq -k - \frac{n(2n-1)g}{2^{n-1}v^2} \sigma^{2(n-1)},$$  \hspace{1cm} (17)

and $|\eta|$ becomes unity at $\sigma = \sigma_{\text{end}}$, which is given by

$$\sigma_{\text{end}} \approx 2 \left( \frac{(1-k)v^2}{n(2n-1)g} \right)^{\frac{1}{2(n-1)}}.$$  \hspace{1cm} (18)

Under the slow-roll approximation, the equation of motion for the inflaton is given by

$$3H \frac{d\sigma}{dt} + V'(\sigma) \approx 0,$$  \hspace{1cm} (19)

or equivalently

$$3H^2 \frac{d\sigma}{dN} + V'(\sigma) \approx 0,$$  \hspace{1cm} (20)
where \( N \) denotes the e-folding number. Solving this equation of motion we obtain

\[
\sigma(N) \approx 2 \left( \frac{k v^2}{n g} \right)^{\frac{1}{2(n-1)}} G(k, n, N)^{-\frac{1}{2(n-1)}},
\]

(21)

where we have defined

\[
G(k, n, N) \equiv e^{2(n-1)N} \left( 1 + \frac{k}{1-k} (2n-1) \right) - 1.
\]

(22)

The curvature perturbation can be expressed in terms of the inflaton potential,

\[
\Delta_R^2 = \frac{1}{12 \pi^2} \frac{V(\sigma)^3}{V'(\sigma)^2} = (2.430 \pm 0.091) \times 10^{-9},
\]

(23)

where we have used the WMAP normalization in the second equality \[13\]. The Hubble parameter during inflation is given by

\[
H_{\text{inf}} \simeq \frac{v^2}{\sqrt{3}} \simeq \sqrt{\frac{\Delta_R^2}{\mathcal{F}(k, n, N)} v_{B-L}^{\frac{n-3}{n-1}}},
\]

(24)

with

\[
\mathcal{F}(k, n, N) \equiv \pi \left( \frac{2^{3n-4} k^{2n-1}}{n G(k, n, N)} \right)^{\frac{1}{2(n-1)}} (1 + G(k, n, N)^{-1}).
\]

(25)

In Fig. 1 we show the function \( \mathcal{F}(k, n, N) \) with respect to \( k \) for several values of \( n \) with \( N = 50 \). We can see that \( \mathcal{F}(k, n, N) \) is about 0.01 for the ranges of the parameters of interest. This is not significantly modified for \( N = 40 \) or 60.

Requiring \( g \lesssim O(1) \), we obtain a lower bound on \( v_{B-L} \):

\[
v_{B-L} \gtrsim 2^n \sqrt{3 \Delta_R^2 \mathcal{F}(k, n, N)} v_{B-L}^{\frac{n-3}{n(n-1)}},
\]

(26)

which is shown in Fig. 2. To be concrete we take \( g = 1 \) in the following analysis, and in this case the bound on \( v_{B-L} \) is saturated. Note that the case of \( n = 2 \) is particularly interesting because the \( U(1)_{B-L} \) breaking scale is close to the see-saw scale inferred from the neutrino oscillation.

The Hubble parameter during inflation is shown in Fig. 3. Considering that the soft mass for the SSM particles should be smaller than the Hubble parameter for the successful inflation to take place, the cases of \( n = 2 \) and \( n = 3 \) are interesting, especially from the point of view of explaining the Higgs mass at around 125 GeV.
Figure 1: The behavior of the function $F(k, n, N)$, where we set $N = 50$.

The inflaton mass at the potential minimum is given by

$$m_\sigma \approx \sqrt{2n}v^2 \left( \frac{\nu^2}{g} \right)^{-\frac{1}{n}} = \frac{2\sqrt{3}nH_{\text{inf}}}{v_{\text{B-L}}} \approx 1.7 \times 10^{-6} n \left( \frac{F(k, n, N)}{10^{-2}} \right) v_{\text{B-L}}^{-1} \quad (27)$$

In Fig. 4 we show the inflaton mass at the potential minimum as a function of $k$ for several values of $n$. For $n \geq 3$, the inflaton mass is greater than about $10^{12}$ GeV.

Lastly let us estimate the spectral index $n_s$, which is approximately given by

$$n_s \approx 1 + 2\eta = 1 - 2k \left( 1 + (2n - 1)G(k, n, N)^{-1} \right). \quad (28)$$

We show the spectral index $n_s$ in Fig. 5 as a function of $k$. The limit $k \to 0$ reproduces the result of Ref. [9]. In principle $k$ can be extremely small, which however requires severer fine-tuning of the parameters. If the fine-tuning is just what is needed for the inflation to take place, we may expect $k$ to be of 0.01. Then, the current WMAP 7yr data $n_s = 0.968 \pm 0.012$ [13] is perfectly consistent with $n \geq 3$, independent of the $U(1)_{\text{B-L}}$ breaking scale.

We note that the spectral index is between 0.94 and 0.95 in the case of $n = 2$, which is slightly smaller than the observed value, causing a tension at $2\sigma$ level. However, we should emphasize here that the above result is derived from the potential (14). As we discussed before, there is a finite contribution from the CW potential once the SUSY breaking is taken into account. Let us take account of the effect by adding the following $V_{\text{SB}}(\sigma)$ to the inflaton potential:

$$V_{\text{SB}} = \frac{1}{2}k'v^4\sigma^2 \log \left( \frac{\sigma}{\sigma_0} \right), \quad (29)$$
Figure 2: The lower bound on the $v_{B-L}$ as a function of $k$ for $n = 2, 3, 4$ and 5. We set $N = 50$.

where $k'$ represents the SUSY breaking, and $\sigma_0$ is the renormalization scale. Using the result in Sec. 2 it is given by

$$k' \equiv \frac{1}{6\pi^2 H_{\inf}^2} \left( \sum_i y_{\varphi,i}^4 m_{\varphi,i}^2 - 3g_{B-L}^2 \left( \frac{q_{\varphi}}{2} \right)^2 M_{\chi}^2 \right). \quad (30)$$

In order for the curvature of the potential to be smaller than the Hubble parameter for $\sigma \lesssim \sigma_{\text{end}}$, $k$ and $k'$ should be smaller than $\sim 0.1$. Note that $k$ is redefined here so that the total potential is given by $V(\sigma) + V_{SB}(\sigma)$. The logarithmic correction slightly changes the global shape of the inflaton potential, and as a result the predicted value of $n_s$ is modified while the other inflation parameters are not significantly changed. We have numerically solved the inflaton dynamics and estimated the spectral index at the pivot scale. We have fixed $\sigma_0 = 10^{-7}$ for simplicity. In Fig. 6 we show the contour of $n_s$ in the case of $n = 2$, where the WMAP normalization is satisfied by slightly varying the value of $v$ and $v_{B-L}$. The values of $v_{B-L}$ and $H_{\inf}$ varies from $3 \times 10^{15}$ GeV to $4 \times 10^{15}$ GeV, and $1 \times 10^6$ GeV to $3 \times 10^6$ GeV, respectively, in the region shown in Fig. 6. We can see that $n_s$ can be increased up to about 0.98 in the presence of the CW correction\footnote{We have confirmed that $n_s$ can be increased to $\sim 0.99$ by further increasing $k$ and $k'$. This is one of the main differences from the previous works on the two-field new inflation.} As we increase
Figure 3: The Hubble parameter during inflation as a function of $k$ for $n = 2, 3, 4$ and $5$ and $N = 50$. $v_{\text{B-L}}$ is given by Fig. 2.

$k'$, the mass near the origin becomes more negative, while the potential becomes flatter as the inflaton goes away from the origin. This two effects explains the behavior of $n_s$ in Fig. 6. We have also confirmed that the total e-folding number is greater than 100 in the region shown in the figure.

Note that $k' = O(0.01)$ requires one of the right-handed neutrinos to have a mass comparable to the inflaton VEV and that the inequality on the SUSY breaking is saturated. It is interesting that including the CW correction gives a better fit to the observed value of $n_s$ in the case of $n = 2$ where the suggested see-saw scale of $O(10^{15})$ GeV is close to the inflaton VEV. Note also that the addition of the potential (29) may create a local minimum along the inflaton trajectory, which spoil the successful inflaton dynamics.

In order to avoid this, we demand

$$k' \left[ \frac{2 - n}{2(n - 1)} - \log \left( \frac{\tilde{\sigma}}{\sigma_0} \right) \right] + k > 0,$$

where we have defined

$$\tilde{\sigma} \equiv 2 \left( \frac{k' v^2}{2 n(n - 1) g} \right)^{\frac{1}{2(n-1)}}.$$

This condition is violated in the upper left shaded region of Fig. 6.
Figure 4: The inflaton mass as a function of $k$ for $n = 2, 3, 4$ and $5$ and $N = 50$. $v_{B-L}$ is given by Fig. 2.

In the case of $n = 3$, there is a little hierarchy between the inflaton VEV of $O(0.1) = O(10^{17})$ GeV and the see-saw scale of $O(10^{15})$ GeV. This tension can be nicely explained by changing the $Z_3$ assignment as $\Phi(+1)$ and $\Phi(-1)$. Then the couplings of $\Phi$ to the right-handed neutrinos are given by

$$W = \frac{y_{\Phi,i}}{2} (\Phi \Phi) \Phi N_i N_i,$$

in order to satisfy the $Z_3$ symmetry. Then we can explain this hierarchy naturally, $10^{15}$ GeV $\sim 10^{-2} \cdot 10^{17}$ GeV.

4 Reheating

After the inflation, the inflaton must release its energy into radiation including the SM particles, which is called the reheating. In gauge-singlet inflation models, it is highly non-trivial if the inflaton successfully reheats the SM sector. In the supergravity framework, it was shown in Ref. [20] that the inflaton is coupled to any sector via the Planck-suppressed interactions if the inflaton has non-zero VEV, providing a robust lower bound on the reheating temperature. At the same time, however, the inflaton would decay into unwanted...
Figure 5: The spectral index for $n = 2, 3, 4$ and 5 in the SUSY limit. The shaded region shows the 1σ allowed range, $n_s = 0.968 \pm 0.012$, by the WMAP 7yr data [13]. Note that $n_s$ is independent of $v_{B-L}$.

Figure 6: Contours of the spectral index for $n = 2$, taking account of the SUSY breaking represented by $k'$ (See Eq. (29)). In the upper left shaded region, the inflation does not end successfully. Note that $k, k' \lesssim 0.1$ must be satisfied in order for the curvature of the potential to be smaller than the Hubble parameter for $\sigma \lesssim \sigma_{\text{end}}$. 
relics such as gravitinos at a non-negligible rate \cite{21, 22, 23}, causing severe cosmological problem.

In our present model, the inflaton is charged under the U(1)$_{B-L}$ symmetry, and it naturally has a coupling to the right-handed neutrinos,

$$W = \frac{y_{\Phi,i}}{2} \Phi N_i N_i,$$  \quad (34)

where $N_i$ denotes the right-handed neutrino chiral superfield of the $i$-th generation, and $y_{\Phi,i}$ corresponds to $\sqrt{2} y_{\phi,i}$ in Eq. (5). After inflation, $\Phi$ develops a VEV, and the U(1)$_{B-L}$ gets spontaneously broken. The U(1)$_{B-L}$ breaking naturally gives rise to the heavy Majorana mass $M_i \equiv y_{\Phi,i} v_{B-L} / \sqrt{2}$ for the right-handed neutrinos, as required by the see-saw mechanism \cite{24} for the light neutrino mass. The above interaction induces the inflaton decay into a pair of the right-handed neutrinos, suppressing the gravitino production.\footnote{In fact, low-scale inflation with a sizable coupling to the visible sector is favored since it suppresses the non-thermal gravitino production \cite{21, 22, 23}.}

The decay into the right-handed sneutrinos proceeds at the same rate \cite{25}. Let us comment on this process, because it is often claimed that this decay process is suppressed compared to that into right-handed neutrinos. Taking the $F$-term of $\Phi$ in the interaction (34) and expanding it in terms of $\chi$, we obtain

$$\mathcal{L} \supset -m_\sigma \frac{y_{\Phi,i}}{2\sqrt{2}} \chi \tilde{N}_i \tilde{N}_i + \text{h.c..}$$  \quad (35)

where we have used $W_{\chi\phi} \simeq -m_\sigma$. At the first sight, it seems that this interaction does not induce the inflaton decay, however, it was shown in Ref. \cite{21} that $\chi$ and $\phi$ gets almost maximally mixed due to the constant term in the superpotential. The mass eigenstates are given by $(\phi \pm \chi^\dagger)/\sqrt{2}$. Therefore, through the mixing, the inflaton decays into the right-handed sneutrinos.

The decay rate of the inflaton into the lightest right-handed (s)neutrinos is given by

$$\Gamma_{\text{inf}}(\text{inflaton} \rightarrow N_1 N_1, \tilde{N}_1 \tilde{N}_1) \simeq \frac{|y_{\phi1}|^2}{64\pi} m_\sigma = \frac{1}{32\pi} \frac{M_1^2 m_\sigma}{v_{B-L}^2},$$  \quad (36)

for $m_\sigma > 2M_1$. Here we have taken account of the mixing between $\phi$ and $\chi$ \cite{21}. The reheating temperature is defined as

$$T_R = \left( \frac{\pi^2 g_*}{90} \right)^{-\frac{1}{4}} \sqrt{\Gamma_\Phi M_p},$$  \quad (37)
Figure 7: The reheating temperature as a function of $M_1$. We set $g_* = 228.75$, $k = 0.01$ and $N = 50$.

where $g_*$ counts the relativistic degrees of freedom at the reheating. In Fig. 7 we show the reheating temperature as a function of $M_1$ for $n = 2, 3, 4,$ and $5$. We set $k = 0.01$ and $N = 50$, and consider the inflation model in the SUSY limit since the effect of the SUSY breaking on the reheating temperature is small. Note that the reheating temperature is so high that non-thermal leptogenesis may work for $M_1 > 10^9$ GeV and $n \geq 2$. We will come back to this issue in Sec. 6.1.

5 The SM-like Higgs boson mass

The SUSY breaking scale is bounded above by the inflationary dynamics. If it is saturated, the typical SUSY breaking scale is of $O(10^5)$ GeV to $O(10^6)$ GeV for $n = 2$, and $O(10^9)$ GeV to $O(10^{10})$ GeV for $n = 3$. It would be difficult to directly produce such heavy SUSY particles at collider experiments. However, we may be able to see a hint for such high-scale SUSY breaking from the large radiative corrections to the SM-like Higgs boson mass. In order to calculate the SM-like Higgs boson mass, we need to specify $\tan \beta$, the SUSY mass spectrum, and the stop mixing parameter, where $\tan \beta$ is the ratio of the up- and down-type Higgs boson VEVs. In the following we set the stop mixing parameter to be zero for simplicity.
The possible mass spectrum can be broadly divided into the following two cases: (1) high-scale SUSY with all the SUSY particles having a mass comparable to \( \tilde{m} \), or (2) split spectrum in which the sfermion mass is of order \( \tilde{m} \) while the gauginos and the higgsino are at around the weak scale (or slightly higher). The first possibility corresponds to the gravity mediation, which requires a singlet SUSY breaking field to give a gaugino mass. The latter can be realized in simple anomaly mediation \[15\] with a generic form of the Kähler potential for \( n = 2 \). In the case of \( n = 3 \), we need a certain mechanism to turn off the anomaly mediation contribution to the gaugino mass. In fact, if we take the hint for the Higgs at around 125 GeV seriously, only \( n = 2 \) is allowed for the case (2). Also, in the case of \( n = 2 \), the allowed region is similar for the cases (1) and (2). Therefore we consider the case of (1) with \( n = 2 \) and \( n = 3 \) in the following.

We have calculated the SM-like Higgs boson mass following Ref. \[16\]. The contours of the Higgs boson mass \( m_H \) are shown in Fig. 8. We can see that the Higgs boson at around 125 GeV suggested by the recent ATLAS and CMS experiments can be explained for \( \tan \beta = 3 - 5 \) and \( \tan \beta = 1 \sim 1.5 \) for \( n = 2 \) and \( n = 3 \), respectively.

The Higgs mass at about 125 GeV suggests a relatively high (but not extremely high) SUSY breaking in the minimal extension of the SSM. For \( \tan \beta \gtrsim 1 \), it varies from \( 10^4 \) GeV to \( 10^{10} \) GeV \[16\]. It is a puzzle why the SUSY should appear at such scale, which is higher than the electroweak scale making the fine-tuning severe, while it is much smaller than the fundamental energy scale such as the GUT or Planck scales. Our scenario provides a possible solution to this issue: this may be due to the inflationary selection. Namely, the apparent fine-tuning could be a result of combination of the \( U(1)_{B-L} \) new inflation and a bias toward high-scale SUSY in the landscape.

6 Cosmological and phenomenological implications

In this section, we discuss various cosmological and phenomenological implications of our scenario. Before going further, let us briefly mention the cosmology in the case of \( n = 3 \). In this case the SUSY breaking scale is rather high: \( O(10^8) \) GeV–\( O(10^{10}) \) GeV, if the upper bound is saturated. See Eq. [9] and Fig. 3. The production of the SUSY particles including the gravitino is suppressed, if the reheating temperature is (much)
Figure 8: The contours of the SM-like Higgs boson mass in the plane of $\tilde{m}$ and $\tan \beta$, corresponding to the cases of $n = 2$ and $n = 3$ for which the Hubble parameter is about $10^6$ GeV and $10^{10}$ GeV, respectively.

lower than the SUSY breaking scale. According to Fig. 7, this is the case for $M_1 \ll 10^{12}$ GeV. Thus there is no cosmological problem associated with the SUSY particles. A plausible candidate for DM will be the QCD axion, although it may be possible that the incomplete thermalization of the lightest supersymmetric particle (LSP) accounts for the DM abundance.

In the following we discuss the cosmology and phenomenology, focusing on the case of $n = 2$, unless otherwise stated. Some of the discussion below can be straightforwardly applied to the case of $n \geq 3$.

6.1 Leptogenesis

In the present model, the right-handed neutrinos are non-thermally produced by the inflaton decay. Let us see if the decay of right-handed neutrinos can yield the right amount of the baryon asymmetry, $n_B/s \sim 8 \times 10^{-11}$ [13]. The abundance of the lightest right-handed neutrino is given by

$$\frac{n_{N_1}}{s} = \frac{3}{2} \frac{T_R}{m_\sigma}$$

(38)
Assuming that $N_1$ immediately decays after produced by the inflaton decay, the lepton number generated by the $N_1$ decay is [11]

$$
\frac{n_L}{s} \simeq 3 \times 10^{-10} \left( \frac{T_R}{10^6 \text{ GeV}} \right) \left( \frac{M_1}{m_\sigma} \right) \left( \frac{0.05 \text{ eV}}{m_{\nu_3}} \right) \delta_{\text{eff}},
$$

(39)

where $m_{\nu_3}$ denotes the mass of the heaviest left-handed neutrino and $\delta_{\text{eff}}$ the effective CP phase. The lepton asymmetry is related to the baryon asymmetry as $n_B/s = -(8/23)n_L/s$ through the sphaleron process. We find that the correct amount of baryon asymmetry is (marginally) generated for $M_1/m_\sigma \sim 0.4$, $T_R \simeq 2 \times 10^6 \text{ GeV}$ and $|\delta_{\text{eff}}| \simeq 1$ in the case of $n = 2$. It is possible to enhance the baryon asymmetry in several ways. So far we have set $g = 1$ for simplicity. If $g \approx 0.1$, for instance, the reheating temperature can be increased by a factor 2 as long as $M_1 \sim m_\sigma$. Alternatively, if the right-handed neutrinos are degenerate, the lepton asymmetry can be enhanced [26]. The (non-)thermal leptogenesis by the $U(1)_{B-L}$ Higgs boson decay has been recently studied in detail in Ref. [27], where the parameters are motivated by the hybrid inflation [28].

Note that the above argument assumes that there is no additional entropy production. Later we will show that this is indeed the case even in the presence of the Polonyi field. For $n \geq 3$, the reheating temperature can be higher, and the right amount of the baryon asymmetry can be produced for a broader parameter range.

### 6.2 Gravitino problem

The gravitinos are produced both thermally and non-thermally at the reheating, and its abundance is tightly constrained by cosmology. For $m_{3/2} \lesssim 30 \text{ TeV}$, the lifetime is shorter than about 1 sec and the energetic particles produced by the gravitino decay changes the helium-4 abundance through affecting the proton-neutron conversion process [29]. For $m_{3/2} \gtrsim 30 \text{ TeV}$, on the other hand, the lifetime is so short that it decays before BBN, and there is no constraint coming from BBN. Instead, the LSPs produced by the gravitino decay contribute to the DM density if the R-parity is conserved. These constraints are summarized as

$$
Y_{3/2} \equiv \frac{n_{3/2}}{s} \lesssim \begin{cases} 
5 \times 10^{-13} & \text{for } 10 \text{ TeV} \lesssim m_{3/2} \lesssim 30 \text{ TeV}, \\
4 \times 10^{-13} \left( \frac{1 \text{ TeV}}{m_{\text{LSP}}} \right) & \text{for } m_{3/2} \gtrsim 30 \text{ TeV},
\end{cases}
$$

(40)
where $m_{\text{LSP}}$ denotes the LSP mass. Notice that the second constraint assumes the R-parity. If the R-parity is violated by a small amount, the LSP can decay before BBN, and there will be no cosmological constraint on the gravitino abundance for $m_{3/2} \gtrsim 30\,\text{TeV}$.

The LSP mass depends on the SUSY breaking mediation. In the gravity mediation we expect that the gravitino mass is comparable to the sfermion and gaugino masses, collectively denoted by $\tilde{m}$. If the upper bound on $\tilde{m}$ (see (9)) is saturated, we expect $m_{\text{LSP}} = O(100)\,\text{TeV} - O(1)\,\text{PeV}$. Suppose that the LSP is the lightest neutralino. In this case the thermal relic abundance exceeds the observed DM density, and either late-time entropy production or the R-parity breaking is needed. We note however that, if the LSP mass is higher than the reheating temperature, the LSP overproduction may be avoided. If the gravitino is the LSP of mass $O(100)\,\text{TeV}$, the bound (40) should read with $m_{\text{LSP}} = m_{3/2}$. The gravitino is mainly produced by the decay of the next-lightest supersymmetric particles, and the gravitino abundance likely exceeds the DM density. This problem can be avoided again by either late-time entropy production or the R-parity violation.

In the anomaly mediation with a generic Kähler potential, the gravitino mass is comparable to the sfermion mass $\tilde{m} = O(100)\,\text{TeV} - O(1)\,\text{PeV}$, while the gaugino mass is suppressed (see footnote 3), and we expect $m_{\text{LSP}} = O(100)\,\text{GeV} - O(1)\,\text{TeV}$. In the case of the Higgsino or Wino-like LSP, its thermal relic abundance can be smaller than the observed one.

In the following we consider thermal and non-thermal production of the gravitinos separately and show that in both cases the gravitino abundance satisfies the cosmological bound (40).

### 6.2.1 Thermal production

Gravitinos are produced by scatterings of particles in thermal bath during the reheating process. The abundance is estimated to be [30, 31, 32]

$$Y_{3/2}^{(TP)} \simeq 2 \times 10^{-16} \left(1 + \frac{m_{\tilde{g}}^2}{3m_{3/2}^2} \right) \left( \frac{T_R}{10^6\,\text{GeV}} \right),$$

where $m_{\tilde{g}}$ denotes the gluino mass and we have omitted the logarithmic dependence on $T_R$ as well as terms that depend on the other gaugino masses. Note that the definition of
$T_R$ is given by (37).

Let us consider the case of the gravity mediation, in which the gluino as well as the LSP have a mass comparable to $m_{3/2}$ of $O(100)\,\text{TeV}$ to $O(1)\,\text{PeV}$. The bound (40) is marginally satisfied for $m_{\text{LSP}} = 1\,\text{PeV}$ and $T_R = 10^6\,\text{GeV}$. In the anomaly mediation, the bound is relaxed because of the suppressed gaugino masses. Note that the bound disappears if the R-parity is broken.

### 6.2.2 Non-thermal production

Gravitinos are generically produced non-thermally by the inflaton decay [21, 22, 23]. The gravitino production rate depends on the SUSY breaking mechanism. Let us first consider the gravity mediation. In the simple Polonyi model, there is a singlet SUSY breaking field $z$ of mass $m_z \sim m_{3/2}$. The inflaton decays into a pair of gravitinos through the following interaction in the Kähler potential

$$K = \frac{1}{2}(c_\Phi|\Phi|^2 + c_{\bar{\Phi}}|\bar{\Phi}|^2)zz + \text{h.c.},$$

where $\phi$ denotes an inflaton field. The gravitino production rate is

$$\Gamma_{3/2} = \frac{1}{64\pi} \tilde{c}^2 v_{B-L}^2 m_{\sigma}^3,$$

where $\tilde{c} \equiv (c_\Phi + c_{\bar{\Phi}})/2$. The resultant gravitino abundance is

$$Y_{3/2}^{(\text{NTP})} \sim 2 \times 10^{-14} \tilde{c}^2 \left(\frac{T_R}{10^6\,\text{GeV}}\right)^{-1} \left(\frac{v_{B-L}}{3 \times 10^{15}\,\text{GeV}}\right)^2 \left(\frac{m_{\sigma}}{5 \times 10^9\,\text{GeV}}\right)^2,$$

where we have set $g_* = 200$. The bound (40) can be satisfied for the LSP mass of 100 TeV and $\tilde{c} \lessapprox 0.3$. In the dynamical SUSY breaking, the $z$ can have a mass much heavier than $m_{3/2}$. In this case the gravitino production rate is similar to (44). If the $z$ is not an elementary field but a composite one at the scale of the inflaton mass, the gravitino production rate can be suppressed by a factor of $O(10^2)$ or so [23]. In this case the bound (40) can be satisfied. As mentioned before, however, the thermal relic abundance of the LSP is generically too large in this case. Once we introduce the R-parity violation to avoid the LSP overproduction, there is no bound on the gravitino abundance.

In the anomaly mediation, no singlet SUSY breaking field is necessary, and the gravitino production rate is suppressed by a factor of $O(10^2)$ compared to (44) [23]. In addition,
the LSP mass is suppressed compared to the case of gravity mediation. Therefore the bound (10) can be satisfied without introduction of the R-parity violation, if the thermal relic abundance of the Higgsino or Wino LSP is sufficiently small.

6.3 Dark Matter

Here we discuss DM candidates in our model. Among various possibilities, we consider the neutralino LSP and the QCD axion. Especially in the presence of the R-parity violation, the latter will be a plausible DM candidate, and we will study its cosmological constraints in detail.

6.3.1 Neutralino DM

In the gravity mediation, the LSP has a mass of $O(100)\,\text{TeV}$ or so, and its thermal relic abundance exceeds the observed DM abundance. If the reheating temperature is much lower than $O(100)\,\text{TeV}$, the LSP abundance can be suppressed, which however makes it difficult for leptogenesis to work. The simplest solution to the overabundance of the neutralino LSPs is to break R-parity by a small amount. Then the LSP is no longer stable, and it can decay before BBN. Of course the LSP cannot be DM in this case, and we need another DM candidate.

In the anomaly mediation, the neutralino LSP is expected to be as light as $O(100)\,\text{GeV}$-$O(1)\,\text{TeV}$, while the sfermion masses are much heavier. In this case the neutralino LSP with a sizable Higgsino or Wino fraction can account for the present DM abundance. In fact, it is well-known that thermal relic of the Wino LSP with mass of $2.7\,\text{TeV}$ can account for the DM [34], while the gravitino and scalar fermions lie at $O(1)\,\text{PeV}$. It is intriguing that the PeV-scale SUSY inferred from the $U(1)_{B-L}$ new inflation is compatible with the Wino DM [9].

6.3.2 Axion

Here we consider the axion cosmology. The axion is a pseudo Nambu-Goldstone boson in association with the spontaneous breakdown of a global $U(1)_{\text{PQ}}$ symmetry, so called the Peccei-Quinn (PQ) symmetry [35, 36]. The PQ mechanism is known as the most plausible solution to the strong CP problem in QCD. There are several ways to implement the PQ
mechanism. For simplicity we assume there is another sector in which the PQ symmetry is spontaneously broken. The breaking scale of the $U(1)_{PQ}$ symmetry is bounded below by the axion emission from red giant stars, and as a result, the axion mass is extremely light. Thus the axion is stable in a cosmological time scale, so the candidate for DM.

In the early Universe the axion gets coherently excited, and its abundance is given by

$$\Omega_a h^2 \simeq 0.2 \left( \frac{f_a}{10^{12} \text{ GeV}} \right)^{1.18} \theta^2,$$

where $f_a$ denotes the PQ symmetry breaking scale and $\theta$ the initial misalignment angle of the axion. Thus it accounts for the present DM abundance for $f_a \sim 10^{11} - 10^{12}$ GeV without tuning on the angle $\theta$. If we allow the fine-tuning of the misalignment $\theta \lesssim O(10^{-3})$, $f_a$ can be increased up to the GUT scale.

Notice that the PQ scale is higher than the inflation scale and the reheating temperature. Thus the PQ symmetry is likely broken already during inflation, it is not restored after that. In this case the axion obtains quantum fluctuations during inflation and it contributes to the CDM isocurvature perturbation, which is constrained by the observation of the CMB anisotropy. Assuming that the axion is a dominant component of DM, the magnitude of the CDM isocurvature perturbation is estimated as

$$S_c = \frac{2\delta \theta}{\theta} = \frac{H_{\text{inf}}}{\pi f_a \theta} \sim 3 \times 10^{-7} \left( \frac{H_{\text{inf}}}{10^6 \text{ GeV}} \right) \left( \frac{f_a \theta}{10^{12} \text{ GeV}} \right)^{-1}.$$  

This satisfies the observational constraint from WMAP+BAO+H$_0$:

$$|S_c| \lesssim 1.4 \times 10^{-5} \ (95\% \ C.L.)$$

The upper bound on $|S_c|$ will be improved by a factor 2 or so by the Planck satellite alone.

In the case of $n \geq 3$, the Hubble parameter is greater than $10^{10}$ GeV. Therefore the axion isocurvature perturbation excludes the axion as a DM candidate as long as the PQ symmetry is broken during and after inflation.

Finally we comment on cosmology of the supersymmetric partners of the axion, saxion ($s$) and axino ($\tilde{a}$). The saxion generically obtains a mass of order the gravitino mass, and it decays into the axion pair or the SSM particles such as gluons, Higgs boson,

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7 In fact this depends on the stabilization mechanism of the saxion.
and SM fermions, depending on the detailed model structure. The saxion is generated in a form of coherent oscillations and its abundance is given by

$$\frac{\rho_s}{s} = \frac{1}{8} T_R \left( \frac{s_i}{M_p} \right)^2 \simeq 2 \times 10^{-8} \text{GeV} \left( \frac{T_R}{10^6 \text{GeV}} \right) \left( \frac{f_a}{10^{12} \text{GeV}} \right)^2 \left( \frac{s_i}{f_a} \right)^2, \quad (48)$$

where $s_i$ is the initial amplitude of the saxion. If the saxion mainly decays into a pair of axions, its lifetime is

$$\tau_s = \left( \frac{1}{64\pi} \frac{m_s^3}{f_a^2} \right)^{-1} \simeq 1 \times 10^{-13} \text{sec} \left( \frac{m_s}{100 \text{ TeV}} \right)^{-3} \left( \frac{f_a}{10^{12} \text{ GeV}} \right)^2, \quad (49)$$

where $m_s$ denotes the saxion mass. Thus the saxion decays before it dominates the Universe for the parameters shown in the parentheses. Even for $s_i \sim f_a \sim 10^{16} \text{GeV}$, the saxion does not dominate if $m_s \sim 1 \text{ PeV}$.

The axino is produced thermally during the reheating \[41\]. Its abundance is given by\[8\]

$$Y_{\tilde{a}} \equiv \frac{n_{\tilde{a}}}{s} \simeq 2 \times 10^{-7} g_s^6 \ln \left( \frac{1.108}{g_s} \right) \left( \frac{f_a}{10^{12} \text{ GeV}} \right)^{-2} \left( \frac{T_R}{10^6 \text{ GeV}} \right), \quad (50)$$

where $g_s$ is the strong coupling constant. The axino mass depends on how the saxion is stabilized. Let us assume that the axino mass is comparable to $\tilde{m}$ and that the axino is unstable, because otherwise the axino density will easily exceed the DM abundance. Then the axino lifetime is given by \[47\]

$$\tau_{\tilde{a}} = \left( \frac{\alpha_s^2}{16\pi^3} \frac{m_{\tilde{a}}^3}{f_{\tilde{a}}^2} \right)^{-1} \simeq 3 \times 10^{-11} \text{sec} \left( \frac{m_{\tilde{a}}}{100 \text{ TeV}} \right)^{-3} \left( \frac{f_a}{10^{12} \text{ GeV}} \right)^2, \quad (51)$$

where it is assumed that the axino mainly decays into the gluino and gluon. Hence, for $m_{\tilde{a}} \gtrsim O(10) \text{ TeV}$, the axino also decays before it dominates the Universe. There are no late-time entropy production processes from these additional particles.

### 6.4 Polonyi problem

Now we consider the cosmology of the SUSY breaking sector. In the gravity mediation, there is a singlet SUSY breaking field $z$, so called the Polonyi field. The Polonyi field causes

\[8\]In the DFSZ axion model \[42, 43\], the axino abundance is saturated for a high reheating temperature \[44\], and its expression is given in Ref. \[45\] taking also account of the Higgsino decay \[46\].
a cosmological problem as we shall briefly explain below. In the anomaly mediation, on
the other hand, such a singlet field is not necessary, therefore there is no Polonyi problem.

First, let us review the Polonyi problem in the original sense [48]. We assume that
the Polonyi has only Planck suppressed interactions, and its potential is approximated
by a quadratic potential up to the Planck scale. The Polonyi begins to oscillate at
\( H \sim m_z \) and its abundance is given by
\[
\frac{\rho_s}{s} = \frac{1}{8} T_R \left( \frac{z_i}{M_p} \right)^2 \simeq 1 \times 10^5 \text{GeV} \left( \frac{T_R}{10^6 \text{GeV}} \right) \left( \frac{z_i}{M_p} \right)^2,
\]
where \( z_i \) is the initial amplitude of the Polonyi. Although the Polonyi decays before BBN
for \( m_z \gtrsim O(10) \text{ TeV} \), it releases a huge amount of entropy because the Polonyi dominates
the Universe soon after the reheating. Therefore any pre-existing baryon asymmetry is
diluted, in particular, the leptogenesis scenario does not work. In addition, the LSPs
produced by the Polonyi decay may overclose the Universe. One of the attractive solutions
to the Polonyi problem is to introduce an enhanced coupling between \( \chi \) and the
Polonyi field [49, 50, 51]. However, since the Hubble parameter during inflation is close
to the Polonyi mass, it is not easy to completely solve the Polonyi problem by this mech-
anism [51].

Second, we consider the case where the F-term of the Polonyi has a dynamical ori-
gin [52]. Note that the Polonyi field itself must be an elementary singlet to give a sizable
mass to gauginos. In this set-up, the Polonyi has a larger SUSY breaking mass at the
potential minimum, which relaxes the Polonyi problem mentioned above. However, as
noted in Ref. [53], the Polonyi field may be driven to the Planck scale, because the po-
tential becomes flat at scales beyond the dynamical scale \( \Lambda \). To see this, let us write the
potential as
\[
V(z) \simeq \begin{cases} 
m_2 |z|^2 & \text{for } |z| < \Lambda \\
3m_{3/2}^2 M_p^2 & \text{for } |z| > \Lambda,
\end{cases}
\]
where \( m_z \sim \sqrt{3m_{3/2} M_p/(4\pi)^2} \) is the Polonyi mass around the origin and the dynamical
scale \( \Lambda \) and the gravitino mass is related by \((\Lambda/4\pi)^2 \sim 3m_{3/2}^2 M_p^2\). We have set coupling
constants of \( z \) to be order unity, for simplicity. In general there exists a linear term in \( z \)
during inflation and it may destabilize the Polonyi field if the inflation scale is too high.
Let us consider the Kähler potential $K = cz + c^* z^*$ with constant $c$ of order $M_p$. In the supergravity, it yields the following term in the scalar potential during inflation

$$V_{\text{lin}}(z) \simeq (cz + c^* z^*) \frac{V_{\text{inf}}}{M_p^2} = 3H_{\text{inf}}^2 (cz + c^* z^*).$$

In order for this linear term not to destabilize the Polonyi field, we need $V_{\text{lin}}(\Lambda) < 3m_{3/2}^2 M_p^2$. This condition is written as

$$H_{\text{inf}} \lesssim 2 \times 10^8 \text{ GeV} \left( \frac{m_{3/2}}{10^3 \text{ TeV}} \right)^{3/4} \left( \frac{M_p}{|c|} \right)^{1/2}.$$  \hspace{1cm} (55)

This is satisfied for $n = 2$, but not for $n \geq 3$. See Fig. 3. The Polonyi problem is absent in the case of $n = 2$ in the the dynamical SUSY breaking scenario. Therefore, as long as we consider the gravity mediation, the case of $n = 2$ is favored.

### 6.5 Moduli stabilization and the dynamical origin of the inflation scale

It has been known that the inflation scale should be smaller than the gravitino mass,

$$H_{\text{inf}} \lesssim m_{3/2},$$

in order not to destabilize the moduli in the simple class of the modulus stabilization models [55, 14]. In our scenario, the SUSY breaking is bounded above by the inflation, and so, it is interesting to see if the inequality (56) can be satisfied.

We estimated the effect of the constant term in the superpotential on the inflaton dynamics in Ref. [9], assuming $|k_5|$ is of order unity, and concluded that $m_{3/2}$ should be one order of magnitude smaller than the Hubble parameter during inflation. In fact, this upper bound can be relaxed to be consistent with (56) if $|k_5| \gg 1$ and $k_5 < 0$. Then the shift of $\chi$ due to the constant term becomes much smaller than the Planck scale, and the analysis so far can be applied to the case of $m_{3/2} \gtrsim H_{\text{inf}}$. Considering that $\tilde{m}$, which is considered to be comparable to the gravitino mass, cannot exceed $H_{\text{inf}}$ (see (56)), the inequality (56) can be marginally satisfied in our set-up, namely,

$$H_{\text{inf}} \sim m_{3/2}.$$  \hspace{1cm} (57)
The reason why the successful inflation is possible even when (57) is satisfied is that the flatness of the inflaton potential is ensured by the $Z_n$ discrete symmetry in our model (11). This should be contrasted to other inflation models such as the hybrid inflation [28] and the single-field new inflation [18] in which the inflaton is charged under a continuous or discrete R-symmetry, and therefore the inflaton potential necessarily receives a correction linear in the inflaton field once the constant term which breaks the R-symmetry is included [54].

As mentioned before, the enhanced coupling of $\chi$ has been considered in context of the adiabatic solution to the moduli problem [50], and it can suppress the modulus abundance so that there will be no significant entropy dilution by the modulus decay [51]. This is especially the case if the modulus has a SUSY mass much heavier than $m_{3/2}$ as in the KKLT model [55].

Interestingly, such an enhancement naturally arises if the inflationary scale, $v$, in Eq. (11) has a dynamical origin. It is straightforward to apply the IYIT model [52] to generate the F-term of $\chi$. Then there is generically a coupling like $K \supset -|\chi|^4/\Lambda_I^2$, where $\Lambda_I$ is the dynamical scale. Note that since the inflaton $\Phi$ and $\bar{\Phi}$ do not participate in the strong dynamics, there is no large contribution to the inflaton mass. The dynamical scale $\Lambda_I$ is intriguingly close to $\Lambda$ for SUSY breaking, and so, both may be related to each other.

Thus, a slight enhancement of the coupling of $\chi$, or equivalently lowering the cut-off scale of the $\chi$’s interaction may be the key to establish a successful moduli cosmology.

7 Conclusions

In this paper we have studied the dynamics of the recently proposed new inflation in detail, where the inflaton is the Higgs field responsible for the breaking of $U(1)_{B-L}$ symmetry. Importantly, we have shown that the soft SUSY breaking is bounded above for the successful inflation. This is because otherwise the CW potential would make the inflaton potential too steep. Interestingly, in the case of $n = 2$, the inflaton VEV, which determines the $U(1)_{B-L}$ breaking scale, is intriguingly close to the see-saw scale of order $10^{15}$ GeV. The upper bound on the SUSY breaking is then about $O(100)$ TeV to $O(1)$ PeV. We have also found that the residual CW correction can increase the predicted spectral index in
consistent with the WMAP data. (Note that the spectral index will be $0.94 - 0.95$ without the CW correction, which causes a tension at $2\sigma$ level.) Furthermore we have discussed various implications of our model: the SM-like Higgs boson mass at about 125 GeV can be easily explained; non-thermal leptogenesis works successfully; thermal and non-thermal gravitino problem can be avoided; the DM candidates are either the lightest neutralino (e.g. the Wino of mass 2.7 TeV) or the QCD axion; the Polonyi/moduli problem can be solved; the constraint on the inflation scale from the modulus stabilization can be marginally satisfied. Thus, our inflation model based on the minimal B-L extension of SSM has surprisingly many positive implications in cosmology and phenomenology.

**Acknowledgment**

This work was supported by the Grant-in-Aid for Scientific Research on Innovative Areas (No. 21111006) [KN and FT], Scientific Research (A) (No. 22244030 [KN and FT] and No.21244033 [FT]), and JSPS Grant-in-Aid for Young Scientists (B) (No. 21740160) [FT]. This work was also supported by World Premier International Center Initiative (WPI Program), MEXT, Japan.

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