

Natural Gravitino Dark Matter and Thermal Leptogenesis in Gauge-Mediated Supersymmetry-Breaking Models

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Abstract

We point out that there is no cosmological gravitino problem in a certain class of gauge-mediated supersymmetry-breaking (GMSB) models. The constant term in the superpotential naturally causes small mixings between the standard-model and messenger fields, which give rise to late-time decays of the lightest messenger fields. This decay provides an exquisite amount of entropy, which dilutes the thermal relics of the gravitinos down to just the observed mass density of the dark matter. This remarkable phenomenon takes place automatically, irrespective of the gravitino mass and the reheating temperature of inflation, once the gravitinos and messenger fields are thermalized in the early Universe. In this class of GMSB models, there is no strict upper bound on the reheating temperature of inflation, which makes the standard thermal leptogenesis the most attractive candidate for the origin of the observed baryon asymmetry in the present Universe.

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1 Introduction

The minimal supersymmetric standard model (MSSM) is the most promising candidate for physics beyond the Standard Model (SM), since it naturally solves the “hierarchy problem” and leads to a successful unification of the gauge coupling constants [1]. Because SUSY is not observed in the real world, it should be broken around the TeV scale. Once we allow generic soft SUSY-breaking terms, we must face hundreds of new parameters, which makes the rate of flavour-changing neutral-current (FCNC) interactions many orders of magnitude larger than the present experimental bounds. In order to obtain a successful low-energy effective theory, various mediation mechanisms of SUSY-breaking effects have been proposed.

The scenario most commonly considered in phenomenology is the minimal gravity-mediated SUSY-breaking (mSUGRA) models. This scenario is very simple and aesthetically attractive, since gravity does exist in nature. In the mSUGRA models, we also have a promising candidate for dark matter. That is the lightest SUSY particle (LSP), which is usually the lightest neutralino. The severest difficulty in this scenario is the lack of a natural explanation for the suppression of the FCNC interactions. A specific form of the soft SUSY-breaking masses must be imposed by hand in order to suppress the FCNC interactions.

It has been argued that the suppression of the FCNC interactions can be naturally obtained in brane-world SUSY-breaking scenarios, such as anomaly- [2] and gaugino- [3] mediated SUSY-breaking models. In these scenarios, the fields relevant to SUSY breaking are assumed to reside on the hidden brane, which is geometrically separated from the visible brane where the SM fields are localized. Recently, however, a crucial observation has been made in Ref. [4], where the authors have found that the separation of the visible and hidden branes in a higher-dimensional space-time is not sufficient for suppressing the FCNC interactions. Consequently, we need additional *ad hoc* assumptions in those models.¹

At present, the most attractive scenario seems to be the gauge-mediated SUSY breaking (GMSB) [6]. In the GMSB models, the suppression of the FCNC interactions is realized in an automatic way, just because SUSY breaking occurs at a very low-energy scale. Furthermore, a whole spectrum of the superparticles in the MSSM sector is com-

¹There is an attempt to realize anomaly-mediation models in a four-dimensional framework [5].

pletely determined by only a few parameters, which allows us to discriminate the GMSB models from other candidates in the future collider experiments.

Unfortunately, from a cosmological perspective, there exists a big difficulty in the GMSB models. That is the so-called cosmological gravitino problem. In these models, we have no natural explanation for the dark matter in the present Universe. Thermal relics of the gravitino, the LSP in the GMSB models, overclose the Universe once they are thermalized in the early Universe. In order to avoid the overproduction of the gravitinos, there is a severe upper bound on the reheating temperature of inflation T_R , which is about $T_R \lesssim 10^6$ GeV for $m_{3/2} = 10$ MeV for instance, and it even reaches $T_R \lesssim 10^3$ GeV in the case of the lighter gravitinos $m_{3/2} \lesssim 100$ keV [7]. Furthermore, we have to fine-tune the reheating temperature just below this upper bound to explain the required mass density of the dark matter. We also have to generate the observed baryon asymmetry at just the same reheating temperature. Therefore, for successful cosmology in the GMSB models, we need incredible fine-tunings of various parameters, which apparently belong to independent physics, such as SUSY breaking, inflation and baryo/leptogenesis.² This is a big drawback of the GMSB models with respect to the standard mSUGRA scenario.

In this letter, we point out that there are indeed no such difficulties in a certain class of GMSB models. We consider direct gauge-mediation models where the SUSY-breaking effects are directly transmitted to the messenger sector without loop suppressions [9]. In this class of models, the specific form of the superpotential of the messenger sector is usually provided by the R -symmetry: since it is violated by the constant term in the superpotential $\langle W \rangle$, which is anyhow required to cancel the cosmological constant, it is quite natural to expect that there are small mixings between the SM and messenger multiplets induced by the condensation $\langle W \rangle$. As a result, the lightest messenger particle decays into the SM particle and gaugino through the mixings. As we will see, the resultant late-time decays of the lightest messengers provide an exquisite amount of entropy, which dilutes the thermal relics of the gravitinos down to just the observed mass density of the dark matter in the present Universe.

Surprisingly, this miracle turns out to be true almost irrespective of the mass of the gravitino and the reheating temperature, once the gravitinos and messenger particles are

²If there exist extra matter multiplets of a SUSY-invariant mass of the order of the “ μ -term”, the observed baryon asymmetry and gravitino dark matter can be simultaneously explained in a way totally independent of the reheating temperature [8].

thermalized in the early Universe. Consequently, by assuming natural mixings between the SM and messenger fields, the severe upper bound on the reheating temperature is completely eluded. This fact makes it much easier to construct a realistic inflationary scenario in the GMSB models. This result also has an important implication on the origin of the observed baryon asymmetry. We will see that the standard thermal leptogenesis, through out-of-equilibrium decays of the right-handed Majorana neutrinos [10], is now the most attractive mechanism to generate the observed baryon asymmetry in the GMSB models.

2 Required Amount of Entropy

In the GMSB models, the longitudinal component of the gravitino (\sim Goldstino) ψ interacts fairly strongly with the SM particles. The total production cross section of ψ by scattering processes is given by [7]

$$\langle \Sigma_{\text{scatt}} v_{\text{rel}} \rangle \approx 5.9 \frac{g_3^2 m_{\tilde{G}}^2}{m_{3/2}^2 M_*^2}, \quad (1)$$

where $\langle \rangle$ denotes the thermal average; $M_* = 2.4 \times 10^{18}$ GeV is the reduced Planck scale, $m_{\tilde{G}}$ is the mass of the gluino and g_3 is the coupling constant of the $SU(3)_C$ gauge group. Then, the resultant interaction rate is given by

$$\Gamma_{\text{scatt}} \approx \langle \Sigma_{\text{scatt}} v_{\text{rel}} \rangle n_{\text{rad}}, \quad (2)$$

with $n_{\text{rad}} = (\zeta(3)/\pi^2) T^3$ being the number density for one massless degree of freedom.

By the scattering interactions, the gravitinos are kept in thermal equilibrium if $\Gamma_{\text{scatt}}/H \gtrsim 1$, where H is the Hubble parameter of the expanding Universe. The corresponding freeze-out temperature of ψ is estimated to

$$T_f \approx 1 \text{ TeV} \left(\frac{g_*(T_f)}{230} \right)^{1/2} \left(\frac{m_{3/2}}{10 \text{ keV}} \right)^2 \left(\frac{1 \text{ TeV}}{m_{\tilde{G}}} \right)^2, \quad (3)$$

where $g_*(T_f)$ denotes the effective massless degrees of freedom when the cosmic temperature $T = T_f$.³ If the reheating temperature of inflation T_R is higher than T_f ,

³For the light gravitino $m_{3/2} \lesssim 10$ keV, the decay processes of SUSY particles are comparable with scattering ones and keep the gravitinos in thermal equilibrium until the temperature drops below the superparticle-mass scale.

gravitinos are in thermal equilibrium in the early Universe. Here, we have assumed the radiation-dominated Universe at the freeze-out time of the gravitinos. We will justify this assumption later, even in the presence of the messenger particles.

If there is no additional entropy production, the resultant yield of the thermal gravitino is estimated to

$$Y_{3/2} \left(\equiv \frac{n_{3/2}}{s} \right) = \frac{45}{2\pi^2 g_*(T_f)} \frac{\zeta(3)}{\pi^2} \left(\frac{3}{2} \right), \quad (4)$$

where $n_{3/2}$ is the number density of gravitinos and s is the entropy density. In terms of the density parameter, it is written as

$$\Omega_{3/2} h^2 \simeq 5.0 \times \left(\frac{m_{3/2}}{10 \text{ keV}} \right) \left(\frac{230}{g_*(T_f)} \right), \quad (5)$$

where h is the present Hubble parameter in units of $100 \text{ km sec}^{-1} \text{ Mpc}^{-1}$, and $\Omega_{3/2} \equiv \rho_{3/2}/\rho_c$; $\rho_{3/2}$ and ρ_c are the energy density of the gravitino and the critical density in the present Universe, respectively. Since the observed mass density of the dark matter is $\Omega_{\text{DM}} h^2 \simeq 0.1\text{--}0.2$, the required dilution factor via the late-time entropy production is given by

$$\Delta \simeq 33 \times \left(\frac{m_{3/2}}{10 \text{ keV}} \right) \left(\frac{230}{g_*(T_f)} \right) \left(\frac{0.15}{\Omega_{\text{DM}} h^2} \right). \quad (6)$$

The above entropy should be supplied after the freeze-out time of the gravitinos.

3 Decay of the Lightest Messenger and $\Omega_{3/2}$

In this section, we discuss the decay of the lightest messenger and the resultant mass density of gravitino dark matter. Further constraints will be discussed in the next section. In this letter, we adopt the simplest messenger sector, which consists of a pair of chiral supermultiplets $\Phi + \bar{\Phi}$, describing a Dirac fermion of mass M and two complex scalar fields of mass squared $M^2 \pm F$, where F is the F -term SUSY-breaking component of the mass of the messenger multiplets. In order to preserve the success of the gauge-coupling unification, we assume that $\Phi + \bar{\Phi}$ transform as $\mathbf{5} + \bar{\mathbf{5}}$ under the $\text{SU}(5)_{\text{GUT}}$ gauge group.⁴

Under this setup, the gauginos and SUSY particles in the MSSM sector obtain the following soft SUSY-breaking masses, M_a and m_{soft}^2 , via gauge interactions at the one-

⁴Adopting a pair of $\mathbf{10} + \bar{\mathbf{10}}$ messenger multiplets does not change the basic results in this letter.

and two-loop level, respectively:

$$M_a \simeq \frac{g_a^2}{16\pi^2} \Lambda, \quad m_{\text{soft}}^2 \simeq 2 \sum_a \left\{ C_a \left(\frac{g_a^2}{16\pi^2} \right)^2 \right\} \Lambda^2, \quad (7)$$

where g_a ($a = 1, 2, 3$) are gauge coupling constants and C_a are the quadratic Casimir. Here, $\Lambda \equiv F/M$ determines the overall scale of the soft SUSY-breaking masses; $\Lambda \simeq 10^5$ GeV is required to obtain the correct size of soft masses. The mass of the messenger particles M and the gravitino mass $m_{3/2}$ are related by the formula:

$$M = \frac{\sqrt{3} k M_*}{\Lambda} m_{3/2} \simeq 4.2 \times 10^8 k \left(\frac{m_{3/2}}{10 \text{ keV}} \right) \left(\frac{10^5 \text{ GeV}}{\Lambda} \right) \text{ GeV}. \quad (8)$$

Here, we have defined $k \equiv F/F_{\text{DSB}}$, where F_{DSB} denotes the original F -term in the dynamical SUSY-breaking sector. In direct gauge-mediation models, this ratio is of the order of 1, i.e. $k \lesssim 1$.

In the present work, we assume that there is the following mixing term between the SM and messenger multiplets through the small R -symmetry-breaking effects caused by the constant term in the superpotential, $\langle W \rangle$:⁵

$$\delta W = f \frac{\langle W \rangle}{M_*^2} \mathbf{5}_M \bar{\mathbf{5}} = f m_{3/2} \mathbf{5}_M \bar{\mathbf{5}}, \quad (9)$$

where f denotes some unknown coefficient of the order of 1 and the subscript M denotes the multiplet that belongs to the messenger sector. In the present scenario, the lightest messenger superparticle is most likely the scalar component of a weak doublet. By virtue of the above small mixing term, the lightest messenger can decay into a SM lepton and a gaugino.⁶ The decay rate is estimated to be

$$\Gamma_M \simeq \frac{g_2^2}{16\pi} \left(\frac{f m_{3/2}}{M} \right)^2 M. \quad (10)$$

The resultant decay temperature of the lightest messenger is given by

$$T_d \simeq 68 \text{ MeV} \times \frac{f}{\sqrt{k}} \left(\frac{10}{g_*(T_d)} \right)^{1/4} \left(\frac{m_{3/2}}{10 \text{ keV}} \right)^{1/2} \left(\frac{\Lambda}{10^5 \text{ GeV}} \right)^{1/2}. \quad (11)$$

⁵Here, we assume, for instance, the R -charge for $\bar{\mathbf{5}}$ and $\mathbf{5}_M$ to be +1 and -1, respectively.

⁶We have neglected the decay channels through small Yukawa interactions.

Here, we have used the relation in Eq. (8). By comparing with Eq. (3), one can see that the decays of the lightest messenger always take place after the freeze-out of the gravitinos. Therefore, the entropy production associated with the decays of the lightest messengers dilutes the thermal relics of the gravitinos.

Now, let us estimate the amount of entropy produced by the late-time decays of the lightest messengers. We will also justify the assumption of the radiation-dominated Universe made in deriving Eq. (3). Assuming the stability of the lightest messengers, their relic density has been estimated in Ref. [11] as

$$Y_M \left(\equiv \frac{n_M}{s} \right) \approx 3.65 \times 10^{-10} \left(\frac{M}{10^6 \text{ GeV}} \right), \quad (12)$$

$$\Omega_M h^2 \approx 10^5 \left(\frac{M}{10^6 \text{ GeV}} \right)^2, \quad (13)$$

where n_M/s denotes the frozen-out value of the yield of the lightest messengers and Ω_M is the corresponding density parameter. After the freeze-out of the lightest messengers, the total energy density of the Universe is given by

$$\rho = \frac{\pi^2}{30} g_*(T) T^4 + \frac{2\pi^2}{45} g_*(T) T^3 M Y_M, \quad (14)$$

where the first and second terms represent the contributions from the radiation and from the lightest messengers, respectively. The thermal relics of the lightest messengers begin to dominate the energy density at

$$T_C = \frac{4}{3} M Y_M \simeq 84 \text{ GeV} \times k^2 \left(\frac{10^5 \text{ GeV}}{\Lambda} \right)^2 \left(\frac{m_{3/2}}{10 \text{ keV}} \right)^2, \quad (15)$$

where we have used the relation in Eq. (8). From Eqs. (3) and (15), one can see that the matter-dominated Universe starts well after the freeze-out time of the gravitinos, and hence the assumption made in the derivation of Eq. (3) is justified.

By assuming the instantaneous decays of the lightest messengers, which is accurate enough for the present purpose, we can easily estimate the dilution factor from energy conservation as

$$\begin{aligned} \Delta_M \left(\equiv \frac{s_{\text{after}}}{s_{\text{before}}} \right) &= \left(1 + \frac{4}{3} Y_M \frac{M}{T_d} \right)^{3/4} \\ &\simeq 1.0 \times 10^2 \left(\frac{M}{10^8 \text{ GeV}} \right)^{3/2} \left(\frac{10 \text{ MeV}}{T_d} \right)^{3/4}, \end{aligned} \quad (16)$$

where s_{before} (s_{after}) denotes the entropy density of the Universe before (after) the decays of the lightest messengers. By substituting Eqs. (8) and (11) into Eq. (16), one can see that Δ_M miraculously coincides the required dilution factor Δ in Eq. (6). From Eqs. (5), (8), (11) and (16), we can derive the final gravitino abundance as follows:

$$\begin{aligned} \Omega_{3/2} h^2 &= \Omega_{3/2} h^2 \Big|_{\text{initial}} \times \frac{1}{\Delta_M} \\ &\simeq 0.18 \times f^{3/4} \left(\frac{10}{g_*(T_d)} \right)^{3/16} \left(\frac{230}{g_*(T_f)} \right) \left(\frac{\Lambda/k}{3 \times 10^5 \text{ GeV}} \right)^{15/8} \left(\frac{10 \text{ keV}}{m_{3/2}} \right)^{1/8}. \end{aligned} \quad (17)$$

Astonishingly, a natural parameter $\Lambda/k \approx 10^5 \text{ GeV}$, which is needed so as to give the correct order of magnitude to the soft SUSY-breaking masses, leads to just the mass density of the gravitinos required to be the dominant dark matter in the present Universe. Furthermore, the resultant mass density of the gravitinos has only negligible dependence on its mass $m_{3/2}$. Consequently, the gravitino dark matter is an almost automatic result of the direct gauge-mediation models with a natural mixing term between the SM and messenger multiplets given in Eq. (9).

Before closing this section, we briefly comment on the property of the resultant gravitino dark matter. The free-streaming length of the gravitino is given by [12]

$$R_f \approx 0.1 \left(\frac{\Omega_{3/2} h^2}{0.15} \right)^{1/3} \left(\frac{1 \text{ keV}}{m_{3/2}} \right)^{4/3} \text{ Mpc}. \quad (18)$$

The late-time entropy production does not change the result in this approximation. For a more accurate estimation, see Ref. [13]. The gravitino of a mass in the range $m_{3/2} \simeq (1\text{--}1.5) \text{ keV}$, whose free-streaming length is about $R_f \simeq 0.1 \text{ Mpc}$, is an interesting warm dark matter candidate,⁷ which may reconcile the predictions of the cold dark matter with observations [15, 13]. Heavier gravitinos with mass $m_{3/2} > (\text{a few}) \text{ keV}$ serve as the cold dark matter. The gravitino with mass $m_{3/2} \lesssim 1 \text{ keV}$ is now disfavoured by observations of Lyman- α forest [16] and the history of cosmological reionization [17].

⁷The possibility of the gravitino warm dark matter with a small entropy production has been discussed in Ref. [14], where the authors introduce the extra matter of mass $M_X \simeq 10^{12} \text{ GeV}$ to obtain the required lifetime of the lightest messenger particle. For their scenario to work, we need an *ad hoc* tuning on M_X . Therefore, their model does not solve the fine-tuning problem for the gravitino dark matter stressed in the introduction of this letter.

4 Further Constraints

In this section, we consider the non-thermal gravitino production from decays of the next-to-lightest superparticles (NLSPs). This contribution may spoil the successful prediction in the previous section, since a huge number of NLSPs are produced in the decays of the lightest messengers below their freeze-out temperature. In addition, if the decay process of the NLSP takes place when $T \lesssim 5$ MeV, it may also spoil the success of the Big Bang nucleosynthesis (BBN) [18].

Let us first consider the constraint from the BBN. The NLSP, which is the bino or stau in the GMSB models, decays into its superpartner and a gravitino, with the following decay width:

$$\Gamma_\chi \simeq \frac{1}{48\pi} \frac{m_\chi^5}{m_{3/2}^2 M_*^2}, \quad (19)$$

where m_χ is the mass of the NLSP. The corresponding decay temperature of the NLSP is given by

$$T_\chi = \left(\frac{90}{\pi^2 g_*(T_\chi)} \right)^{1/4} \sqrt{\Gamma_\chi M_*} \simeq 5 \text{ MeV} \left(\frac{m_\chi}{100 \text{ GeV}} \right)^{5/2} \left(\frac{1 \text{ MeV}}{m_{3/2}} \right). \quad (20)$$

In order not to spoil the success of the BBN, the decays of the NLSPs should be completed before the onset of the BBN, $T_\chi \gtrsim 5$ MeV. If we take the natural range of the mass of the NLSP, for instance $m_\chi \lesssim 250$ GeV, $m_{3/2} \lesssim 10$ MeV is required for the success of the BBN.

Let us now estimate the contribution of the non-thermal gravitinos to the total mass density of the dark matter. If $\Gamma_M > \Gamma_\chi$, the produced NLSPs have enough time to annihilate before they decay into gravitinos. This case corresponds to

$$m_{3/2} \gtrsim 27 \text{ keV} \frac{1}{f^{2/3}} \left(\frac{m_\chi}{100 \text{ GeV}} \right)^{5/3} \left(\frac{3 \times 10^5 \text{ GeV}}{\Lambda/k} \right)^{1/3}. \quad (21)$$

In this case, the resultant yield of the NLSPs before they decay is given by the following value, to a good approximation [19]:

$$\frac{n_\chi}{s} = \sqrt{\frac{45}{8\pi^2 g_*(T_d)} \frac{\langle \sigma v \rangle_\chi^{-1}}{M_* T_d}}, \quad (22)$$

where $\langle \sigma v \rangle_\chi$ denotes the s -wave annihilation cross section of the NLSP. The subsequent decays of the NLSPs produce the same number of gravitinos. Therefore, combined with

Eq. (11), the resultant abundance of the non-thermal gravitinos is given by

$$\Omega_{3/2}^{\text{NT}} h^2 \simeq 1.3 \times 10^{-2} \frac{\sqrt{k}}{f} \left(\frac{10}{g_*(T_d)} \right)^{\frac{1}{4}} \left(\frac{m_{3/2}}{10 \text{ MeV}} \right)^{\frac{1}{2}} \left(\frac{10^5 \text{ GeV}}{\Lambda} \right)^{\frac{1}{2}} \left(\frac{10^{-11} \text{ GeV}^{-2}}{\langle \sigma v \rangle_\chi} \right). \quad (23)$$

In the case of the stau NLSP, the s -wave annihilation cross section is approximately given by $\langle \sigma v \rangle_\chi \simeq 10^{-7} \text{ GeV}^{-2} (100 \text{ GeV}/m_\chi)^2$. Even in the case of the bino NLSP, $\langle \sigma v \rangle_\chi \gtrsim 10^{-11} \text{ GeV}^{-2}$ is naturally obtained for relatively large $\tan\beta$, which is about $\tan\beta \gtrsim 15$ (30) for $m_\chi \simeq 100$ (250) GeV, via small higgsino contamination.

The above estimation is not valid if $\Gamma_M < \Gamma_\chi$, since the produced NLSPs immediately decay into gravitinos. In this case, we have to follow the full evolution by solving coupled Boltzmann equations.⁸ The result is given in Fig. 1. In this calculation, we have conser-

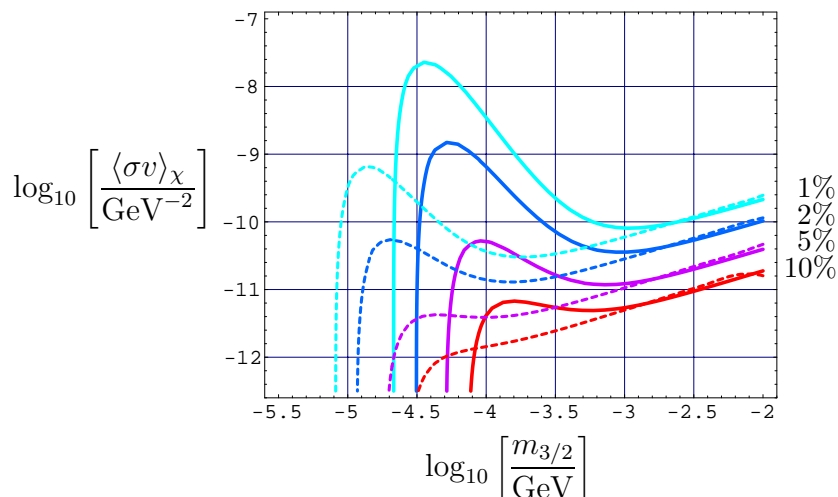


Figure 1: Contour plot of the abundance of the non-thermal gravitinos $\Omega_{3/2}^{\text{NT}} h^2$. Solid (dashed) lines correspond to $\Omega_{3/2}^{\text{NT}} h^2 = 10, 5, 2$ and 1% of the total mass density of the dark matter ($\Omega_{\text{DM}} h^2 \simeq 0.15$), for $m_\chi = 250$ (100) GeV, from the bottom up, respectively. The previous estimation given in Eq. (23) is valid in the region where the contour lines are almost straight.

vatively assumed that the number of NLSPs produced per decay of the lightest messenger particle is given by $N_\chi \approx M/E_{\text{th}}$, where $E_{\text{th}} \approx m_\chi^2/T_d$ is the threshold energy to produce the NLSPs by scatterings with the thermal backgrounds with $T = T_d$. One can see that the contribution of the non-thermal gravitinos is smaller than several per cent of the total mass density of the dark matter as long as $\langle \sigma v \rangle_\chi \gtrsim 10^{-11} \text{ GeV}^{-2}$.

⁸Significant annihilations take place also in this case.

In summary, the prediction of the gravitino dark matter given in Eq. (17) is not spoiled in most of the parameter region. The relevant constraint comes only from the BBN, which is satisfied as long as $m_{3/2} \lesssim 10$ MeV for a reasonable range of m_χ . Finally, we briefly comment on another interesting aspect. The non-thermal gravitinos produced from the NLSPs behave as a hot/warm component of dark matter, which contributes several per cent to the total mass density of dark matter in the case of the bino NLSP with moderate $\tan\beta$. Therefore, the present model can naturally realize a mixed dark matter scenario. By virtue of the late-time production of the NLSPs, an unnaturally large hierarchy between the slepton and the bino masses is not required to realize a mixed dark matter scenario [20].

5 Thermal Leptogenesis

As we have mentioned in the introduction, the present model has important implications on the origin of the baryon asymmetry in the present Universe. By virtue of the entropy production by the decays of the lightest messengers, we obtain the required mass density of the dark matter without any adjustments of the reheating temperature. Consequently, the standard thermal leptogenesis [10, 21] now becomes the most promising candidate for the origin of the observed baryon asymmetry in the GMSB models.

Let us start by introducing the relevant terms in the superpotential:

$$W = \frac{1}{2}M_{Ri}N_iN_i + h_{i\alpha}N_iL_\alpha H_u, \quad (24)$$

where N_i ($i = 1, 2, 3$) denote the heavy right-handed Majorana neutrinos of mass M_{Ri} ; L_α ($\alpha = e, \mu, \tau$) and H_u denote the lepton doublets and the Higgs doublet that couples to up-type quarks, respectively. The lepton-number asymmetry per decay of a right-handed neutrino N_i is given by [10, 22]

$$\begin{aligned} \epsilon_i &\equiv \frac{\sum_\alpha \Gamma(N_i \rightarrow L_\alpha + H_u) - \sum_\alpha \Gamma(N_i \rightarrow \overline{L_\alpha} + \overline{H_u})}{\sum_\alpha \Gamma(N_i \rightarrow L_\alpha + H_u) + \sum_\alpha \Gamma(N_i \rightarrow \overline{L_\alpha} + \overline{H_u})} \\ &= -\frac{1}{8\pi} \frac{1}{(hh^\dagger)_{ii}} \sum_{k \neq i} \text{Im} \left[\left\{ (hh^\dagger)_{ik} \right\}^2 \right] \left[\mathcal{F}_V \left(\frac{M_{Rk}^2}{M_{Ri}^2} \right) + \mathcal{F}_S \left(\frac{M_{Rk}^2}{M_{Ri}^2} \right) \right], \quad (25) \end{aligned}$$

where N_i , L_α and H_u ($\overline{L_\alpha}$ and $\overline{H_u}$) symbolically denote fermionic or scalar components of corresponding supermultiplets (and their antiparticles); $\mathcal{F}_V(x)$ and $\mathcal{F}_S(x)$ represent the

contributions from vertex and self-energy diagrams, respectively [23]:

$$\mathcal{F}_V(x) = \sqrt{x} \ln \left(1 + \frac{1}{x} \right), \quad \mathcal{F}_S(x) = \frac{2\sqrt{x}}{x-1}. \quad (26)$$

For hierarchical right-handed neutrinos $M_{R1} \ll M_{R2}, M_{R3}$, the lepton asymmetry is dominantly supplied by decays of the lightest right-handed (s)neutrinos. In the following discussion, we assume that this is the case for simplicity. In this case, the expression of the asymmetry parameter ϵ_1 is given by

$$\epsilon_1 = \frac{3}{8\pi} \frac{M_{R1} m_{\nu 3}}{\langle H_u^0 \rangle^2} \delta_{\text{eff}}, \quad (27)$$

where $m_{\nu 3}$ is the mass of the heaviest left-handed neutrino and δ_{eff} is an effective CP-violating phase. If the lightest right-handed (s)neutrinos are in thermal equilibrium, the resultant lepton asymmetry is given by the following formula:

$$\left| \frac{n_L}{s} \right| = \frac{1}{\Delta} \times \frac{45}{2\pi^2 g_*(T_B)} \frac{\zeta(3)}{\pi^2} \left(\frac{3}{2} + 2 \right) |\epsilon_1| \kappa, \quad (28)$$

where T_B is the freeze-out temperature of N_1 , and $g_*(T_B) \approx 270$; ⁹ κ denotes the fraction of the produced asymmetry that survives washout processes by lepton-number-violating interactions after N_1 decay. Here, we have assumed that there is no entropy production before the decays of the lightest messengers, and set $\Delta_M = \Delta$ for simplicity. For $\kappa \sim 1$, the following out-of-equilibrium condition should be satisfied [21]:

$$\widetilde{m}_1 = \frac{8\pi \langle H_u^0 \rangle^2}{M_{R1}^2} \Gamma_{N_1} = (hh^\dagger)_{11} \frac{\langle H_u^0 \rangle^2}{M_{R1}} \lesssim 5 \times 10^{-3} \text{ eV}. \quad (29)$$

The lepton asymmetry produced by the right-handed (s)neutrino decays is subsequently converted into the baryon asymmetry by the sphaleron effects:

$$\frac{n_B}{s} = C \frac{n_L}{s}, \quad (30)$$

where C is a number of $\mathcal{O}(1)$, which takes the value $C = -8/23$ in the MSSM.

⁹In this expression, we have assumed that a pair of messenger multiplets ($\mathbf{5} + \overline{\mathbf{5}}$) are in thermal equilibrium. After the decoupling of the messenger particles, the contents of the MSSM give rise to $g_* \approx 230$. If we take into account the effective degrees of freedom in the dynamical SUSY-breaking sector, $g_*(T_B)$ would be larger than this value, but it changes the resultant asymmetry by only a factor of at most $\mathcal{O}(1)$.

From Eqs. (6), (27), (28) and (30), we can derive the baryon asymmetry in the present Universe. In terms of the density parameter it is written as

$$\Omega_B h^2 \simeq 0.02 \left(\frac{10 \text{ keV}}{m_{3/2}} \right) \left(\frac{g_*(T_f)}{230} \right) \left(\frac{270}{g_*(T_B)} \right) \left(\frac{M_{R1}}{10^{10} \text{ GeV}} \right) \left(\frac{m_{\nu 3}}{0.06 \text{ eV}} \right) \kappa \delta_{\text{eff}}. \quad (31)$$

In Fig. 2, we show the lower bound of M_{R1} (and hence, it is the lower bound on the reheating temperature T_R) to obtain the required baryon asymmetry. Here, we have assumed the following relation for simplicity:

$$\left(\frac{g_*(T_f)}{230} \right) \left(\frac{270}{g_*(T_B)} \right) \left(\frac{m_{\nu 3}}{0.06 \text{ eV}} \right) \kappa \delta_{\text{eff}} \leq 1. \quad (32)$$

As can be seen, the observed baryon asymmetry $\Omega_B h^2 \simeq 0.02$ can be naturally explained

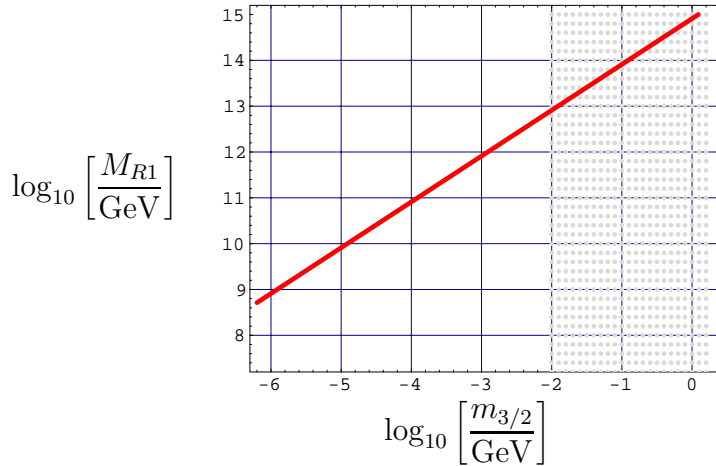


Figure 2: The lower bound on M_{R1} (and hence T_R) to obtain the observed baryon asymmetry. Here, we have assumed the hierarchical right-handed neutrinos ($M_{R1} \ll M_{R2}, M_{R3}$) and the relation given in Eq. (32). The shaded region is disfavoured by the BBN constraint.

by the thermal leptogenesis in a wide range of the gravitino mass.¹⁰ From Eq. (8), one can see that the required lower mass bound on the right-handed neutrino is always larger than the mass of the messenger fields, and then the thermalization condition for the messenger particles is always satisfied. We should stress that we do not have to fine-tune the couplings of the inflaton to the SM particles to reduce the reheating temperature, which is usually imposed in SUSY inflationary models because of the cosmological gravitino

¹⁰In the simplest chaotic inflationary scenario, for example, $\mathcal{O}(1)$ couplings of the inflaton to the SM fields lead to $T_R \approx 10^{13}$ GeV.

problem. We should also stress that there is an interesting possibility to determine $m_{3/2}$ directly in future collider experiments if the gravitino is lighter than about 100 keV [24]. Note that such a light gravitino has been considered as unlikely because of the cosmological gravitino problem. However, as we have seen, it is now indeed *favoured* from cosmological perspectives.

6 Conclusions and Discussion

In this letter, we have pointed out that there is no fine-tuning problem to obtain the required mass density of the dark matter in a certain class of GMSB models. By virtue of the small mixing between the SM fields, which is induced by the R -symmetry-breaking effects, the lightest messenger particle has a finite lifetime and provides an exquisite amount of entropy, which dilutes the thermal relics of the gravitinos down to just the mass density required for the dark matter. This phenomenon takes place automatically, regardless of the gravitino mass and the reheating temperature of inflation as long as the gravitinos and messenger fields are thermalized in the early Universe. There is no severe upper bound on the reheating temperature in this class of GMSB models, which makes the standard thermal leptogenesis very attractive as the origin of the observed baryon asymmetry.

The present scenario should have important implications also on other candidates for the origin of the present baryon asymmetry. It would be very interesting to reanalyse those models with the disappearance of the cosmological gravitino problem taken into account.

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