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Lower bound on the gravitino mass $m_{3/2} > O(100)$ 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*

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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, $m_{3/2} > O(100)$ TeV, assuming that the R-symmetry breaking is dominantly induced by the inflaton and the inflaton dynamics is not affected by the supersymmetry-breaking sector, which can be guaranteed by symmetry or compositeness of the supersymmetry-breaking sector. This lower limit is consistent with the observed Higgs mass of 126 GeV when the masses of the top squarks are of order the gravitino mass scale. 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, assuming leptogenesis and that the wino is the lightest supersymmetric particle.

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I. 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 the SUSY-breaking scale to be the Planck scale in order 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 [1–4] between the weak scale and the Planck scale or the scale of the grand unified theory.

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 nonvanishing F-term VEV at the same time if the symmetry is a continuous one [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 [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,[1] which evokes a scalar potential used in new inflation models [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 [9–14]. It is also remarkable that the domain-wall problem [15] associated with the discrete R-symmetry breaking is automatically solved when the R-symmetry-breaking field plays the role of the inflaton.[2]

We here emphasize that new inflation models tend to predict a small tensor fraction due to their small inflation scales [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 [17–19].[3]

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, $m_{3/2} > O(100)$ TeV, assuming that the R-symmetry breaking is dominantly induced by the inflaton and the inflaton dynamics is not affected by the

[1]Reference [9] pointed out not only the presence of the so-called η problem in supergravity inflation models but also the importance of the R symmetry to have flat potentials necessary for the inflation to occur.

[2]This situation is analogous to the original new inflation model [7,8], where an inflaton is identified with a grand unified theory-breaking field and the monopole problem is solved.

[3]Simple large-field inflation models such as the chaotic inflation models with a quadratic or a quartic potential [20] are, on the other hand, now slightly disfavored at least at the 1σ level, which requires some extensions [21–25].

SUSY-breaking sector.[4] 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 top squarks 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, assuming leptogenesis [31] and that the wino is the lightest SUSY particle.

II. BRIEF REVIEW OF 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 ϕ [9,10]. Here, we assume that ϕ is a singlet except for the R symmetry with an R charge 2. Assuming $N = 2n$, the superpotential of ϕ is given by

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

where the ellipses represent higher-power terms of ϕ. 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 use the unit in which the reduced Planck scale $M_{\rm PL} \simeq 2.4 \times 10^{18}$ GeV is unity. The parameters v^2 and g are taken real and positive without loss of generality.[5] At supersymmetric vacua, the Z_{2nR} symmetry is spontaneously broken down to the Z_{2R} symmetry by the VEV of ϕ,

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

which leads to the VEV of the superpotential,[6]

$$\langle W\rangle \simeq \frac{n}{n+1} v^2 \left(\frac{v^2}{g}\right)^{1/n} e^{2\pi i m/n}. \tag{3}$$

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 ϕ is set close to its origin by, for example, a positive Hubble induced mass term of preinflation [11] and the slow-roll conditions are satisfied, ϕ 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 this purpose, let us note that the Kähler potential of ϕ is given by

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

where the ellipses denote higher-power terms of ϕ, whose contributions to the dynamics of ϕ 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 ϕ is given by

$$\begin{aligned} V(\phi) &= |v^2 - g\phi^n|^2 - kv^4|\phi|^2 + \cdots \\ &= v^4 - (gv^2\phi^n + \text{H.c.}) - kv^4|\phi|^2 \cdots. \end{aligned} \tag{5}$$

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

$$V(\varphi, \theta) = v^4 - \frac{k}{2}v^4\varphi^2 - \frac{g}{2^{n/2-1}}v^2\varphi^n \cos(n\theta) + \cdots. \tag{6}$$

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

As we have mentioned, we assume that the initial condition of φ is close to 0, i.e., $|\varphi| \ll 1$. We further suppose that the initial condition of the angular direction θ is given by $\theta = 0(\text{m}od2\pi/n)$ for the time being. Since $\theta = 0(\text{m}od2\pi/n)$ is the minimum of the potential along the angular direction, $\theta = 0(\text{m}od2\pi/n)$ is kept during the inflation. Along the inflaton trajectory, the first and second slow-roll parameters are given, respectively, by

$$\begin{aligned} \epsilon &\equiv \frac{1}{2}\left(\frac{\partial V/\partial\varphi}{V}\right)^2 = \frac{1}{2}\left(k\varphi + \frac{ng}{2^{n/2-1}}\frac{\varphi^{n-1}}{v^2}\right)^2, \\ \eta &\equiv \frac{\partial^2 V/\partial\varphi^2}{V} = -k - \frac{n(n-1)g}{2^{n/2-1}}\frac{\varphi^{n-2}}{v^2}. \end{aligned} \tag{7}$$

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

[4]These two assumptions are realized in SUSY-breaking models such that fields in the SUSY-breaking sector are charged or composite. This is because couplings between the SUSY-breaking sector and the inflaton as well as the vacuum expectation value of the SUSY-breaking field are suppressed due to the charge or the compositeness. Note that such SUSY-breaking models are free from the Polonyi problem [6].

[5]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 a $\mathrm{U}(1)_R$ symmetry under which ϕ has a charge of $2/(n+1)$, and the $\mathrm{U}(1)_R$ symmetry is dynamically broken by a condensation of a (composite) chiral field with a $\mathrm{U}(1)_R$ charge of $2 - 2(n+1)$ [10].

[6]At the vacuum, the potential energy of the inflaton is as large as $-3|\langle W\rangle|^2$ and negative. This negative contribution must be canceled with the positive contribution to the potential energy from SUSY breaking. As we have mentioned in the Introduction, we assume that couplings between the inflaton and the SUSY-breaking sector are suppressed and hence the dynamics of the inflaton is not affected by the SUSY-breaking sector. As we will see, Planck-suppressed interactions do not affect the inflaton dynamics in the parameter space of interest.

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

$$\varphi_{\rm end} = \left(\frac{2^{(n-2)/2} v^2}{n(n-1)g}\right)^{1/(n-2)}, \tag{8}$$

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

$$N_e(\varphi) = \int_{\varphi_{\rm end}}^{\varphi} \frac{V}{\partial V/\partial\varphi} {\rm d}\varphi. \tag{9}$$

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

$$\varphi^{n-2}(N_e) = \frac{2^{(n-2)/2} k v^2}{ng} (e^{k(n-2)N_e} - 1 + k(n-1)e^{k(n-2)N_e})^{-1}. \tag{10}$$

In order to compare model predictions with cosmic microwave background observations, let us calculate the properties of the curvature perturbation. The spectrum of the curvature perturbation $\mathcal{P}_\zeta$ and its spectral index n_s are given by

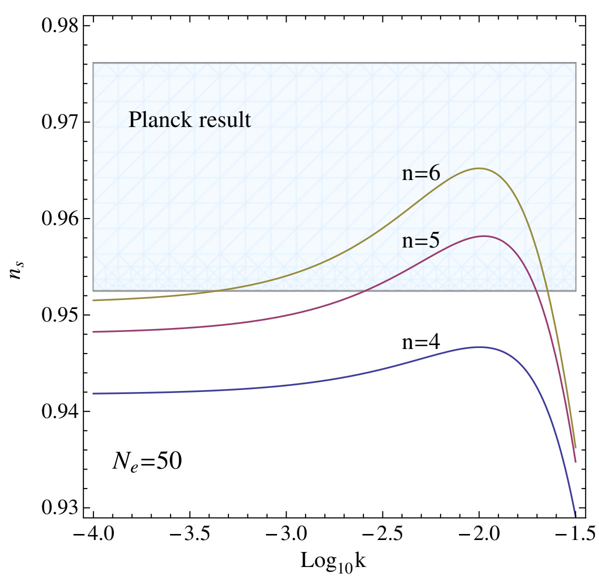

FIG. 1 (color online). The spectral index of the curvature perturbation n_s for $n = 4, 5, 6$ with $N_e = 50$. The colored region shows the 95% C.L. region favored by the Planck experiment, $n_s = 0.9643 \pm 0.012$ [19].

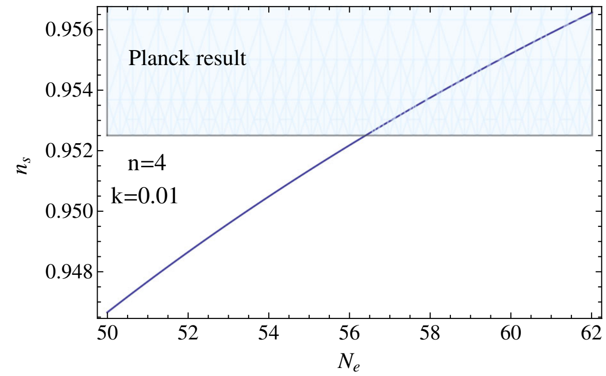

FIG. 2 (color online). The spectral index of the curvature perturbation n_s for $n = 4$ with various N_e. The colored region shows the 95% C.L. limit from the Planck experiment, $n_s = 0.9643 \pm 0.012$ [19].

$$\mathcal{P}_\zeta = \frac{1}{24\pi^2}\frac{V}{\epsilon} = \frac{1}{24\pi^2}(n^2 g^2 k^{-2(n-1)} v^{4(n-3)} \times (e^{k(n-2)N_e} - 1)^{2(n-1)})^{\frac{1}{n-2}} e^{-2k(n-2)N_e}, \tag{11}$$

$$n_s = 1 - 6\epsilon + 2\eta = 1 - 2k\left(1 + \frac{n-1}{(1+k(n-1))e^{k(n-2)N_e} - 1}\right), \tag{12}$$

respectively. In Fig. 1, we show the prediction for 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 $k_* = 0.002\ {\rm Mpc}^{-1}$ at 95% C.L. It can be seen that the model with $n \le 4$ 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 for $n = 4$. The figure shows that the model with $n = 4$ is still

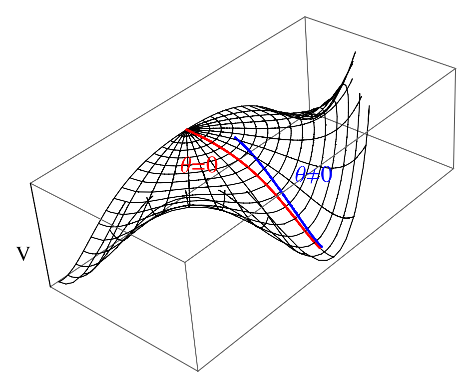

FIG. 3 (color online). A schematic picture for the scalar potential of ϕ. The two lines show the trajectories of the inflaton with angular initial condition with either $\theta = 0$ or $\theta \neq 0$. The later trajectory feels a steeper potential, and hence the spectral index becomes more red-tilted.

consistent with the Planck experiment for $N_e \gtrsim 56$.[7] We will discuss impacts of the observed spectral index on the gravitino mass in the next section.

Before closing this section, let us discuss more general initial conditions for the inflaton field, $\theta \neq 0(\mathrm{m}od 2\pi/n)$. In particular, we are interested in how the spectral index is affected, since $n = 4$ is severely constrained for $\theta = 0(\mathrm{m}od 2\pi/n)$ 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 $\mathrm{Re}(\phi) > 0$. For a fixed number of e-foldings, a nonzero angle θ leads to a larger corresponding field value for φ. 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 $\theta \neq 0(\mathrm{m}od 2\pi/n)$, the model with $n = 4$ is still disfavored unless N_e is large. For a rigorous discussion with the δN formalism [32–34], see Ref. [35].

III. 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

$$v^2 = \left((24\pi^2 \mathcal{P}_\zeta)^{n-2}(ng)^{-2}\left(\frac{k}{e^{k(n-2)N_e} - 1}\right)^{2(n-1)} e^{2k(n-2)^2 N_e}\right)^{\frac{1}{2(n-3)}}, \tag{13}$$

which leads to

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

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

$$m_{3/2} = \frac{ng}{n+1}\left(\frac{v^2}{g}\right)^{\frac{n+1}{n}} \simeq \begin{cases} 1.6 \times 10^{2}\ \mathrm{GeV} g^{-3/2} & (n = 4, k = 0.01, N_e = 56), \\ 2.0 \times 10^{7}\ \mathrm{GeV} g^{-4/5} & (n = 5, k = 0.01, N_e = 50), \\ 1.1 \times 10^{9}\ \mathrm{GeV} g^{-5/9} & (n = 6, k = 0.01, N_e = 50). \end{cases} \tag{15}$$

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 [36]

$$N_e = 52 - \ln\left(\frac{10^{12}\ \mathrm{GeV}}{v}\right) \tag{16}$$

for the pivot scale $k_* = 0.002\ \mathrm{Mpc}^{-1}$. 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, $m_{3/2} > \mathcal{O}(10^8)$ 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 as n and g become smaller,

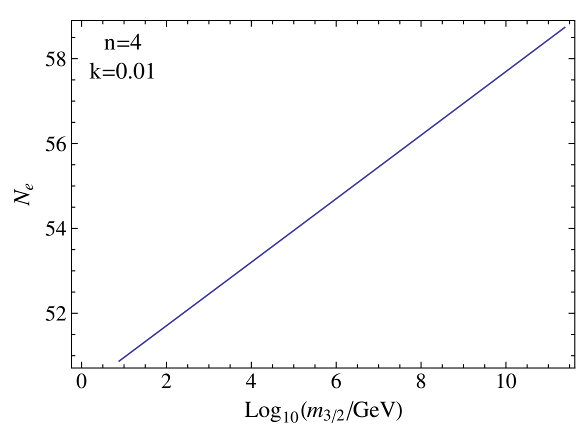

FIG. 4 (color online). A relation between $m_{3/2}$ and N_e for $n = 4$.

[7] In Ref. [14], it was pointed out that the model with $n = 4$ is also consistent with the Planck experiment if there is a small constant term in the superpotential besides the one from the condensation of ϕ. Since we assume that the R symmetry is broken only by the condensation of ϕ, that solution is not applicable.

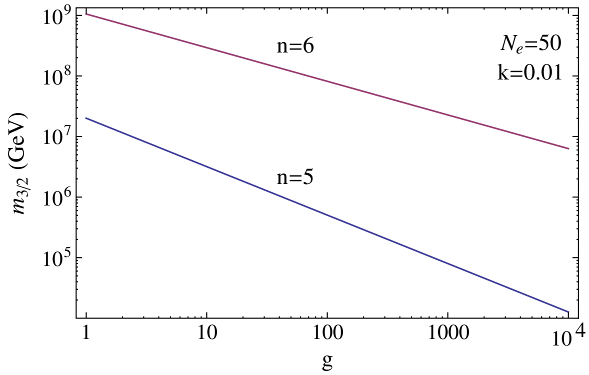

FIG. 5 (color online). The gravitino mass for $n = 5, 6$ with $N_e = 50$, $k = 0.01$.

the gravitino mass becomes larger. 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,

$$\delta K \simeq \frac{5!}{(16\pi^2)^4} g^2 M_*^6 \phi\phi^\dagger, \tag{17}$$

where M_* is the cutoff of the loop integration. By requiring unitarity up to the Planck scale, i.e., $M_* \simeq M_{\rm PL}$, the unitarity limit, $|\delta K| \lesssim \phi\phi^\dagger$, leads to an upper bound on g,[8]

$$g \lesssim (16\pi^2)^2/\sqrt{5!} \simeq 2000. \tag{18}$$

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 is

$$m_{3/2} \gtrsim 100 \text{ TeV} \tag{19}$$

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 [28–30].

Let us finally discuss the effect of Planck-suppressed couplings between the inflaton and the SUSY-breaking sector. Whatever the symmetry of the SUSY-breaking sector is, Planck-suppressed interactions in general give the inflaton a soft squared-mass term as large as $m_{3/2}^2$. In the parameter space of interest, $m_{3/2} \sim O(100)$ TeV, $m_{3/2}$ is much smaller than the Hubble scale during inflation, $H \sim v^2$, and hence the slow-roll condition is not affected by the soft squared-mass term.

[8]This requirement based on $M_* = 1$ is equivalent to the Born unitarity up to the Planck scale.

IV. 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 $N_e = 50$.

A. Baryon asymmetry

1. Thermal leptogenesis

Let us first discuss whether the thermal leptogenesis [31] can be achieved in the R-breaking new inflation model, that is, whether a reheating temperature T_R can be high enough, $T_R \gtrsim 10^9$ GeV [37].

First, let us consider an inflaton decay via Planck-suppressed dimension-five interactions[9] in which the decay width of the inflaton $\Gamma_{\phi,\dim-5}$ is as large as m_ϕ^3, where m_ϕ is the inflaton mass around the vacuum,

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

In this case, the reheating temperature T_R is as large as

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

Therefore, for a successful thermal leptogenesis, we are lead to introduce unsuppressed interactions.[10]

In order to enhance the decay rate of the inflaton, let us consider a superpotential

$$W = \frac{y}{2\ell}\phi^\ell QQ, \tag{22}$$

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

[9]For example, a Kähler interaction $K = \lambda\phi^\dagger QQ$, where Q is some chiral field lighter than the inflaton, provides such a decay channel.

[10]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 nonhierarchical mass spectrum and the neutrino Yukawa matrix is rather tuned, the thermal leptogenesis is possible [38–40].

$$\langle\phi\rangle = \left(\frac{v^2}{g}\right)^{1/5} \simeq 2\times 10^{-3}\left(\frac{g}{1000}\right)^{-3/10}, \quad (23)$$

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

$$\Gamma_\phi = \frac{1}{8\pi} y^2 |\langle\phi\rangle|^{2\ell-2} m_\phi = \frac{\ell^2}{8\pi}\frac{m_Q^2}{|\langle\phi\rangle|^2} m_\phi, \quad (24)$$

where m_Q is the mass of Q. The reheating temperature is given by

$$\begin{aligned} T_R &\simeq \left(\frac{90}{\pi^2 g_*}\right)^{1/4}\sqrt{\Gamma_\phi} \\ &= 1.8\times 10^9\ \mathrm{GeV}\left(\frac{m_Q}{5\times 10^{10}\ \mathrm{GeV}}\right)\left(\frac{\ell}{3}\right) \\ &\quad \times \left(\frac{g}{1000}\right)^{1/5}\left(\frac{g_*}{200}\right)^{-1/4}, \end{aligned} \quad (25)$$

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

In the above-mentioned 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 μ term—should be as small as the gravitino mass, and hence the 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 [41]—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: $m_{N_1} \le m_{N_2} \le m_{N_3}$. The inflaton decays mostly into the heaviest right-handed neutrino as long as the decay is kinematically allowed, that is, $2m_{N_i} < m_\phi$. 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.

2. *Nonthermal 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, the possibility of the nonthermal leptogenesis scenario [12] is also opened.[11] There, the inflaton decays into right-handed neutrinos and the nonequilibrium 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 [42]

$$\begin{aligned} \eta_B &\equiv \frac{n_B}{s} \\ &= 9\times 10^{-11}\left(\frac{T_R}{10^6\ \mathrm{GeV}}\right)\left(\frac{2m_{N_1}}{m_\phi}\right)\left(\frac{m_{\nu 3}}{0.05\ \mathrm{eV}}\right)\frac{1}{\sin^2\beta}\delta_{\mathrm{eff}}, \end{aligned} \quad (26)$$

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

B. 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 [43,44], in which the gaugino masses are generated only by one-loop effects and hence are smaller in comparison with the gravitino, Higgsino, and sfermion masses, and the wino is the LSP. The wino mass M_2 is given by [45]

$$M_2 = \frac{g_2^2}{16\pi^2}(m_{3/2} + L), \quad (27)$$

where g_2 is the SU(2) gauge coupling constant. The first term originates from an anomaly-mediated effect [45–47], while the second term, L, parametrizes a Higgsino threshold correction.[12] As shown in Ref. [43], L is expected to be

[11] If we introduce a Kähler interaction $K = \phi^\dagger NN$ instead of the superpotential given by Eq. (22), the reheating temperature is as large as 10^7 GeV [Eq. (21)] and the nonthermal leptogenesis is possible. In this case, the right-handed neutrinos have 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 the breaking of the charged field and hence are naturally small.

[12] If there is a vector-like matter in addition to the MSSM fields, the gaugino masses receive an additional one-loop correction [48,49]. For a comprehensive discussion on the phenomenology of the gauginos in this case, see Ref. [50].

of order the gravitino mass in pure gravity mediation models/minimal split SUSY models.

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

1. Thermal wino relic

Since the wino has an SU(2) gauge interaction, it is in 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 weakly interacting massive particle scenario. In order for the thermal abundance to not be in excess of the observed cold dark matter value, $\Omega_c h^2 = 0.1196 \pm 0.0031$ [18], it is required that [51]

$$M_2 \lesssim 3\ \text{TeV}. \tag{28}$$

2. Gravitino scattered from thermal bath

Since the gravitino interacts with other 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 [52–54]

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

3. Gravitino from inflaton decay

After SUSY breaking, there is no remaining symmetry that prevents a mixing between the inflaton field and the SUSY-breaking field at the vacuum. This effect induces an inflaton decay into gravitinos [55–62], which provides another source of nonthermal wino dark matter. The branching fraction of the inflaton into gravitinos must be suppressed, since otherwise the universe is over-closed by winos.

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 Λ^2 is a SUSY-breaking scale, which should satisfy $\Lambda^2 = \sqrt{3} m_{3/2}$ in our flat Universe.[13] 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 ϵ 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.[14]

A coupling between the scalar component of Z and its fermionic component ψ—the Goldstino—is provided by the following Kähler potential, which gives a mass to the scalar component of the SUSY-breaking field [57,63]:

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

The D-term of Eq. (33) yields

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

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

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

The entropy yield of the gravitino after the inflaton decay, $Y_{3/2}$, is estimated as

$$Y_{3/2} = 2 \times \frac{\Gamma_{3/2}}{\Gamma_{\text{tot}}}\frac{3T_R}{4m_\phi} = \frac{3}{2}\sqrt{\frac{\pi^2 g_*}{90}}\frac{\Gamma_{3/2}}{m_\phi T_R}, \tag{36}$$

where Γ_{tot} is the total decay width of the inflaton. The wino abundance is given by

$$\Omega_{\text{wino,dec}} h^2 = \left(\frac{M_2}{3.5 \times 10^{-9}\ \text{GeV}}\right) \times Y_{3/2}. \tag{37}$$

In Fig. 6, we show constraints on the gravitino mass and the reheating temperature from the wino abundance in the $(m_{3/2}, T_R)$ 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, $M_2 \simeq 3 \times 10^{-3} m_{3/2}$. The figure shows that the observed dark matter density is mainly explained by the nonthermal contributions. If the coupling constant in the Kähler potential, b, is close to

[13] We have assumed that $|\langle Z\rangle| \ll 1$ to avoid the Polonyi problem.

[14] If not, an inflaton decay into gravitinos is suppressed [57,58]. 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, $m_Z \ll \Lambda$ [63].

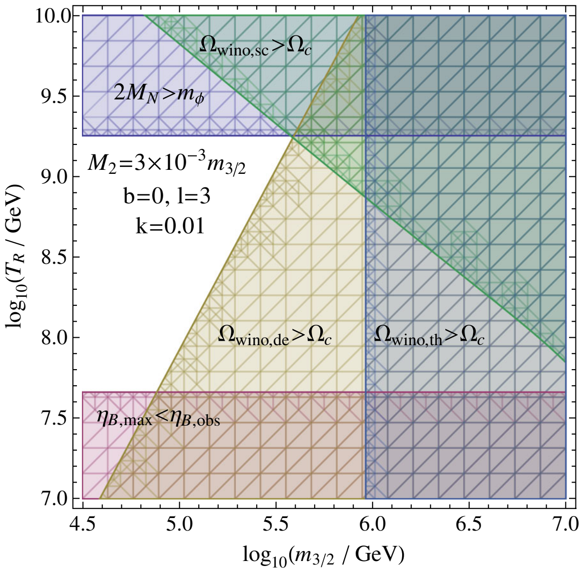

FIG. 6 (color online). Constraint on the gravitino mass and the reheating temperature from the wino abundance and the successful nonthermal leptogenesis scenario. Here, we have assumed the wino mass M_2 in Eq. (27) with $L = 0$.

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 nonthermal 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, $\delta_{\rm eff} = 1$. The result is insensitive to $\tan\beta$ as long as $\tan\beta \gtrsim 1$. 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.

Finally, let us comment about constraints on the wino dark matter. As discussed in Refs. [64,65], the wino dark matter is severely constrained by indirect dark matter searches, in particular by the gamma-ray line search by the H.E.S.S. experiments from the Galactic center. The constraint on the annihilation cross section of dark matter, however, suffers from large ambiguities in the dark matter profile (see, e.g., Ref. [66]) as well as background estimations. For example, the constraint is weakened by a factor of $O(10^{2-3})$ when we take a cored dark matter profile in our Galaxy. To date, the least uncertain constraint is set by the continuum gamma-ray search from dwarf spheroidal galaxies by Fermi-LAT [67], which has excluded wino dark matter with a mass below about 400 GeV and in 2.2–2.5 TeV.

V. SUMMARY AND DISCUSSION

In this paper, we have investigated the 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, $m_{3/2} > \mathcal{O}(100)$ 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, assuming leptogenesis and that the wino is the lightest supersymmetric particle.

As a final remark, let us interpret the gravitino mass from the landscape point of view [68–71]. 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 this case, nature should choose the gravitino mass that 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 from the landscape point of view.[15]

It should be cautioned that there is a hidden parameter in this argument, k, which has been fixed as $k \simeq 0.01$ 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 mass is lowered to

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

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 $m_{3/2} \sim 1$ TeV, in which the electroweak scale is obtained much more naturally than the case with $k \sim 0.01$ and $m_{3/2} \sim 100$ TeV, and the prediction already contradicts the observed value of the spectral index.

This situation is similar to anthropic arguments [74] regarding the electroweak scale. It is argued that an electroweak scale like that realized in nature is required for people to exist in the Universe [75,76]. Therefore, other parameters besides 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 considered to be fixed by some underlying theory.

[15] 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 [72] by decaying gravitino dark matter [73].

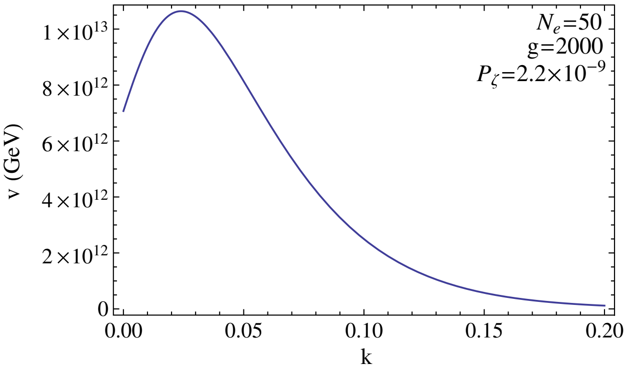

FIG. 7 (color online). A line in $k-v$ space in which $\mathcal{P}_\zeta = 2.2 \times 10^{-9}$.

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 hypersurface of the parameter space corresponding to $\mathcal{P}_\zeta \sim 10^{-9}$ would have been chosen anthropically. In Fig. 7, we show a line in $k-v$ 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 why the spectral index n_s is not too small, such as 0.8, but close to the observed value, i.e., $n_s \sim 0.96$.

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