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## Gravitino problem in supergravity chaotic inflation and SUSY breaking scale after BICEP2

Kazunori Nakayama <sup>a,c,\*</sup>, Fuminobu Takahashi <sup>b,c</sup>, Tsutomu T. Yanagida <sup>c</sup><sup>a</sup> Department of Physics, University of Tokyo, Tokyo 113-0033, Japan<sup>b</sup> Department of Physics, Tohoku University, Sendai 980-8578, Japan<sup>c</sup> Kavli IPMU, TODIAS, University of Tokyo, Kashiwa 277-8583, Japan

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## ABSTRACT

Gravitinos are generically produced by inflaton decays, which place tight constraints on inflation models as well as supersymmetry breaking scale. We revisit the gravitino production from decays of the inflaton and the supersymmetry breaking field, based on a chaotic inflation model suggested by the recent BICEP2 result. We study cosmological constraints on thermally and non-thermally produced gravitinos for a wide range of the gravitino mass, and show that there are only three allowed regions of the gravitino mass:  $m_{3/2} \lesssim 16$  eV,  $m_{3/2} \simeq 10$ –1000 TeV and  $m_{3/2} \gtrsim 10^{13}$  GeV.

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## 1. Introduction

Recently the BICEP2 Collaboration reported a detection of the primordial B-mode polarization of the cosmic microwave background (CMB) [1], which, if confirmed, would provide the strong case for inflation [2,3]. The BICEP2 result can be explained by a large tensor-to-scalar ratio,  $r = 0.20_{-0.05}^{+0.07}$ . Taken at face value, it implies large-field inflation models, where the inflaton field excursion exceeds the Planck scale [4].

Among various large-field inflation models, by far the simplest one is the quadratic chaotic inflation model [5] given by

$$V(\varphi) = \frac{1}{2}m^2\varphi^2, \quad (1)$$

where  $\varphi$  is the inflaton, and the inflaton mass is fixed to be  $m \simeq 2 \times 10^{13}$  GeV by the observed curvature perturbations. The energy density of the Universe during inflation is close to the GUT scale, and therefore, it is conceivable that the inflation model (1) is realized in the framework of supergravity or string theory. The chaotic inflation model in supergravity was proposed in Refs. [6,7], where an approximate shift symmetry on the inflaton was introduced to have good control over inflaton field values greater than the Planck scale.<sup>1</sup>

\* Corresponding author.

E-mail addresses: [kazunori@hep-th.phys.s.u-tokyo.ac.jp](mailto:kazunori@hep-th.phys.s.u-tokyo.ac.jp) (K. Nakayama), [fumi@tuhep.phys.tohoku.ac.jp](mailto:fumi@tuhep.phys.tohoku.ac.jp) (F. Takahashi), [tsutomu.yanagida@ipmu.jp](mailto:tsutomu.yanagida@ipmu.jp) (T.T. Yanagida).

<sup>1</sup> There are various large-field inflation models in the supergravity and superstring theory [8–29].

In order to lead to the standard Big Bang cosmology after inflation, the inflaton must transfer its energy to the standard model (SM) particles, i.e., reheating of the Universe. In supergravity, the inflaton generically decays through various Planck-suppressed interactions, unless the inflaton is charged under unbroken symmetry. Specifically, the inflaton decays into top quarks and Higgs, gluon pairs, right-handed neutrinos, etc. even without introducing ad hoc couplings with the visible sector [30–32]. Some unwanted relics, however, are also produced at the same time. One of such unwanted relics is the gravitino. In fact, it is known that gravitinos are generically produced by decays of the inflaton [31,33–35] and the moduli [36–38], and the solutions to the gravitino overproduction problem were studied in Refs. [39–41].

The amount of gravitinos produced by the inflaton decay depends on the properties of the inflaton and the supersymmetry (SUSY) breaking field. The gravitino production rate is enhanced for a heavier inflaton with a larger coefficient of the linear term in the Kähler potential. In many inflation models, the latter approximately coincides with the vacuum expectation value (VEV) of the inflaton or waterfall fields after inflation. The gravitino overproduction problem becomes acute especially if the SUSY breaking field,  $z$ , is a purely singlet, i.e., the so called Polonyi field, as in the gravity mediation [33,37,38]. In this case the inflaton decay produces too many gravitinos, and as a result, various inflation models are tightly constrained or excluded for a wide range of the gravitino mass [31,33–35].

Moreover, the Polonyi field itself causes a serious cosmological problem [42]; the coherent oscillations of the Polonyi field easily dominate the energy density of the Universe, and typically decay into the SM particles during and after the Big-Bang

nucleosynthesis (BBN), altering the light element abundances in contradiction with observations, and they also produce too many lightest SUSY particles (LSPs). It is possible to consider dynamical SUSY breaking scenarios in which  $z$  is charged under some symmetry and is stabilized with a heavy SUSY breaking mass [41–44]. Then, the Polonyi problem becomes significantly relaxed since the Polonyi field can be stabilized at the enhanced symmetry point during inflation, suppressing the initial oscillation amplitude. Also, the inflaton decay rate into gravitinos and the Polonyi fields can be suppressed [41].

In this letter we revisit the gravitino overproduction problem in the chaotic inflation in light of the recent BICEP2 data, for a wide range of the gravitino mass. In particular we take account of various sources for the gravitino production; thermal production as well as non-thermal one from decays of the inflaton and the Polonyi field. We will show that there are only three allowed regions of the gravitino mass,  $m_{3/2} \lesssim 16$  eV,  $m_{3/2} \simeq 10$ –1000 TeV and  $m_{3/2} \gtrsim 10^{13}$  GeV.

## 2. Chaotic inflation in supergravity

In this section we briefly review the chaotic inflation model given in Refs. [6,7]. We impose a shift symmetry on the inflaton superfield,

$$\phi \rightarrow \phi + iC, \quad (2)$$

where  $C$  is a real transformation parameter. With an abuse of notation, we shall use the same symbol to denote both a chiral superfield and its scalar component, unless noted otherwise. The relevant interactions are given by

$$K_{\text{inf}} = c(\phi + \phi^\dagger) + \frac{1}{2}(\phi + \phi^\dagger)^2 + |X|^2 - k|X|^4 + \dots, \quad (3)$$

$$W_{\text{inf}} = mX\phi, \quad (4)$$

where  $c$  and  $k$  are real constants of order unity and  $X$  is a gauge singlet chiral superfield, and we assume that  $X$  has an  $R$  charge  $+2$  whereas  $\phi$  is neutral. Here we have chosen the origin of  $\phi$  so that the superpotential takes the above form. The inflaton is identified with  $\varphi \equiv \sqrt{2}\text{Im}\phi$ . The Kähler potential respects the shift symmetry (2), which is explicitly broken by the superpotential. Here and in what follows, we adopt the Planck unit in which the reduced Planck mass is set to be unity,  $M_P = 1$ .

The scalar potential in supergravity is given by

$$V_{\text{sugra}} = e^K ((D_i W) K^{i\bar{j}} (D_{\bar{j}} W)^* - 3|W|^2). \quad (5)$$

The inflaton potential is given by (1) even for  $\varphi \gg 1$  because of the shift symmetry. Note that, during inflation,  $X$  is stabilized at the origin  $X = 0$  for  $k \gtrsim \mathcal{O}(1)$ , whereas the real component of  $\phi$  is stabilized at  $\text{Re}\phi \approx -c/2$ . After inflation, both  $X$  and  $\phi$  are stabilized at the origin. For the graceful exit of the inflation we assume  $c$  is at most of order unity [7]. Note that the real component of  $\phi$  starts to oscillate after inflation with an initial amplitude  $\sim c/2$ . For  $c = \mathcal{O}(1)$ , its abundance is comparable to that of the inflaton  $\varphi$ .

One can impose a discrete  $Z_2$  symmetry under which both  $\phi$  and  $X$  flip their sign. Then those terms proportional to odd powers of  $(\phi + \phi^\dagger)$  in the Kähler potential are absent, and in particular,  $c = 0$ . In this case, the inflaton decay into gravitinos can be forbidden. On the other hand, it becomes non-trivial to induce the inflaton decay. As long as the  $Z_2$  symmetry is unbroken, the inflaton would be stable unless we assign the  $Z_2$  charge on the SM particles and its SUSY partners. Otherwise, we would need to break the  $Z_2$  symmetry to induce the inflaton decay. Therefore, introducing a  $Z_2$  symmetry imposes extra conditions on the structure of the underlying theory. Throughout this letter we assume that there

is no such  $Z_2$  symmetry so that there is a linear term in the Kähler potential with  $c = \mathcal{O}(1)$ .

The inflaton decay into the visible sector in the chaotic inflation model was studied in detail in Ref. [32]. In the presence of the linear term in the Kähler potential, the inflaton is automatically coupled to all the fields that appear in the superpotential with gravitational strength. For instance, the inflaton can decay into top (s)quarks and Higgs(ino) at tree level [30], and a pair of gauge bosons and gauginos at one-loop level [31]. Also the inflaton decays into a pair of the right-handed (s)neutrinos if kinematically allowed. The inflaton decay rate into the visible sector is of order  $m^3$ , and the reheating temperature is  $\mathcal{O}(10^9)$  GeV for  $c = \mathcal{O}(1)$ , without introducing ad hoc couplings with the visible sector. This is one of the virtues of the chaotic inflation without the  $Z_2$  symmetry. Therefore we will adopt the inflaton decay rate  $\Gamma_{\text{inf}} \sim 1$  GeV and the reheating temperature  $T_R \sim 10^9$  GeV as reference values in the following analysis.

Lastly let us comment on the mass eigenstates of the inflaton sector after inflation. We will add a constant term  $W_0 \simeq m_{3/2}$  in the superpotential in order to cancel positive contributions from the SUSY breaking. Then,  $\phi$  and  $X$  get maximally mixed with each other to form mass eigenstates [33],

$$\Phi_{\pm} \equiv \frac{\phi \pm X^\dagger}{\sqrt{2}}. \quad (6)$$

Both eigenstates have a mass  $\simeq m$ . This mixing is effective only if the inflaton decay rate is smaller than the gravitino mass. For our reference value of the inflaton decay rate, the critical value is  $m_{3/2} \sim 1$  GeV. For  $m_{3/2} \gtrsim \mathcal{O}(1)$  GeV, the inflaton can decay through the interactions of  $\phi$  and  $X$  with the other sectors, and as we shall see in the next section, the mixing between  $\phi$  and  $X$  is crucial for non-thermal gravitino production. On the other hand, for  $m_{3/2} \lesssim \mathcal{O}(1)$  GeV,  $\phi$  is an effective mass eigenstate until the decay, the mixing between  $\phi$  and  $X$  is irrelevant for the inflaton decay. In this case, although the direct gravitino production from the inflaton becomes ineffective, too many gravitinos are thermally produced for  $T_R \sim 10^9$  GeV except  $m_{3/2} \lesssim 16$  eV [50]. In the next section we will study the gravitino production from various sources in detail.

## 3. Gravitino production

Gravitinos are produced from various sources. First, gravitinos are generically produced non-thermally by inflaton decays. As we shall see shortly, the gravitino production rate sensitively depends on the properties of the SUSY breaking field. Secondly, gravitinos are produced by thermal scatterings in plasmas. Thirdly, the Polonyi field mainly decays into a pair of gravitinos, if kinematically allowed. The Polonyi field is copiously produced by coherent oscillations, and it is also produced by the inflaton decays. We will consider these production processes alternately.

### 3.1. Non-thermal production from inflaton decays

The gravitino production from inflaton decays proceeds through couplings between the inflaton and SUSY breaking field(s). In the following we assume  $m > 2m_{3/2}$ . Otherwise gravitinos are not produced by the inflaton decay.

Let us add a simple extension of the Polonyi model to the inflation model (3):

$$K = K_{\text{inf}} + |z|^2 - \frac{|z|^4}{\Lambda^2},$$

$$W = W_{\text{inf}} + \mu^2 z + W_0, \quad (7)$$

where  $z$  is the SUSY breaking field,  $\Lambda$  is an effective cut-off scale for the quartic coupling, and  $W_0$  is a constant term that is required to make the cosmological constant (almost) zero in the present Universe. The linear term of  $z$  in the superpotential could be generated dynamically, and the quartic coupling in the Kähler potential may or may not be induced by the same dynamics. We however do not specify the origin of  $\mu^2$  to retain generality, and indeed, the following discussion does not depend on it.<sup>2</sup> For simplicity we take  $\mu^2$  and  $W_0$  real in the following.

Let us first consider the original Polonyi model which is obtained in the limit of  $\Lambda \rightarrow \infty$ . In this case  $z$  is a purely singlet, and there is no special point in the field space. One can show that  $z$  is stabilized at  $\langle z \rangle = \sqrt{3} - 1$  with the  $F$ -term,  $F^z = -\sqrt{3}\mu^2$ . In the original Polonyi model, there is the notorious cosmological Polonyi problem [42]. The Polonyi field is copiously produced by coherent oscillations, and easily dominates the Universe after reheating. For the gravitino mass comparable to or lighter than  $\mathcal{O}(1)$  TeV, it decays during or after BBN, altering the light element abundances in contradiction with observations. In fact, even if the Polonyi problem is solved by some mechanism,<sup>3</sup> there is the gravitino overproduction from the inflaton decay. As the SUSY breaking field  $z$  is a purely singlet, the following operator is allowed [33,37, 38],

$$K_{\text{int}} = \frac{1}{2}(\phi + \phi^\dagger)zz + \text{h.c.} \quad (8)$$

This leads to the coupling of the inflaton to goldstinos  $\chi$ ,

$$\mathcal{L} = mX\chi\chi + \text{h.c.} = m\frac{\phi^\dagger - \phi}{\sqrt{2}}\chi\chi + \text{h.c.} \quad (9)$$

where we have expanded the  $F$ -term of the inflaton in terms of  $X$  and used  $m \approx W_{\phi\chi}$ , and  $\chi$  is the fermionic component of  $z$ . Therefore, the inflaton mass eigenstates  $\Phi_\pm$  decay into a pair of gravitinos with a rate of order  $m^3$ , which results in the gravitino overproduction for a wide range of the gravitino mass. Here and in what follows, we assume that the mixing between  $\phi$  and  $X$  is effective as long as the non-thermal gravitino production from inflaton decay is concerned.

The Polonyi problem can be ameliorated if the SUSY breaking field  $z$  is charged under some symmetry as in the dynamical SUSY breaking scenario [43]. The low-energy effective theory is given by Eq. (7) with the cut-off scale  $\Lambda \ll 1$ . The  $F$ -term of  $z$  is given by  $F_z \approx -\mu^2 = \sqrt{3}m_{3/2}$ , and the mass is given by

$$m_z^2 \simeq \frac{12m_{3/2}^2}{\Lambda^2}. \quad (10)$$

The low energy true minimum is located at

$$\langle z \rangle \simeq 2\sqrt{3}\left(\frac{m_{3/2}}{m_z}\right)^2 \simeq \frac{m_{3/2}}{m_z}\Lambda.$$

During inflation the  $z$  can be stabilized in the vicinity of the origin where the symmetry is enhanced, suppressing the initial oscillation amplitude. Still, some amount of coherent oscillations of  $z$  is induced because the low-energy minimum of  $z$  is slightly deviated from the origin. We will return to the (relaxed) Polonyi problem later in this section.

Let us estimate the gravitino production from inflation decays in the case of  $\Lambda \ll 1$ . Since  $z$  is charged under some symmetry, interactions like (8) and  $K_{\text{int}} = (\phi + \phi^\dagger)z + \text{h.c.}$  are forbidden. Note

<sup>2</sup> Actually the situation gets worse if  $\mu^2$  is of dynamical origin and if its dynamical scale is below the inflaton mass, as the inflaton directly decays into hidden gauge sectors, producing many gravitinos [31].

<sup>3</sup> One possibility is the so called adiabatic suppression mechanism [45–47].

however that the effective interactions are induced from higher order terms, since  $z$  develops a small but non-zero VEV. For instance let us consider

$$K = (\phi + \phi^\dagger)\frac{|z|^4}{\Lambda^2}, \quad (11)$$

which leads to the effective interaction (8) with a coefficient suppressed by  $\langle z \rangle^2/\Lambda^2 \sim m_{3/2}^2/m_z^2$ . The contribution to the gravitino pair production rate is therefore significantly suppressed for  $m_z \gg m_{3/2}$ , as expected.

There is another process that becomes relevant for a heavy  $m_z$ , which can be understood as follows. First, the SUSY breaking field  $z$  is known to mainly decay into a pair of gravitinos, as its coupling is enhanced for longitudinal modes (see Eq. (23)). This is an analogue of the standard model Higgs boson which would mainly decay into the longitudinal modes of  $W$  bosons for the Higgs mass heavier than twice the  $W$  boson mass. Secondly, there is a non-zero mixing between  $X$  and  $z$  in supergravity. Thus, the inflaton decays into a pair of gravitinos through the mixing between  $X$  and  $z$ .

To get the feeling of how the decay proceeds, let us write down a part of the interactions leading to the mixing:

$$V_{\text{sugra}} \supset e^K K_\phi W K^{\phi\bar{\phi}} (D_\phi W) + \text{h.c.} \supset m\langle K_\phi \rangle \mu^2 z X^\dagger + \text{h.c.}, \quad (12)$$

where we have expanded  $W$  and  $W_\phi$  with respect to  $z$  and  $X$ , respectively. This induces a mixing between  $z$  and  $X$ , and the mixing angle  $\theta$  depends on the relation between  $m$  and  $m_z$ :

$$\theta \sim \begin{cases} \frac{m_{3/2}}{m} K_\phi & \text{for } m \gg m_z, \\ \frac{m_{3/2}m}{m_z^2} K_\phi & \text{for } m \ll m_z. \end{cases} \quad (13)$$

There is another contribution to the mixing of comparable size from interaction like  $K = (\phi + \phi^\dagger)|z|^2$ . Although not shown here, there is also a contribution from the kinetic mixing,  $K_{\phi z} \sim \langle z \rangle$ . The following results on the gravitino production rate can be understood by combining the above mixing angle and the decay rate of  $z$  into a pair of gravitinos given in Eq. (23). The detailed expressions for the gravitino pair production rate are given in Ref. [38].

Here we summarize the gravitino production rate (cf. Ref. [41]):

$$\Gamma(\Phi \rightarrow 2\psi_{3/2}) \simeq \begin{cases} \frac{1}{8\pi} \frac{m^3}{M_p^2} (\tilde{c}^2 (\frac{m_{3/2}}{m})^2 + c'^2 (\frac{m_{3/2}}{m_z})^4 + \frac{d^2}{8} (\frac{m_z}{m})^4) & \text{for } m \gg m_z, \\ \frac{d^2}{64\pi} \frac{m^3}{M_p^2} & \text{for } m \ll m_z, \end{cases} \quad (14)$$

where we have defined<sup>4</sup>

$$\tilde{c} \equiv \langle K_{\phi z z} \rangle, \quad (15)$$

$$c' \equiv \frac{1}{4} \langle K_{\phi z z z z} \rangle, \quad (16)$$

$$d \equiv \langle K_\phi - K_{\phi z z} \rangle, \quad (17)$$

which are considered to be of order unity. The second term in  $d$  has a comparable contribution to the first term if there is an interaction like  $K \supset (\phi + \phi^\dagger)|z|^2$ .

The inflaton also decays into a pair of the SUSY breaking fields with the rate

<sup>4</sup> Precisely speaking,  $\tilde{c}$  should be defined in the eigenstate basis of the non-analytic mass terms [36]. Our results are not changed by this simplification.

$$\Gamma(\Phi \rightarrow zz^\dagger) = \frac{\tilde{d}^2}{32\pi} \left(\frac{m_z}{m}\right)^4 \frac{m^3}{M_P^2} \left(1 - \frac{4m_z^2}{m^2}\right)^{1/2}$$

for  $m > 2m_z$ ,

where we have defined

$$\tilde{d} \equiv \left\langle K_\phi - K_{\phi\bar{z}\bar{z}} + \frac{K_{z\bar{z}z\bar{z}\phi}}{K_{z\bar{z}z\bar{z}}} \right\rangle. \quad (19)$$

The second and third terms in  $\tilde{d}$  have comparable contributions to the first term if there are operators like  $K \supset (\phi + \phi^\dagger)|z|^2$  and  $(\phi + \phi^\dagger)|z|^4/\Lambda^2$ . Since  $z$  predominantly decays into a pair of gravitinos, this decay process gives a comparable contribution to the final gravitino abundance, with respect to the direct production. We will set  $c' = \tilde{c} = d = \tilde{d} = 1$  in the next section.

From these decay rates, it is obvious that the gravitino production rate is significantly suppressed for a certain range of  $m_z$  satisfying  $m_{3/2} \ll m_z \ll m$ , because of the suppression factor of  $(m_z/m)^4$  and  $(m_{3/2}/m_z)^4$ . The non-thermal gravitino abundance is then given by

$$Y_{3/2}^{(\phi)} = \frac{3T_R}{4m} \frac{2\Gamma(\Phi \rightarrow \bar{z}\bar{z}) + 4\Gamma(\Phi \rightarrow zz^\dagger)}{\Gamma_{\text{tot}}}, \quad (20)$$

where  $\Gamma_{\text{tot}} = (\pi^2 g_{*s}/90)^{1/2} T_R^2/M_P$  is the total decay rate of the inflaton.

### 3.2. Thermal production

Gravitinos are also produced by scatterings of SM particles and their superpartners in thermal bath. The gravitino abundance is proportional to the reheating temperature  $T_R$  [48]:

$$Y_{3/2}^{(\text{th})} \simeq \begin{cases} \min\left[2 \times 10^{-12} \left(1 + \frac{m_{\tilde{g}}^2}{3m_{3/2}^2}\right) \left(\frac{T_R}{10^{10} \text{ GeV}}\right), \frac{0.42}{g_{*s}(T_{3/2})}\right] & \text{for } T_R \gtrsim m_{\text{SUSY}}, \\ 0 & \text{for } T_R \lesssim m_{\text{SUSY}}, \end{cases} \quad (21)$$

where  $m_{\tilde{g}}$  denotes the gluino mass,  $m_{\text{SUSY}}$  is the typical soft SUSY breaking mass, and  $g_{*s}(T_{3/2})$  is the effective degrees of freedom at the gravitino freezeout, if the gravitinos are thermalized [49,50]. We will take  $m_{\tilde{g}} = m_{\text{SUSY}}$  in the next section.

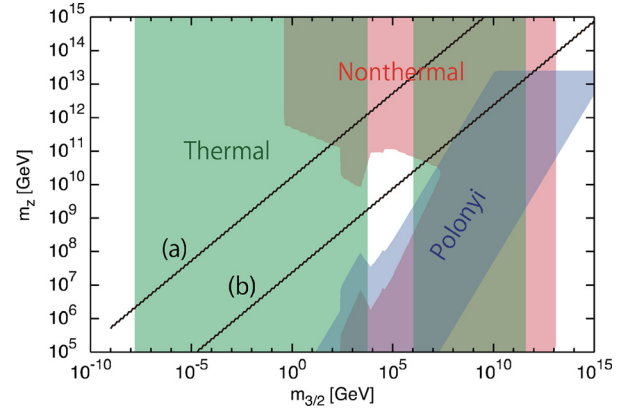
### 3.3. Production from decays of the Polonyi field

The true minimum of the Polonyi field  $z$  is deviated from the position during inflation, and hence it starts to oscillate when the Hubble parameter becomes equal to the Polonyi mass.

If  $m_z \ll m$ , the Polonyi can be stabilized in the vicinity of the origin  $z \sim 0$  due to the Hubble-induced mass term. The low energy true minimum, on the other hand, is located at  $\langle z \rangle = 2\sqrt{3}m_{3/2}^2/m_z^2$ . Thus it starts to oscillate with an amplitude of  $\langle z \rangle$  when the Hubble parameter becomes comparable to  $m_z$ . On the other hand, the coherent oscillations are not induced if  $m_z \gg m$ ,<sup>5</sup> since the Polonyi field adiabatically follows the temporal minimum of the potential in this case [47]. Therefore, the abundance of Polonyi coherent oscillations is estimated as

$$\frac{\rho_z}{s} \simeq \begin{cases} 3T_R \left(\frac{m_{3/2}}{m_z}\right)^4 & \text{for } m_z \ll m, \\ 0 & \text{for } m_z \gg m. \end{cases} \quad (22)$$

The Polonyi field decays into gravitinos with a rate:



**Fig. 1.** The green shaded regions, purple shaded region and blue shaded region are excluded from thermal production ( $Y_{3/2}^{(\text{th})}$ ), non-thermal production by the inflaton decay ( $Y_{3/2}^{(\phi)}$ ) and the Polonyi decay ( $Y_{3/2}^{(z)}$ ), respectively. The two black lines show typical values of  $m_z$  (see text). In this figure we have taken  $T_R = 10^9$  GeV. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

$$\Gamma(z \rightarrow 2\psi_{3/2}) \simeq \frac{1}{96\pi} \frac{m_z^5}{m_{3/2}^2 M_P^2}. \quad (23)$$

The gravitino abundance from the Polonyi decay is estimated as

$$Y_{3/2}^{(z)} \simeq \frac{2}{m_z} \frac{\rho_z}{s}. \quad (24)$$

We will take into account those contributions to the gravitino abundance in the next section.

## 4. SUSY breaking scale inferred from BICEP2

The gravitino abundance is given by the sum of all the contributions considered in the previous section:

$$Y_{3/2} = Y_{3/2}^{(\text{th})} + Y_{3/2}^{(\phi)} + Y_{3/2}^{(z)}. \quad (25)$$

Cosmological effects of the gravitino depend on its mass. If it is stable, the energy density of the gravitino should not exceed the observed dark matter abundance. If it is unstable and if its lifetime is longer than  $\sim 1$  s, the gravitino abundance is severely constrained by BBN. If its lifetime is much shorter than 1 s, the LSPs produced by the gravitino decay is constrained by the dark matter abundance. For concreteness, we assume the anomaly mediation for the mass spectrum of the gauginos for  $m_{3/2} > 100$  TeV and hence the Wino is the LSP.<sup>6</sup> Also we assume the gauge mediation model for  $m_{3/2} \lesssim 100$  GeV in which the NLSP mass is around 1 TeV, so that the decay of NLSP is severely constrained by BBN if  $m_{3/2} \gtrsim 1$  GeV.

Fig. 1 shows the constraints from the gravitino overproduction on the  $(m_{3/2}, m_z)$  plane. In this figure we have taken  $T_R = 10^9$  GeV, for which thermal leptogenesis is possible [51]. Note that the spontaneous inflaton decay through the top Yukawa coupling leads to  $T_R \sim 10^8$  GeV. This is considered to be the lower bound on the reheating temperature [30]. The green, purple and blue shaded regions are excluded by thermal production ( $Y_{3/2}^{(\text{th})}$ ), non-thermal production by the inflaton decay ( $Y_{3/2}^{(\phi)}$ ), and the Polonyi decay ( $Y_{3/2}^{(z)}$ ), respectively. We consider non-thermal production of gravitinos from inflaton decays only for  $m_{3/2} \geq 1$  GeV, since otherwise

<sup>5</sup>  $m$  is close to the Hubble scale at the end of inflation in the chaotic inflation model.

<sup>6</sup> We assume the conservation of R-parity. If the R-parity is violated, the LSP can decay well before BBN and constraints can be significantly relaxed for  $m_{3/2} \gg 10$  TeV.



the mixing between  $\phi$  and  $X$  becomes ineffective.<sup>7</sup> The two black lines show typical values of  $m_z$ . If the mass of  $z$  is generated radiatively,  $\Lambda$  in (3) is written as  $\Lambda = (4\pi/\lambda^2)\Lambda_{\text{dyn}}$  with  $\Lambda_{\text{dyn}}$  being the dynamical SUSY breaking scale, and  $\lambda$  the coupling constant between  $z$  and the strong dynamical sector.<sup>8</sup> The upper line (a) corresponds to the strongly coupled case,  $\lambda = 4\pi$  and the lower line (b) corresponds to the weak coupling case  $\lambda = 1$ .

One can see clearly from the figure that there are only three allowed regions: (i)  $m_{3/2} \lesssim 16$  eV, (ii)  $m_{3/2} \sim 10\text{--}1000$  TeV and (iii)  $m_{3/2} \gtrsim 10^{13}$  GeV. Interestingly, the allowed region (ii) is consistent with the pure gravity mediation scenario [53] / the minimal split SUSY [54,55], which naturally explains the 125 GeV Higgs boson observed at LHC [56]. If the SUSY breaking sector is truly strongly coupled, i.e.  $\lambda = 4\pi$  corresponding to the line labeled by (a), there are only two regions (i) and (iii). The second region (ii) is viable if the coupling of the SUSY breaking field  $z$  is perturbative, e.g.  $\lambda = 1$  (the line labeled by (b)).

## 5. Conclusions

In this paper we revisited the gravitino overproduction problem in the chaotic inflationary Universe scenario, in light of recent BICEP2 result. Taking account of the non-thermal gravitino production from the direct inflaton decay as well as thermal production, and also the effect of Polonyi coherent oscillation, we have shown that there are only three allowed regions of the gravitino mass:  $m_{3/2} \lesssim 16$  eV,  $m_{3/2} \sim 10\text{--}1000$  TeV and  $m_{3/2} \gtrsim 10^{13}$  GeV. It is interesting that, except for the trivial limits of ultra light and ultra heavy gravitino, the gravitino mass of  $\sim 100$  TeV appeared from these considerations, which fits the pure gravity mediation scenario. Interestingly, the inflaton decays into the visible sector even without introducing ad hoc couplings, because there is generically a linear term of the inflaton in the Kähler potential. Therefore the inflaton generically decays into all the fields that appear in the superpotential, and the reheating temperature is naturally as high as  $\sim 10^9$  GeV so that thermal leptogenesis successfully works. The non-thermal leptogenesis is also possible, if the right-handed neutrino mass is close to the inflaton mass [32].

The large tensor-to-scalar ratio observed by BICEP2 indicates a detectable level of stochastic gravitational wave background of the primordial origin around the frequency of  $\sim 1$  Hz, which can be detected by future space-based gravitational wave detectors. In particular, the observation of the shape of the gravitational wave spectrum enables us to determine the reheating temperature, if it is around  $\sim 10^9$  GeV [57,58].

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<sup>7</sup> Still, there are other production processes which are not suppressed for  $m_{3/2} \lesssim 1$  GeV. For instance, the abundance of the real component of  $\phi$  is sizable and it decays into the lowest components of  $z$  with a rate comparable to (18). This however does not change our conclusion, because thermally produced gravitinos exceed the observed dark matter abundance in this case for  $T_R \sim 10^9$  GeV.

<sup>8</sup>  $\lambda$  corresponds to the coupling between  $z$  and dynamical quarks in the IYIT model [52] (see Ref. [41]).

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