Lower Bound on the Gravitino Mass

\[ m_{3/2} > O(100) \text{ TeV} \]

in \( R \)-Symmetry Breaking New Inflation

Keisuke Harigaya\(^1\), Masahiro Ibe\(^2,1\) and Tsutomu T. Yanagida\(^1\)

\(^1\)Kavli IPMU (WPI), TODIAS, University of Tokyo, Kashiwa 277-8583, Japan
\(^2\)ICRR, University of Tokyo, Kashiwa 277-8582, Japan

Abstract

In supersymmetric theories, the \( R \) symmetry plays a unique role in suppressing a constant term in the superpotential. In single chiral field models of spontaneous breaking of a discrete \( R \) symmetry, an \( R \)-breaking field can be a good candidate for an inflaton in new inflation models. In this paper, we revisit the compatibility of the single-field \( R \)-breaking new inflation model with the results of the Planck experiment. As a result, we find that the model predicts a lower limit on the gravitino mass, \( m_{3/2} > O(100) \text{ TeV} \). This lower limit is consistent with the observed Higgs mass of 126 GeV when the masses of the stops are of order the gravitino mass scale.
1 Introduction

In supersymmetric (SUSY) theories, the $R$ symmetry plays a unique role in suppressing a constant term in the superpotential. Without the $R$ symmetry, the constant term is expected to be at the Planck scale, which requires a SUSY breaking scale to be the Planck scale to achieve the almost flat universe. Thus, there is a strong case for the existence of a spontaneously broken $R$-symmetry if SUSY is the solution to the hierarchy problem \cite{1-4} between the weak scale and the Planck scale or the scale of the Grand Unified Theory (GUT).

One caveat of the $R$ symmetry is that a generation of the appropriate vacuum expectation value (VEV) of the superpotential requires a symmetry breaking field to have a Planck scale $A$-term VEV and a non-vanishing $F$ term VEV at the same time if the symmetry is a continuous one \cite{5}. This means that an $R$ symmetry breaking field is nothing but the Polonyi field for the continuous $R$ symmetry. Therefore, by taking the Polonyi problem \cite{6} seriously, the $R$ symmetry which suppresses the constant term of the superpotential should be a discrete one.

Interestingly, the simplest model of spontaneous discrete $R$ symmetry breaking consisting of a single chiral field has a convex but a very flat potential around the origin of the chiral field\cite{1}, which evokes a scalar potential used in new inflation models \cite{7,8}. In fact, the simplest $R$-breaking model satisfies the slow-roll conditions in a wide parameter region, and hence, the $R$-breaking field is a good candidate for an inflaton \cite{9-14}. It is also remarkable that the domain wall problem \cite{15} associated with the discrete $R$ symmetry breaking is automatically solved when the $R$ symmetry breaking field plays a role of the inflaton\footnote{Ref. \cite{9} pointed out not only the presence of the so-called $\eta$ problem in supergravity inflation models but also the importance of the $R$ symmetry to have flat potentials necessary for the inflation to occur.}

We here emphasize that new inflation models tend to predict a small tensor fraction due to their small inflation scales \cite{16}. This property is fairly supported by the upper limit on the tensor fraction of cosmic perturbations set by the recent observations of the cosmic microwave background (CMB) \cite{17,19} \footnote{This situation is analogous to the original new inflation model \cite{7,8}, where an inflaton is identified with a GUT breaking field and the monopole problem is solved.} \footnote{Simple large field inflation models such as the chaotic inflation models with a quadratic or a quartic potential lead to a small tensor fraction.}

\begin{thebibliography}{99}
\bibitem{1} Ref. \cite{9} pointed out not only the presence of the so-called $\eta$ problem in supergravity inflation models but also the importance of the $R$ symmetry to have flat potentials necessary for the inflation to occur.
\bibitem{2} This situation is analogous to the original new inflation model \cite{7,8}, where an inflaton is identified with a GUT breaking field and the monopole problem is solved.
\bibitem{3} Simple large field inflation models such as the chaotic inflation models with a quadratic or a quartic potential lead to a small tensor fraction.
\end{thebibliography}
In this paper, we further investigate the compatibility of the $R$-breaking new inflation model with the results of the Planck experiment [18, 19]. As we will see, the $R$-breaking new inflation model is consistent with all cosmological constraints and observations in a wide parameter region. Furthermore, the model predicts a lower bound on the gravitino mass, $m_{3/2} > O(100)\,\text{TeV}$. This lower limit on the gravitino mass is consistent with the observed Higgs mass of 126 GeV [26, 27] in a class of models in which the masses of the stops are of order the gravitino mass [28–30]. We also show that the baryon asymmetry of the universe as well as the observed dark matter density can be consistently explained along with the $R$-breaking new inflation model.

2 Brief review on the $R$-breaking new inflation model

Let us begin with the simplest model of spontaneous discrete $Z_{NR}$ symmetry breaking consisting of a single chiral field $\phi$ [9, 10]. Here, we assume that $\phi$ is a singlet except for the $R$ symmetry with an $R$ charge 2. Assuming $N = 2n$, the superpotential of $\phi$ is given by,

$$W = v^2 \phi - \frac{g}{n+1} \phi^{n+1} + \cdots,$$

where the ellipses represent higher power terms of $\phi$. We neglect them throughout this paper, since we are interested in the region with $|\phi| \ll 1$. The size of the coupling constant $g$ will be discussed later. Here and hereafter, we take the unit of the reduced Planck scale $M_{PL} \simeq 2.4 \times 10^{18} \,\text{GeV}$ being unity. The parameters $v^2$ and $g$ are taken real and positive without loss of generality. At supersymmetric vacua, the $Z_{2nR}$ symmetry is spontaneously broken down to the $Z_{2R}$ symmetry by the VEV of $\phi$,

$$\langle \phi \rangle \simeq \left( \frac{v^2}{g} \right)^{1/n} \times e^{2\pi i m/n}, \quad (m = 0, 1, \ldots, n-1)$$

potential [20] are, on the other hand, now slightly disfavored at least by 1\sigma level, which requires some extensions [21–25].

In order for the gravitino mass to be far smaller than the Planck scale, $v^2$ must be suppressed. The suppression can be explained, for example, by assuming an $U(1)_R$ symmetry under which $\phi$ has a charge of $2/(n+1)$, and the $U(1)_R$ symmetry be dynamically broken by a condensation of a (composite) chiral field with an $U(1)_R$ charge of $2 - 2(n+1)$ [10].
which leads to the VEV of the superpotential,

$$
\langle W \rangle \simeq \frac{n}{n+1} v^2 \left( \frac{v^2}{g} \right)^{1/n} e^{2\pi m/n}.
$$

(3)

As we emphasized in the introduction, the scalar potential of this model is convex but very flat around $\phi \sim 0$. Thus, if the initial field value of $\phi$ is set close to its origin by, for example, a positive Hubble induced mass term of pre-inflation [11] and the slow-roll conditions are satisfied, $\phi$ automatically brings about the inflation. Therefore, the simplest model of discrete $R$-symmetry breaking is equipped with necessary structures as a model of new inflation.

Now, let us discuss details of the new inflation model. For that purpose, let us note that the Kähler potential of $\phi$ is given by

$$
K = \phi \phi^\dagger + \frac{1}{4} k (\phi \phi^\dagger)^2 + \cdots,
$$

(4)

where the ellipses denote higher power terms of $\phi$, whose contributions to the dynamics of $\phi$ are negligible again. The parameter $k$ is at most of order unity, and we assume $k > 0$ so that $\phi = 0$ is a local maximum (see below). From Eqs. (2) and (4), the scalar potential of the scalar component of $\phi$ is given by

$$
V(\phi) = |v^2 - g\phi^n|^2 - k v^4 |\phi|^2 + \cdots = v^4 - (gv^2\phi^n + \text{h.c.}) - kv^4 |\phi|^2 + \cdots.
$$

(5)

In terms of the radial and the angular components of $\phi$, $\phi = \varphi e^{i\theta}/\sqrt{2}$, the scalar potential is rewritten as,

$$
V(\varphi, \theta) = v^4 - \frac{k}{2} v^4 \varphi^2 - \frac{g}{2^{n/2-1}} v^2 \varphi^n \cos(n\theta) + \cdots.
$$

(6)

It can be seen that for a given $\varphi > 0$, the minimum of the potential is provided by $\theta = 2\pi l/n$ ($l = 0, 1, \cdots, n-1$). In the following, the radial component $\varphi$ plays a role of the inflaton in new inflation.

As we have mentioned, we assume that the initial condition of $\varphi$ is close to 0, i.e. $|\varphi| \ll 1$. We further suppose that the initial condition of the angular direction $\theta$ is given by $\theta = 0 \mod 2\pi/n$ for the time being. Since $\theta = 0 \mod 2\pi/n$ is the minimum of
the potential along the angular direction, $\theta = 0 \mod 2\pi/n$ is kept during the inflation. Along the inflaton trajectory, the first and the second slow-roll parameters are given by

$$
\epsilon \equiv \frac{1}{2} \left( \frac{\partial V}{\partial \varphi} \right)^2 = \frac{1}{2} \left( k \varphi + \frac{ng}{2^{n/2-1}} \frac{\varphi^{n-1}}{v^2} \right)^2, \\
\eta \equiv \frac{\partial^2 V}{\partial \varphi^2} = -k - \frac{n(n-1)g}{2^{n/2-1}} \frac{\varphi^{n-2}}{v^2}.
$$

(7)

Thus, the slow-roll conditions can be actually satisfied for $|\varphi| \ll 1$ as long as $k \ll 1$.

By assuming $|k| \ll 1$, the inflation lasts until the inflaton reaches to

$$
\varphi_{\text{end}} = \left( \frac{2^{(n-2)/2} v^2}{n(n-1)g} \right)^{1/(n-2)},
$$

(8)

at which the slow-roll conditions are violated, $|\eta| \simeq 1$. It should be noted that there is an one-to-one correspondence between the number of $e$-foldings $N_e$ and the field value of $\varphi$ during the inflation via

$$
N_e(\varphi) = \int_{\varphi_{\text{end}}}^{\varphi} \frac{V}{\partial V/\partial \varphi} d\varphi.
$$

(9)

Thus, by taking the inverse of Eq. (9), we obtain

$$
\varphi^{n-2}(N_e) = \frac{2^{(n-2)/2} k v^2}{n g} \left( e^{k(n-2)N_e} - 1 + k(n-1)e^{k(n-2)N_e} \right)^{-1}.
$$

(10)

In order to compare model predictions with CMB observations, let us calculate the properties of the curvature perturbation. The spectrum of the curvature perturbation $P_\zeta$ and its spectral index $n_s$ are given by

$$
P_\zeta = \frac{1}{24\pi^2 \epsilon} = \frac{1}{24\pi^2} \left( n^2 g^2 k^{-2(n-1)} v^{4(n-3)} \left( e^{k(n-2)N_e} - 1 \right)^{2(n-1)} \right)^{\frac{1}{2}} e^{-2k(n-2)N_e},
$$

(11)

$$
n_s = 1 - 6\epsilon + 2\eta = 1 - 2k \left( 1 + \frac{n-1}{(1+k(n-1))e^{k(n-2)N_e} - 1} \right),
$$

(12)

respectively. In Fig. 1, we show the prediction on the spectral index for $n = 4, 5, 6$ and $N_e = 50$. The colored region shows a region favored by the Planck experiment, i.e. $n_s = 0.9643 \pm 0.012$ [19] for the pivot scale $k_\star = 0.002$ Mpc$^{-1}$ at 95\%C.L. It can be seen that the model with $n \leq 4$ is disfavored by the Planck experiment for $N_e = 50$. For $n = 5$, $k \sim 10^{-2}$ is favored. In Fig. 2, we show the $N_e$ dependence of the spectral index.
Figure 1: The spectral index of the curvature perturbation $n_s$ for $n = 4, 5, 6$ with $N_e = 50$. A colored region show the 95% C.L. favored region by the Planck experiment, $n_s = 0.9643 \pm 0.012$ \cite{19}. For $n = 4$. The figure shows that the model with $n = 4$ is still consistent with the Planck experiment for $N_e \gtrsim 56$\footnote{In Ref. \cite{14}, it is pointed out that the model with $n = 4$ is also consistent with the Planck experiment if there are a small constant term in the superpotential beside the one from the condensation of $\phi$. Since we assume that the $R$ symmetry is broken only by the condensation of $\phi$, that solution is not applicable.}. We will discuss impacts of the observed spectral index to the gravitino mass in the next section.

Before closing this section, let us discuss more general initial conditions for the inflaton field, $\theta \neq 0 \pmod{2\pi/n}$. In particular, we are interested in how the spectral index is affected, since $n = 4$ is severely constrained for $\theta = 0 \pmod{2\pi/n}$ by the Planck results. In Fig. 3, we show a schematic picture of the shape of the inflaton potential for $n = 4$. For a better presentation, we show only the region with $\text{Re}(\phi) > 0$. For a fixed number of e-foldings, a non-zero angle $\theta$ leads to a larger corresponding field value for $\phi$. As a result, the curvature of the inflaton trajectory becomes negatively larger, and the spectral index becomes more red-tilted. Therefore, even if we consider the initial condition with $\theta \neq 0 \pmod{2\pi/n}$, the model with $n = 4$ is still disfavored unless $N_e$ is large.
Figure 2: The spectral index of the curvature perturbation $n_s$ for $n = 4$ with various $N_e$. The colored region show the 95% C.L. limit from the Planck experiment, $n_s = 0.9643 \pm 0.012$ [19].

3 Lower bound on the gravitino mass

In this section, we put a lower bound on the gravitino mass $m_{3/2}$ in the $R$-breaking new inflation models based on the results obtained in the previous section. From Eq. (11), the parameter $v^2$ is expressed by the curvature perturbation, $\mathcal{P}_\zeta \simeq 2.2 \times 10^{-9}$ [19], as

$$v^2 = \left(24\pi^2 \mathcal{P}_\zeta \right)^{n-2} (ng)^{-2} \left( \frac{k}{e^{k(n-2)N_e} - 1} \right)^{2(n-1)} e^{2k(n-2)N_e} \frac{1}{\pi(n-1)},$$

which leads to

$$v \simeq \begin{cases} 
9.0 \times 10^{11} \text{ GeV} \ g^{-1/2} & (n = 4, k = 0.01, N_e = 56), \\
6.2 \times 10^{13} \text{ GeV} \ g^{-1/4} & (n = 5, k = 0.01, N_e = 50), \\
2.5 \times 10^{14} \text{ GeV} \ g^{-1/6} & (n = 6, k = 0.01, N_e = 50).
\end{cases}$$

(14)

It should be noted that $v$ does not depend on $k$ significantly. As a result, the gravitino mass $m_{3/2}$ is given by

$$m_{3/2} = \frac{ng}{n + 1} \left( \frac{v^2}{g} \right)^{\frac{n+1}{n}}$$

$$\simeq \begin{cases} 
1.6 \times 10^2 \text{ GeV} \ g^{-3/2} & (n = 4, k = 0.01, N_e = 56), \\
2.0 \times 10^7 \text{ GeV} \ g^{-4/5} & (n = 5, k = 0.01, N_e = 50), \\
1.1 \times 10^9 \text{ GeV} \ g^{-5/9} & (n = 6, k = 0.01, N_e = 50).
\end{cases}$$

(15)
Figure 3: A schematic picture for the scalar potential of $\phi$. The two lines show the trajectories of the inflaton with angular initial condition with either $\theta = 0$ or $\theta \neq 0$. The later trajectory feels steeper potential, and hence, the spectral index becomes more red-tilted.

As we have shown in the previous section, the model with $n = 4$ is consistent with the Planck experiment only if $N_e \gtrsim 56$. This requires a very large $v^2$, which in turn puts a lower bound on the gravitino mass. To see this, let us remind ourselves that $N_e$ is given by the inflation scale as

$$N_e = 52 - \ln \left( \frac{10^{12} \text{ GeV}}{v} \right),$$

for the pivot scale $k_* = 0.002 \text{ Mpc}^{-1}$. Here, we have assumed an instantaneous reheating after the inflation, which brings about the largest $N_e$ for a fixed inflation scale. From Eqs. (14), (15) and (16), we obtain a relation between $m_{3/2}$ and $N_e$, which is shown in Fig. 4. From the figure and the constraint $N_e \gtrsim 56$, we obtain a lower bound on the gravitino mass, $m_{3/2} > \mathcal{O}(10^8) \text{ GeV}$.

Next, let us discuss the model with $n > 4$. In Fig. 5, we show the gravitino mass for $n = 5, 6$ with $N_e = 50$, $k = 0.01$. In can be seen that the larger and smaller $n$ and $g$ are, the larger the gravitino mass is. Hence, we can derive a lower bound on $m_{3/2}$ from an upper bound on $g$ for the model with $n = 5$.

It should be noted that there is an upper bound on $g$ from the unitarity limit, which can be extracted by considering the leading radiative correction to the Kähler potential due to the coupling $g$,

$$\delta K \simeq \frac{5!}{(16\pi^2)^4} g^2 M_6 \phi \phi^\dagger,$$

for the pivot scale $k_* = 0.002 \text{ Mpc}^{-1}$. Here, we have assumed an instantaneous reheating after the inflation, which brings about the largest $N_e$ for a fixed inflation scale. From Eqs. (14), (15) and (16), we obtain a relation between $m_{3/2}$ and $N_e$, which is shown in Fig. 4. From the figure and the constraint $N_e \gtrsim 56$, we obtain a lower bound on the gravitino mass, $m_{3/2} > \mathcal{O}(10^8) \text{ GeV}$.

Next, let us discuss the model with $n > 4$. In Fig. 5, we show the gravitino mass for $n = 5, 6$ with $N_e = 50$, $k = 0.01$. In can be seen that the larger and smaller $n$ and $g$ are, the larger the gravitino mass is. Hence, we can derive a lower bound on $m_{3/2}$ from an upper bound on $g$ for the model with $n = 5$.

It should be noted that there is an upper bound on $g$ from the unitarity limit, which can be extracted by considering the leading radiative correction to the Kähler potential due to the coupling $g$,
where $M_*$ is the cutoff of the loop integration. By requiring the unitarity up to the Planck scale, i.e. $M_* \simeq M_{\text{PL}}$, the unitarity limit, $|\delta K| \lesssim \phi \phi^*$, leads to an upper bound on $g^P$

$$g \lesssim \frac{(16\pi^2)^2}{\sqrt{5!}} \simeq 2000.$$  \hspace{1cm} \text{(18)}

By substituting this upper limit into Eqs. (15) and (18), we obtain a lower bound on the gravitino mass, $m_{3/2} \gtrsim 100$ TeV for $n > 4$.

In summary, we find that the lower bound on the gravitino mass;

$$m_{3/2} \gtrsim 100 \text{ TeV},$$  \hspace{1cm} \text{(19)}

in the $R$-breaking new inflation model. For $n = 4$, the (much higher) lower limit on the gravitino mass is obtained to achieve the observed spectral index, while the milder limit for $n > 4$ is obtained from the size of the curvature perturbation. As stressed in the introduction, this lower bound is consistent with the observed Higgs mass of 125 GeV \cite{28, 30}.

4 \hspace{0.5cm} Baryon asymmetry and dark matter density

In this section, we argue that the baryon asymmetry as well as the dark matter density in the present universe can be explained consistently with the $R$-breaking inflation model. In the following, we concentrate on the model with $n = 5$, $k \simeq 0.01$ and $N_e = 50$.\footnote{This requirement based on $M_* = 1$ is equivalent to the Born unitarity up to the Planck scale.}
4.1 Baryon asymmetry

Thermal leptogenesis

Let us first discuss whether the thermal leptogenesis \[32\] can be achieved in the $R$-breaking new inflation model, that is, whether a reheating temperature $T_R$ can be high enough, $T_R \gtrsim 10^9 \text{ GeV} \ [33]$. 

First, let us consider an inflaton decay via Planck-suppressed dimension five interactions\[7\] in which the decay width of the inflaton $\Gamma_{\phi, \text{dim}-5}$ is as large as $m_{\phi}^3$, where $m_{\phi}$ is the inflaton mass around the vacuum,

$$m_{\phi} = 5g \left( \frac{v^2}{g} \right)^{4/5} \simeq 1.4 \times 10^{11} \text{ GeV} \left( \frac{g}{1000} \right)^{-1/5}.$$ (20)

In this case, a reheating temperature $T_R$ is as large as

$$T_R \sim \sqrt{\Gamma_{\phi, \text{dim}-5}} \sim 10^7 \text{ GeV} \left( \frac{g}{1000} \right)^{-3/10} \ll 10^9 \text{ GeV}.$$ (21)

Therefore, for a successful thermal leptogenesis, we are lead to introduce unsuppressed interactions\[8\]

---

\[7\] For example, a Kähler interaction $K = \lambda_{\phi} Q Q$, where $Q$ is some chiral field lighter than the inflaton, provides such decay channel.

\[8\] If the dimension five interaction saturates the unitarity bound, $\lambda \sim 4\pi$, $T_R$ is as large as $10^8$ GeV. When the right-handed neutrinos have a non-hierarchical mass spectrum and the neutrino Yukawa matrix is rather tuned, the thermal leptogenesis is possible \[34\,36\].
In order to enhance the decay rate of the inflaton, let us consider a superpotential

\[ W = \frac{y}{2\ell} \phi^\dagger QQ, \]  

(22)

where \( Q \) is some chiral field lighter than the inflaton and \( y \) is a coupling constant. Due to large \( \langle \phi \rangle \),

\[ \langle \phi \rangle = \left( \frac{v^2}{g} \right)^{1/5} \simeq 2 \times 10^{-3} \left( \frac{g}{1000} \right)^{-3/10}, \]  

(23)

the decay of the inflaton by this interaction is effective even if \( \ell > 1 \). The decay width of \( \phi \) by this operator is given by

\[ \Gamma_{\phi} = \frac{\ell^2}{8\pi} \langle \phi \rangle \langle \phi \rangle^{2\ell-2} m_{\phi} = \frac{\ell^2}{8\pi} \left( \frac{g}{1000} \right)^{1/5} \left( \frac{g_*}{200} \right)^{-1/4} m_{Q}, \]  

(24)

where \( m_Q \) is the mass of \( Q \). A reheating temperature is given by

\[ T_R \simeq \left( \frac{90}{\pi^2 g_*} \right)^{1/4} \sqrt{\Gamma_{\phi}} = 1.8 \times 10^9 \text{GeV} \left( \frac{m_Q}{5 \times 10^{10} \text{GeV}} \right) \left( \frac{\ell}{3} \right) \left( \frac{g}{1000} \right)^{1/5} \left( \frac{g_*}{200} \right)^{-1/4}, \]  

(25)

where \( g_* \) is the effective degree of freedom of the radiations. It can be seen that the thermal leptogenesis is marginally possible.

In the mentioned above reheating scenario, we have introduced a matter field \( Q \). Note that we cannot identify \( Q \) with the minimal supersymmetric standard model (MSSM) higgs doublets, since a Dirac mass term of the MSSM higgs doublets, the so-called \( \mu \) term, should be as small as the gravitino mass, and hence a reheating temperature is not high enough (see Eq. (25)).

An interesting idea is to identify \( Q \) with the right-handed neutrinos, \( N_i \) \( (i = 1, 2, 3) \) [9]. In this case, the masses of the right-handed neutrinos, which should be far smaller than the Planck scale in order to obtain the observed masses of the left-handed neutrinos by the seesaw mechanism [37], are controlled by the \( Z_{2nR} \) symmetry rather than the \( B - L \) symmetry.

For example, let us arrange the right-handed neutrinos by their masses; \( m_{N_1} \leq m_{N_2} \leq m_{N_3} \). The inflaton decays mostly into the heaviest right-handed neutrino as long as the decay is kinematically allowed, that is, \( 2m_{N_1} < m_{\phi} \). If the inflaton decays mostly into \( N_2 \) or \( N_3 \) and the resulting reheating temperature is larger enough than \( m_{N_1} \), the thermal leptogenesis is marginally possible.
Non-thermal leptogenesis

We have shown that the thermal leptogenesis is marginally possible in the $R$-breaking new inflation model with $n = 5$. Interestingly, when we identify $Q$ with the right-handed neutrinos, a possibility of the non-thermal leptogenesis scenario [12] is also opened. There, the inflaton decays into right-handed neutrinos and the non-equilibrium decay of the right-handed neutrinos with a $CP$ violation generates lepton numbers.

For simplicity, let us assume that the inflaton decays mostly into the lightest right-handed neutrino $N_1$. The entropy yield of the baryon number is given by [38]

$$\eta_B \equiv \frac{n_B}{s} = 9 \times 10^{-11} \left( \frac{T_R}{10^6 \text{GeV}} \right) \left( \frac{2m_{N_1}}{m_\phi} \right) \left( \frac{m_{\nu 3}}{0.05 \text{eV}} \right) \frac{1}{\sin^2 \beta \delta_{\text{eff}}},$$

where $m_{\nu 3}$ is the mass of the heaviest left-handed neutrino, and $\beta$ is defined by the vacuum expectation values of the up-type and down-type higgs doublets, $H_u$ and $H_d$, as $\tan \beta = \langle H_u \rangle / \langle H_d \rangle$. $\delta_{\text{eff}}$ represents a degree of the $CP$ violation, which is given by the Yukawa couplings of the right-handed neutrinos, and expected be of order one. Compared with the observed value, $\eta_{B-\text{obs}} \simeq 8.5 \times 10^{-11}$ [18], an appropriate baryon asymmetry can be generated in the non-thermal leptogenesis scenario.

4.2 Dark matter density

In the MSSM, there is a candidate for dark matter, the lightest supersymmetric particle (LSP). Here, we assume pure gravity mediation models/minimal split SUSY models [39, 40], in which the gaugino masses are generated only by one-loop effects and hence smaller in comparison with the gravitino, higgsino and sfermion masses, and the wino is the LSP. The wino mass $M_2$ is given by [41]

$$M_2 = \frac{g_2^2}{16\pi^2} \left( m_{3/2} + L \right),$$

9 If we introduce a Kähler interaction $K = \phi^4 NN$ instead of the superpotential given by Eq. (22), a reheating temperature is as large as $10^7$ GeV (Eq. (21)) and the non-thermal leptogenesis is possible. In this case, the right-handed neutrinos has an $R$ charge of one, and the masses of the right-handed neutrinos are in general of order the Planck scale. In order to obtain $m_N < m_\phi$ as well as the observed masses of the left-handed neutrinos, some tunings are necessary. If we further assume that the scale $v^2$ is given by a breaking of some charged field, the masses of the right-handed neutrinos are also given by breaking of the charged field and hence is naturally small.
where \( g_2 \) is the \( SU(2) \) gauge coupling constant. The first term originates from an anomaly mediated effect \([41-43]\), while the second term, \( L \), parametrizes a higgsino threshold correction. As shown in Ref. \([39]\), \( L \) is expected to be of order the gravitino mass in pure gravity mediation models/minimal split SUSY models.

There are three sources for wino productions, a thermal wino relic, non-thermal production of gravitinos from a thermal bath, and gravitino production from the inflaton decay. We explain them in the following.

**Thermal wino relic**

Since the wino has an \( SU(2) \) gauge interaction, it is in a thermal equilibrium in the early universe. As the temperature of the universe decreases, the wino abundance freezes out and remains as a dark matter since the wino is the LSP. This is nothing but the conventional WIMP scenario. In order for the thermal abundance not to excess the observed cold dark matter value, \( \Omega_c h^2 = 0.1196 \pm 0.0031 \) \([18]\), it is required that \([47]\)

\[
M_2 \lesssim 3 \text{ TeV.} \tag{28}
\]

**Gravitino scattered from thermal bath**

Since the gravitino interacts with another light fields only through Planck-suppressed interactions, once it is scattered from a thermal bath, it does not interact with the thermal bath again, and eventually decays into the wino. A contribution to the wino abundance from this process is given by \([48-50]\)

\[
\Omega_{\text{wino,sc}} h^2 \simeq 0.12 \left( \frac{M_2}{200 \text{ GeV}} \right) \left( \frac{T_R}{10^{10} \text{ GeV}} \right). \tag{29}
\]

**Gravitino from inflaton decay**

After SUSY breaking, there is no remaining symmetry which prevents a mixing between the inflaton field and the SUSY breaking field at the vacuum. This effect induces an

\footnote{If there is a vector-like matter in addition to the MSSM fields, the gaugino masses receive a one-loop correction further \([44,45]\). For a comprehensive discussion on the phenomenology of the gauginos in that case, see Ref. \([46]\).}
inflaton decay into gravitinos \([51,58]\), which provides another source of non-thermal wino dark matter.

As an example, let us take the following effective superpotential for the SUSY breaking field \(Z\),

\[
W_{\text{eff}} = \Lambda^2 Z, \tag{30}
\]

where \(\Lambda^2\) is a SUSY breaking scale, which should satisfy \(\Lambda^2 = \sqrt{3}m_{3/2}\) in our flat universe.\(^{11}\) By calculating the scalar potential of the scalar components of \(Z\) and \(\delta \phi \equiv \phi - \langle \phi \rangle\) including supergravity effects, we obtain a mixing term,

\[
V_{\text{mix}} = \sqrt{3}(1 - b)m_\phi \langle \phi \rangle m_{3/2} \delta \phi Z^\dagger + \text{h.c.}, \tag{31}
\]

where \(b\) is a coupling constant in the Kähler potential, \(K \supset bZZ^\dagger \phi \phi^\dagger\). A mixing angle \(\epsilon\) between the scalar components of \(Z\) and \(\delta \phi\) is given by

\[
\epsilon = \sqrt{3}(1 - b)m_\phi \langle \phi \rangle m_{3/2} / m_Z^2, \tag{32}
\]

where \(m_Z\) is the mass of the SUSY breaking field. Here it is assumed that \(m_Z \gg m_\phi\), which is the case with typical dynamical SUSY breaking models.\(^{12}\)

A coupling between the scalar component of \(Z\) and its fermionic component \(\psi\), the goldstino, is provided by the following Kähler potential which gives a mass to the scalar component of the SUSY breaking field \([53,59]\),

\[
K \supset -\frac{m_Z^2}{12m_{3/2}^2} ZZ^\dagger ZZ^\dagger. \tag{33}
\]

The \(D\) term of Eq. (33) yields

\[
\mathcal{L} \supset -\frac{\sqrt{3}}{6} \frac{m_Z^2}{m_{3/2}^2} Z^\dagger \psi \psi + \text{h.c.} \tag{34}
\]

From Eqs (31) and (33), the decay rate is given by

\[
\Gamma_{3/2} \equiv \Gamma_{\phi \to 2\psi_{3/2}} \simeq \Gamma_{Z \to 2\phi, m_Z = m_\phi} |\epsilon|^2 = \frac{(b - 1)^2}{32\pi} m_\phi^3 \langle \phi \rangle^2. \tag{35}
\]

\(^{11}\)We have assumed that \(|\langle Z\rangle| \ll 1\) to avoid the Polonyi problem.

\(^{12}\)If not, an inflaton decay into gravitinos is suppressed \([53,54]\). An inflaton decay into SUSY breaking sector fields, which are expected to exist in general dynamical SUSY breaking models, can be also suppressed by separating the dynamical scale and the mass of \(Z\), \(m_Z \ll \Lambda\) \([59]\).
The entropy yield of the gravitino after the inflaton decay, $Y_{3/2}$, is estimated as

$$Y_{3/2} = 2 \times \frac{\Gamma_{3/2}}{\Gamma_{\text{tot}}} \frac{3 T_R}{4 m_{\phi}} = \frac{3}{2} \frac{\pi^2 g_*}{90} \frac{\Gamma_{3/2}}{m_{\phi} T_R},$$

(36)

where $\Gamma_{\text{tot}}$ is a total decay width of the inflaton. The wino abundance is given by

$$\Omega_{\text{wino,dec}} h^2 = \left( \frac{M_2}{3.5 \times 10^{-9} \text{ GeV}} \right) \times Y_{3/2}.$$  

(37)

In Fig. 6, we show constraints on the gravitino mass and the reheating temperature from the wino abundance in a $(m_{3/2}, T_R)$ plane, which is obtained by Eqs. (28), (29) and (37). Here, we have assumed that the wino mass is given by the purely anomaly mediated effect, $M_2 \approx 3 \times 10^{-3} m_{3/2}$. The figure shows that the observed dark matter density is mainly explained by the non-thermal contributions. If the coupling constant in the Kähler potential, $b$, is close to unity, the mixing between the SUSY breaking field and the inflaton is suppressed and hence the contribution from the inflaton decay is small.

We have also shown constraints from the baron asymmetry in the non-thermal leptogenesis scenario. The reheating temperature is identified with the one given in Eq. (25). In the lowest colored region, the generated baryon asymmetry is smaller than the observed value even if the CP violation is maximum, $\delta_{\text{eff}} = 1$. The result is insensitive to $\tan \beta$ as long as $\tan \beta > 1$. It can be seen that there is a portion of parameter space in which the baryon asymmetry as well as the dark matter density in the present universe is explained.

5 Summary and discussion

In this paper, we have investigated a compatibility of the supersymmetric $R$-breaking new inflation model with the results of the Planck experiment. We have shown that a lower bound on the gravitino mass, $m_{3/2} > \mathcal{O}(100) \text{ TeV}$, is obtained from the result of the Planck experiment. We have also shown that the baryon asymmetry as well as the dark matter density in the present universe can be explained consistently with the $R$-breaking inflation model.

As a final remark, let us interpret the gravitino mass from the landscape point of view [60–63]. In the landscape of vacua, it is possible that the gravitino mass is biased to low energy scales in order to obtain the electroweak scale as naturally as possible. In
Figure 6: Constraint on the gravitino mass and the reheating temperature from the wino abundance and the successful non-thermal leptogenesis scenario. Here, we have assumed the wino mass $M_2$ in Eq. (27) with $L = 0$. In this case, the nature should choose the gravitino mass which saturates the lower bound given by Eq. (19). Therefore, the gravitino mass, $m_{3/2} \simeq 100$ TeV, is a prediction in the $R$-breaking new inflation model in the landscape point of view.

It should be cautioned that there is a hidden parameter in this argument, $k$, which has been fixed $k \simeq 0.01$ to account for the observed spectral index. From the anthropic point of view, however, there seems no reason for the spectral index to be close to unity as observed. If we allow for a spectral index as large as 0.8, for example, then the gravitino

13 If there is a severer bound on $g$ than the unitarity bound, a larger gravino mass, such as PeV, is predicted from the landscape point of view. This argument may support the explanation of the PeV IceCube neutrino events by decaying gravitino dark matter.
mass is lowered down to,

\[ m_{3/2} \simeq 1.9 \times 10^3 \text{ GeV} \times \left( \frac{g}{2000} \right)^{-4/5} \quad (n = 5, k = 0.1, N_e = 50), \tag{38} \]

which is much smaller than 100 TeV.

This shows that our landscape argument is self-consistent only if the parameter \( k \) is fixed to be close to 0.01 by some underlying theory. If not, the landscape argument predicts that \( k \sim 0.1 \) and \( m_{3/2} \sim 1 \text{ TeV} \), in which the electroweak scale is obtained much more naturally than the case with \( k \sim 0.01 \) and \( m_{3/2} \sim 100 \text{ TeV} \), and the prediction already contradicts with the observed value of the spectral index.

This situation is similar to anthropic arguments \cite{64} on the electroweak scale. It is argued that an electroweak scale of the one realized in the nature is required for the people to exist in the universe \cite{67,68}. There, other parameters other than the Higgs boson mass in the standard model such as the gauge coupling constants and the Yukawa couplings are fixed to the observed value. The anthropic prediction on the electroweak scale is viable only if all such couplings are consider to be fixed by some underlying theory.

Instead of fixing the parameter \( k \), we may move ahead with the landscape point of view under an additional assumption. Suppose that the parameter with the positive mass dimension in the superpotential, \( v \), is strongly biased to larger mass scales. However, \( v \) is anthropically required to be sufficiently small in order to generate a small cosmological perturbation, \( \mathcal{P}_\zeta \sim 10^{-9} \). Consequently, the maximum \( v \) on the hyper-surface of the parameter space corresponding to \( \mathcal{P}_\zeta \sim 10^{-9} \) would have been chosen anthropically. In Fig. \cite{7} we show a line in a \( k - v \) space in which \( \mathcal{P}_\zeta = 2.2 \times 10^{-9} \). It can be seen that \( k \sim 10^{-2} \), which is consistent with the observed spectral index, gives the maximum \( v \).

Note that the result is insensitive to the parameter \( g \). It is remarkable that a high energy biased \( v \) explains the reason why the spectral index \( n_s \) is not too small such as 0.8 but close to the observed value, i.e. \( n_s \sim 0.96 \).

**Acknowledgments**

This work is supported by Grant-in-Aid for Scientific research from the Ministry of Education, Science, Sports, and Culture (MEXT), Japan, No. 22244021 (T.T.Y.), No. 24740151.
Figure 7: A line in a $k - v$ space in which $P_\zeta = 2.2 \times 10^{-9}$.

(M.I), and also by World Premier International Research Center Initiative (WPI Initiative), MEXT, Japan. The work of K.H. is supported in part by a JSPS Research Fellowships for Young Scientists.

References

[1] L. Maiani. in Proceedings: Summer School on Particle Physics, Paris, France (1979).
[2] M. J. G. Veltman, Acta Phys. Polon. B 12, 437 (1981).
[3] E. Witten, Nucl. Phys. B 188, 513 (1981).
[4] R. K. Kaul, Phys. Lett. B 109, 19 (1982), and references therein.
[5] M. Dine, G. Festuccia and Z. Komargodski, JHEP 1003, 011 (2010) [arXiv:0910.2527 [hep-th]].
[6] G. D. Coughlan, W. Fischler, E. W. Kolb, S. Raby and G. G. Ross, Phys. Lett. B 131, 59 (1983); M. Ibe, Y. Shinbara and T. T. Yanagida, Phys. Lett. B 639, 534 (2006) [hep-ph/0605252]; see also K. Harigaya, M. Ibe, K. Schmitz and T. T. Yanagida, Phys. Lett. B 721, 86 (2013) [arXiv:1301.3685 [hep-ph]].
[7] A. D. Linde, Phys. Lett. B 108, 389 (1982).
[8] A. Albrecht and P. J. Steinhardt, Phys. Rev. Lett. 48, 1220 (1982).
[9] K. Kumekawa, T. Moroi and T. Yanagida, Prog. Theor. Phys. 92, 437 (1994) [hep-ph/9405337].
[10] K. -I. Izawa and T. Yanagida, Phys. Lett. B 393, 331 (1997) [hep-ph/9608359].
[11] K. I. Izawa, M. Kawasaki and T. Yanagida, Phys. Lett. B 411, 249 (1997) [hep-ph/9707201].
[12] M. Ibe, K. -I. Izawa, Y. Shinbara and T. T. Yanagida, Phys. Lett. B 637, 21 (2006) [hep-ph/0602192].
[13] M. Ibe, Y. Shinbara and T. T. Yanagida, Phys. Lett. B 642, 165 (2006) [hep-ph/0608127].
[14] F. Takahashi, arXiv:1308.4212 [hep-ph].
[15] Y. B. Zeldovich, I. Y. Kobzarev and L. B. Okun, Zh. Eksp. Teor. Fiz. 67, 3 (1974) [Sov. Phys. JETP 40, 1 (1974)].
[16] D. H. Lyth, Phys. Rev. Lett. 78, 1861 (1997) [hep-ph/9606387].
[17] G. Hinshaw et al. [WMAP Collaboration], Astrophys. J. Suppl. 208, 19 (2013) arXiv:1212.5226 [astro-ph.CO].
[18] P. A. R. Ade et al. [Planck Collaboration], arXiv:1303.5076 [astro-ph.CO].
[19] P. A. R. Ade et al. [Planck Collaboration], arXiv:1303.5082 [astro-ph.CO].
[20] A. D. Linde, Phys. Lett. B 129, 177 (1983).
[21] R. Kallosh and A. Linde, JCAP 1011, 011 (2010) arXiv:1008.3375 [hep-th];
R. Kallosh, A. Linde and T. Rube, Phys. Rev. D 83, 043507 (2011) arXiv:1011.5945 [hep-th].
[22] F. Takahashi, Phys. Lett. B 693, 140 (2010) arXiv:1006.2801 [hep-ph]; K. Nakayama
and F. Takahashi, JCAP 1011, 009 (2010) arXiv:1008.2956 [hep-ph].
[23] K. Harigaya, M. Ibe, K. Schmitz and T. T. Yanagida, Phys. Lett. B 720, 125 (2013) arXiv:1211.6241 [hep-ph].
[24] D. Croon, J. Ellis and N. E. Mavromatos, Physics Letters B 724, , 165 (2013) arXiv:1303.6253 [astro-ph.CO].
[25] K. Nakayama, F. Takahashi and T. T. Yanagida, Phys. Lett. B 725, 111 (2013) arXiv:1303.7315 [hep-ph]; K. Nakayama, F. Takahashi and T. T. Yanagida, JCAP 1308, 038 (2013) arXiv:1305.5099 [hep-ph].
[26] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716, 1 (2012) arXiv:1207.7214 [hep-ex].

[27] S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 716, 30 (2012) arXiv:1207.7235 [hep-ex].

[28] Y. Okada, M. Yamaguchi and T. Yanagida, Prog. Theor. Phys. 85, 1 (1991); Y. Okada, M. Yamaguchi and T. Yanagida, Phys. Lett. B 262, 54 (1991).

[29] J. R. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. B 257, 83 (1991).

[30] H. E. Haber and R. Hempfling, Phys. Rev. Lett. 66, 1815 (1991).

[31] A. R. Liddle and D. H. Lyth, Phys. Rept. 231, 1 (1993) astro-ph/9303019.

[32] M. Fukugita and T. Yanagida, Phys. Lett. B 174, 45 (1986).

[33] W. Buchmuller, P. Di Bari and M. Plumacher, Annals Phys. 315, 305 (2005) hep-ph/0401240.

[34] M. Flanz, E. A. Paschos, U. Sarkar and J. Weiss, Phys. Lett. B 389, 693 (1996) hep-ph/9607310.

[35] A. Pilaftsis, Phys. Rev. D 56, 5431 (1997) hep-ph/9707235.

[36] S. Blanchet and P. Di Bari, Nucl. Phys. B 807, 155 (2009) arXiv:0807.0743 [hep-ph].

[37] T. Yanagida, in Proceedings of the Workshop on Unified Theory and Baryon Number of the Universe, eds; O. Sawada and A. Sugamoto (KEK, 1979) p.95; M. Gell-Mann, P. Ramond and R. Slansky, in Supergravity, eds.; P. van Niewwenhuizen and D. Freedman (North Holland, Amsterdam, 1979). See also P. Minkowski, Phys. Lett. B 67, 421 (1977).

[38] M. Ibe, T. Moroi and T. Yanagida, Phys. Lett. B 620, 9 (2005) hep-ph/0502074.

[39] M. Ibe, T. Moroi and T. T. Yanagida, Phys. Lett. B 644, 355 (2007) hep-ph/0610277; M. Ibe and T. T. Yanagida, Phys. Lett. B 709, 374 (2012) arXiv:1112.2462 [hep-ph]; M. Ibe, S. Matsumoto and T. T. Yanagida, Phys. Rev. D 85, 095011 (2012) arXiv:1202.2253 [hep-ph].

[40] N. Arkani-Hamed, IFT Inaugural Conference (2011), http://www.ift.uam.es/workshops/Xmas11/?q=node/2; N. Arkani-Hamed, A. Gupta, D. E. Kaplan, N. Weiner and T. Zorawski, arXiv:1212.6971 [hep-ph].
[41] G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP 9812, 027 (1998) [hep-ph/9810442].

[42] M. Dine and D. MacIntire, Phys. Rev. D 46, 2594 (1992) [hep-ph/9205227].

[43] L. Randall and R. Sundrum, Nucl. Phys. B 557, 79 (1999) [hep-th/9810155].

[44] A. E. Nelson and N. J. Weiner, hep-ph/0210288.

[45] K. Nakayama and T. T. Yanagida, Phys. Lett. B 722, 107 (2013) arXiv:1302.3332 [hep-ph]].

[46] K. Harigaya, M. Ibe and T. T. Yanagida, arXiv:1310.0643 [hep-ph].

[47] J. Hisano, S. Matsumoto, M. Nagai, O. Saito and M. Senami, Phys. Lett. B 646, 34 (2007) hep-ph/0610249.

[48] M. Kawasaki and T. Moroi, Prog. Theor. Phys. 93, 879 (1995) [hep-ph/9403364, hep-ph/9403061]; M. Kawasaki, K. Kohri, T. Moroi and A. Yotsuyanagi, Phys. Rev. D 78, 065011 (2008) arXiv:0804.3745 [hep-ph]].

[49] T. Gherghetta, G. F. Giudice and J. D. Wells, Nucl. Phys. B 559, 27 (1999) hep-ph/9904378.

[50] M. Ibe, R. Kitano, H. Murayama and T. Yanagida, Phys. Rev. D 70, 075012 (2004) hep-ph/0403198; M. Ibe, R. Kitano and H. Murayama, Phys. Rev. D 71, 075003 (2005) hep-ph/0412200.

[51] M. Kawasaki, F. Takahashi and T. T. Yanagida, Phys. Lett. B 638, 8 (2006) hep-ph/0603265.

[52] T. Asaka, S. Nakamura and M. Yamaguchi, Phys. Rev. D 74, 023520 (2006) hep-ph/0604103.

[53] M. Dine, R. Kitano, A. Morisse and Y. Shirman, Phys. Rev. D 73, 123518 (2006) hep-ph/0604140.

[54] M. Endo, K. Hamaguchi and F. Takahashi, Phys. Rev. D 74, 023531 (2006) hep-ph/0605091.

[55] M. Kawasaki, F. Takahashi and T. T. Yanagida, Phys. Rev. D 74, 043519 (2006) hep-ph/0605297.
[56] M. Endo, M. Kawasaki, F. Takahashi and T. T. Yanagida, Phys. Lett. B 642, 518 (2006) [hep-ph/0607170].
[57] M. Endo, F. Takahashi and T. T. Yanagida, Phys. Lett. B 658, 236 (2008) [hep-ph/0701042].
[58] M. Endo, F. Takahashi and T. T. Yanagida, Phys. Rev. D 76, 083509 (2007) arXiv:0706.0986 [hep-ph].
[59] K. Nakayama, F. Takahashi and T. T. Yanagida, Phys. Lett. B 718, 526 (2012) arXiv:1209.2583 [hep-ph].
[60] R. Bousso and J. Polchinski, JHEP 0006, 006 (2000) hep-th/0004134.
[61] S. Kachru, R. Kallosh, A. D. Linde and S. P. Trivedi, Phys. Rev. D 68, 046005 (2003) hep-th/0301240.
[62] L. Susskind, In *Carr, Bernard (ed.): Universe or multiverse?* 247-266 hep-th/0302219.
[63] F. Denef and M. R. Douglas, JHEP 0405, 072 (2004) hep-th/0404116.
[64] S. Weinberg, Phys. Rev. Lett. 59, 2607 (1987).
[65] M. G. Aartsen et al. [IceCube Collaboration], Phys. Rev. Lett. 111, 021103 (2013) arXiv:1304.5356 [astro-ph.HE]].
[66] B. Feldstein, A. Kusenko, S. Matsumoto and T. T. Yanagida, Phys. Rev. D 88, 015004 (2013) arXiv:1303.7320 [hep-ph].
[67] V. Agrawal, S. M. Barr, J. F. Donoghue and D. Seckel, Phys. Rev. Lett. 80, 1822 (1998) hep-ph/9801253.
[68] T. E. Jeltema and M. Sher, Phys. Rev. D 61, 017301 (2000) hep-ph/9905494.