DESY 07-156

Dark Matter in Gauge Mediated Supersymmetry Breaking using Metastable Vacua

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arXiv:0710.1561v2 [hep-ph] 11 Oct 2007

Abstract

We point out that, in a class of gauge mediation models using metastable supersymmetry breaking vacua, the minimum of the supersymmetry breaking field in the early universe is dynamically deviated from the one in the low energy. The deviation induces coherent oscillations of the supersymmetry breaking field, which decays into the gravitinos. For certain parameters, it can produce a right amount of the gravitinos to account for the observed dark matter.

I. INTRODUCTION

Gauge mediation (GM) of supersymmetry (SUSY) breaking [1] is a natural solution to the phenomenological problems such as excessive flavor-changing neutral currents. In spite of the successes, building a realistic model is a rather non-trivial task. According to the argument of Nelson and Seiberg, an exact $U(1)_R$ symmetry is required if the superpotential is generic [2]. This observation strongly limits possible models, and a lot of efforts have been devoted to building a realistic one.

Recently, Murayama and Nomura has proposed drastically simplified models, focusing on metastable vacua in the SUSY breaking sector [3]. Even though the entire superpotential does not possess the exact $U(1)_R$ symmetry, an accidental one exists near a local SUSY breaking minimum [4]. Such a scenario is viable as long as the metastable vacua have a sufficiently long life time.

In a class of the metastable SUSY breaking models, there are the local SUSY breaking vacua near the origin of the SUSY breaking fields, where the accidental $U(1)_R$ symmetry exists. On the other hand, the breaking of $U(1)_R$ symmetry is necessarily involved in the messenger sector to mediate the SUSY breaking to the gauginos in the supersymmetric standard model (SSM). Since the SUSY breaking field (denoted by S) linearly couples to the messenger fields, the breaking of the $U(1)_R$ symmetry induces a linear term of S in the Kähler potential.

Such a linear term forces the SUSY breaking field S to deviate from its minimum in the low energy, while the inflaton field dominates the energy of the universe. When the Hubble parameter becomes comparable to the mass of S, it starts to oscillate coherently around the SUSY breaking vacuum, and then decays into the gravitinos. In the GM models, the gravitino is stable and behaves as dark matter (DM) in the universe as long as the R-parity is preserved. Thus, DM is generally produced from the SUSY breaking field in a class of the GM models using the metastable vacua. In this letter, we study the gravitino production in this scheme.

Before proceeding to the details, let us comment on the previous works. The cosmological evolution of the SUSY breaking field has been studied in the models with the metastable SUSY breaking vacua in Refs. [5]. Those literatures assumed so high reheating temperature as to make the SUSY breaking sector to be in thermal equilibrium. Then the gravitino is

expected to reach thermal equilibrium, and as a result, its abundance exceeds the observed DM abundance or it erases the density fluctuation too much, unless the gravitino mass is smaller than 16 eV [6]. For a wide range of the gravitino mass, i.e., from 16 eV to $O(10)$ GeV, therefore, the SUSY breaking sector should not be thermalized as long as the standard thermal history of the universe is assumed. In this letter, we do not pursue this possibility and assume that the SUSY breaking sector never reaches thermal equilibrium. On the other hand, Ibe and Kitano discussed the gravitino production from the SUSY breaking field in the GM models using the metastable vacua [7]. They assumed a different thermal history taking a different set of the parameters; in particular, the reheating temperature in our scenario is as high as 10^{8-10} GeV, while they considered the low reheating temperature. Furthermore, the interaction that shifts the SUSY breaking field from its minimum in the low energy is different from that considered in Ref. [7].

II. MODEL

In this section we provide a model of the gauge mediation using the metastable vacua. To be explicit, we adopt a model by Murayama and Nomura given in Ref. [3]. The Kähler potential and the superpotential for a gauge singlet chiral field S and the messengers f and $\bar{f}$ are written as

$$K = |S|^2 - \frac{|S|^4}{4\Lambda^2} + |f|^2 + |\bar{f}|^2, \tag{1}$$

$$W = -\mu^2 S + \kappa S f \bar{f} + M f \bar{f}, \tag{2}$$

where the higher order corrections of $O(|S|^6/\Lambda^4)$ are omitted for simplicity in the Kähler potential. In the following we take f and $\bar{f}$ to be in $\mathbf{5} + \mathbf{5}^*$ representation of $SU(5)_{\rm GUT}$, and μ^2, κ and M are set to be real and positive without loss of generality. We assign the charges $U(1)_R$ symmetry as $R[S] = 2$ and $R[f] = R[\bar{f}] = 0$. Then one can see that the messenger mass term explicitly violates the $U(1)_R$ symmetry to Z_2.

For our purpose we do not need to specify the UV physics above a scale Λ that provides the second term in Eq. (1) as well as the first term in Eq. (2). We simply note here that there are many explicit models that actually lead to this low energy effective theory. (See Ref. [3] for examples.) In particular, we do not give the SUSY breaking mechanism explicitly here, which is assumed to be such that the first term in Eq. (2) is somehow produced.

From the Kähler potential and the superpotential given above, one can show that there is a SUSY minimum at

$$S \;=\; -\frac{M}{\kappa}, \quad f \;=\; \bar{f} \;=\; \frac{\mu}{\sqrt{\kappa}}. \tag{3}$$

On the other hand, SUSY is broken at $S = f = \bar{f} = 0$, which is a metastable local minimum as long as

$$M^2 \;>\; \kappa\mu^2 \tag{4}$$

is satisfied, since otherwise one of the messenger scalars becomes tachyonic. Note that the second term in Eq. (1) produces a positive mass of S, $m_S = \mu^2/\Lambda$, around the origin. The SUSY breaking scale is dictated by the first term in Eq. (2), and the F-term of S is given by $F_S \simeq \mu^2$. Requiring a vanishing cosmological constant, we can relate the SUSY breaking scale to the gravitino mass $m_{3/2}$ as

$$\mu \;=\; \left(\sqrt{3}\, m_{3/2} M_P\right)^{\frac{1}{2}} \simeq 2 \times 10^9\,\mathrm{GeV} \left(\frac{m_{3/2}}{1\,\mathrm{GeV}}\right)^{\frac{1}{2}}, \tag{5}$$

where $M_P = 2.4 \times 10^{18}\,\mathrm{GeV}$ is the reduced Planck scale. The SUSY breaking effects are transmitted to the visible sector by the messenger loops. The integration of the messengers give rise to the gaugino masses as [1]

$$m_i \simeq \frac{\alpha_i}{4\pi}\frac{\kappa\mu^2}{M} \qquad \text{for} \;\; i = 1, 2, 3. \tag{6}$$

Here, $m_{1,2,3}$ and $\alpha_{1,2,3}$ are the gaugino masses and the gauge coupling constants for $U(1)_Y, SU(2)_L$ and $SU(3)_C$ in the SSM. We have used the $SU(5)_{\mathrm{GUT}}$ normalization for the $U(1)_Y$ gauge coupling constant. We can express $\kappa\mu^2/M$ in terms of the gluino mass m_3:

$$\frac{\kappa\mu^2}{M} \;\simeq\; 1 \times 10^2\,\mathrm{TeV} \left(\frac{\alpha_3}{0.1}\right)^{-1} \left(\frac{m_3}{1\mathrm{TeV}}\right). \tag{7}$$

From (5) and (7), one can express M as

$$M \;\simeq\; 3 \times 10^{13}\,\mathrm{GeV}\,\kappa \left(\frac{\alpha_3}{0.1}\right) \left(\frac{m_3}{1\mathrm{TeV}}\right)^{-1} \left(\frac{m_{3/2}}{1\mathrm{GeV}}\right). \tag{8}$$

We also assume $m_S \lesssim \Lambda$, or equivalently,

$$\mu \;\lesssim\; \Lambda, \tag{9}$$

since we consider the dynamics of S (e.g. coherent oscillations and decay), which should be described within the low energy effective theory. Using (9), we obtain an upper-bound on

m_S,

$$m_S = \frac{\mu^2}{\Lambda} \simeq 40\,\mathrm{TeV}\left(\frac{m_{3/2}}{1\,\mathrm{GeV}}\right)\left(\frac{\Lambda}{10^{14}\mathrm{GeV}}\right)^{-1}, \lesssim 2\times 10^9\,\mathrm{GeV}\left(\frac{m_{3/2}}{1\,\mathrm{GeV}}\right)^{\frac{1}{2}}. \quad (10)$$

Lastly, let us discuss radiative corrections to the Kähler potential. Integrating the messenger loop, the relevant corrections are given by

$$K^{(1)} = K^{(1)}_{nh} + K^{(1)}_{h}, \quad (11)$$

with

$$K^{(1)}_{nh} = -\frac{5M^2}{16\pi^2}\left\{\frac{1}{2}\left(\frac{\kappa}{M}\right)^3|S|^2(S+S^\dagger) - \frac{1}{6}\left(\frac{\kappa}{M}\right)^4|S|^2(S^2+S^{\dagger 2})+\cdots\right\}, \quad (12)$$

$$K^{(1)}_{h} = -\frac{5M^2}{16\pi^2}\left\{\frac{\kappa}{M}(S+S^\dagger)+\cdots\right\}, \quad (13)$$

where we have separated the holomorphic terms and the non-holomorphic ones. Note that S is assumed to be much smaller than M/κ so that S sits far away from the SUSY minimum (see (3)). As explained in Introduction, the reason why such corrections appear is that the R-symmetry is explicitly broken by the messenger mass term.

One can check that the radiative corrections (12) reproduce the result of the Coleman-Weinberg potential for S given in Ref. [3], up to the order explicitly shown in (12):

$$V^{(1)}_{nh} = \frac{5\mu^4}{16\pi^2}\left\{\frac{\kappa^3}{M}(S+S^\dagger) - \frac{\kappa^4}{2M^2}(S^2+S^{\dagger 2})+\cdots\right\} \quad (14)$$

To avoid the mass of S to become tachyonic due to the radiative corrections, we require [3]

$$M \gtrsim \frac{\kappa^2}{4\pi}\Lambda, \quad (15)$$

throughout this letter.

As far as the SUSY breaking sector is concerned, the linear term in the Kähler potential (13) does not modify the scalar potential significantly. In the very early universe, however, such a linear term makes the minimum of the scalar potential to deviate from the origin. Therefore, it is crucial for cosmological evolution of S to take into account the linear term in the Kähler potential, as we will show in the next section.

III. COSMOLOGY

Now we consider the cosmological evolution of the SUSY breaking field, S. First let us give a sketch how S is deviated from the origin due to the linear term in the Kähler potential, and estimate the cosmic abundance. While the F-term of the inflaton dominates the universe, the scalar potential of S is approximately given by [a]

$$\begin{aligned} V(S) &\simeq e^K (3H^2 M_P^2), \\ &\simeq 3H^2 \left(|S|^2 - \frac{5\kappa}{16\pi^2} M (S + S^\dagger) + \cdots \right), \end{aligned} \tag{16}$$

where we have assumed that S does not couple to the inflaton in the Kähler potential for simplicity [b]. The scalar potential has a minimum given by

$$S_c = \frac{5\kappa}{16\pi^2} M. \tag{17}$$

If this minimum S_c exceeds Λ, there is no stable minimum in the low-energy theory in the early universe. Depending on the UV theory, the system may settle down at the SUSY minimum in this case. To avoid such a situation, we impose $S_c < \Lambda$ in the following.

When the Hubble parameter becomes comparable to the mass of S, it starts to oscillate around the minimum, $S \simeq 0$, with an initial amplitude S_c. The abundance of S is estimated as

$$\begin{aligned} \frac{n_S}{s} &\simeq \frac{3T_R}{4} \left(\frac{5\kappa}{16\pi^2} \right)^2 \frac{m_S M^2}{3 m_S^2 M_P^2}, \\ &\simeq 1 \times 10^{-10}\, \kappa^4 \left(\frac{\alpha_3}{0.1} \right)^2 \left(\frac{m_3}{1\mathrm{TeV}} \right)^{-2} \left(\frac{m_{3/2}}{1\mathrm{GeV}} \right) \left(\frac{T_R}{10^8\mathrm{GeV}} \right) \left(\frac{\Lambda}{10^{14}\mathrm{GeV}} \right), \end{aligned} \tag{18}$$

where n_S is the number density of S, s is the entropy density, T_R denotes the reheating temperature, and we have used (8) to eliminate M. We have assumed here that the reheating has not completed when $H = m_S$. This assumption is indeed reasonable, since otherwise too many gravitinos are produced by thermal scatterings in plasma.

Several comments are in order. First, we have assumed that the Hubble parameter during inflation, H_I, is larger than the mass of S, i.e., $H_I > m_S$. If m_S is larger than

[a] The linear term of S generically appear in the scalar potential due to the supergravity effects [7]. One can neglect its effect on the dynamics of S, as long as $\kappa \gtrsim 0.04\,(\alpha_3/0.1)^{1/2}(1\mathrm{TeV}/m_3)^{1/2}(m_{3/2}/1\mathrm{GeV})^{1/2}(\Lambda/10^{14}\mathrm{GeV})$. We assume that this inequality is satisfied in the following analysis.

[b] Even in the presence of the interactions, the following argument does not change qualitatively.

H_I, the deviation of S is suppressed by $(H_I/m_S)^2$. Considering the upper-bound on m_S given by (10), however, the assumption $H_I > m_S$ is valid except for low-scale inflation models. Second, although our model is given only below the scale Λ, this does not limit the application of the above arguments only to the inflation models with $H_I < \Lambda$. For $H_I < \Lambda$, the above scalar potential (16) is obviously valid both during and after inflation. On the other hand, for $H_I > \Lambda$, one cannot use (16) during inflation. After inflation, however, the Hubble parameter decreases and becomes smaller than Λ at certain point. If the system can be then described by the low energy effective theory, S will quickly settle down at the potential minimum (17) [c].

Next let us consider the decay processes of S. The SUSY breaking field S will decay into a pair of the gravitinos, and the decay rate is [8, 9]

$$\Gamma_{3/2} \simeq \frac{1}{96\pi}\frac{m_S^5}{m_{3/2}^2 M_P^2} \simeq \frac{1}{32\pi}\frac{m_S^3}{\Lambda^2}, \tag{19}$$

where we have used (5). If m_S is smaller than $2M$, S decays into the gauge sector through the messenger loop [d]. The relevant interactions are

$$\mathcal{L} \simeq -\frac{\alpha_i}{4\pi}\frac{\kappa}{M}S\Big[-\frac{1}{4}F^{(i)}_{\mu\nu}F^{(i)\mu\nu} + \frac{i}{8}\epsilon^{\mu\nu\rho\sigma}F^{(i)}_{\mu\nu}F^{(i)}_{\rho\sigma} - \frac{\kappa F_S}{M}\bar{\lambda}^{(i)}\mathcal{P}_L\lambda^{(i)}\Big] + \text{h.c.}, \tag{20}$$

where we neglected terms with higher orders of $\kappa\langle S\rangle/M$. In particular, S decays into the gluons and gluinos, and the decay rates are

$$\Gamma_g \simeq \frac{\alpha_3^2\kappa^2}{64\pi^3}\frac{m_S^3}{M^2} \tag{21}$$

and [7]

$$\Gamma_{\tilde{g}} \simeq \frac{\kappa^2}{\pi}\frac{m_3^2\, m_S}{M^2}, \tag{22}$$

respectively. Note that the decay rate into the gluinos is smaller than that into the gluons, if m_S is much larger than the gluino mass [10],

$$m_S \gg \frac{8\pi}{\alpha_3}m_3, \tag{23}$$

[c] Depending on the details of the UV theory, it is possible that the position of S is larger than Λ and the system cannot be described by the model given by (1) and (2) even for $H < \Lambda$. Then, the abundance of S will generically become larger than (18), and so, our estimate is conservative.

[d] The decay into the scalars does not change our argument significantly.

and vice versa. Using (19), (21) and (22), we obtain the branching ratio of the gravitino production,

$$B_{3/2} \simeq \frac{1}{1+r} \tag{24}$$

with

$$r \equiv \frac{8}{3}\left(\frac{m_3\,\Lambda}{m_{3/2}\,M_P}\right)^2 + \frac{32}{9}\left(\frac{4\pi}{\alpha_3}\right)^2\left(\frac{m_3\,\Lambda}{m_{3/2}\,M_P}\right)^4. \tag{25}$$

Therefore, the SUSY breaking field will dominantly decay into the gravitinos if $r \lesssim 1$, or equivalently,

$$\Lambda \lesssim 2\times 10^{14}\,\mathrm{GeV}\left(\frac{\alpha_3}{0.1}\right)^{\frac{1}{2}}\left(\frac{m_3}{1\mathrm{TeV}}\right)^{-1}\left(\frac{m_{3/2}}{1\mathrm{GeV}}\right). \tag{26}$$

We will assume this inequality is met in the following analysis. The decay temperature of S is given by

$$\begin{aligned} T_d &\equiv \left(\frac{\pi^2 g_*}{10}\right)^{-\frac{1}{4}}\sqrt{\Gamma_{3/2}M_P}, \\ &\simeq 4\,\mathrm{GeV}\left(\frac{g_*}{100}\right)^{-\frac{1}{4}}\left(\frac{m_{3/2}}{1\mathrm{GeV}}\right)^{\frac{3}{2}}\left(\frac{\Lambda}{10^{14}\mathrm{GeV}}\right)^{-\frac{5}{2}}, \end{aligned} \tag{27}$$

where g_* counts the relativistic degrees of freedom at the decay.

Now we can estimate the gravitino abundance. Since S dominantly decays into a pair of the gravitinos, the gravitino abundance is given by

$$Y_{3/2} \simeq 2\times 10^{-10}\,\kappa^4\left(\frac{\alpha_3}{0.1}\right)^2\left(\frac{m_3}{1\mathrm{TeV}}\right)^{-2}\left(\frac{m_{3/2}}{1\mathrm{GeV}}\right)\left(\frac{T_R}{10^8\mathrm{GeV}}\right)\left(\frac{\Lambda}{10^{14}\mathrm{GeV}}\right). \tag{28}$$

The density parameter of the gravitino is

$$\Omega_{3/2}h^2 \simeq 0.06\,\kappa^4\left(\frac{\alpha_3}{0.1}\right)^2\left(\frac{m_3}{1\mathrm{TeV}}\right)^{-2}\left(\frac{m_{3/2}}{1\mathrm{GeV}}\right)\left(\frac{T_R}{10^8\mathrm{GeV}}\right)\left(\frac{\Lambda}{10^{14}\mathrm{GeV}}\right). \tag{29}$$

where h is the present Hubble parameter in units of 100 km/s/Mpc. Therefore, a right amount of the gravitinos can be produced by the decay of the SUSY breaking field, S. Note that, for the non-thermally produced gravitinos to be a dominant component of DM, κ should not be suppressed. The reason is as follows. For a small κ, the messenger mass is also small to keep the size of the soft masses in the SSM sector (see (8)). Since the shift of the S field is proportional to the breaking of $U(1)_R$ symmetry, i.e., the messenger mass, the gravitino abundance is suppressed for a small value of κ. Furthermore, depending on the reheating temperature and the mass spectrum of the SSM particles, the thermal production of the gravitinos and the NLSP decay may also give sizable contributions [11, 12, 13, 14, 15].

Note also that the relatively high reheating temperature is favored, which may accommodate the thermal leptogenesis scenario [17].

Finally, let us estimate the free streaming length of the gravitinos produced by the decay of S. The comoving free streaming length λ_{FS} at matter-radiation equality is defined by

$$\lambda_{FS} \equiv \int_{t_D}^{t_{\rm eq}} \frac{v_{3/2}(t)}{a(t)} dt, \tag{30}$$

where $a(t)$ is the scale factor, and t_D and $t_{\rm eq} (\sim 2 \times 10^{12}\,{\rm sec})$ denote the time at the S decay and at matter-radiation equality, respectively. $v_{3/2}$ is the velocity of the gravitino, given by

$$v_{3/2}(t) = \frac{|\mathbf{p}_{3/2}|}{E_{3/2}} \simeq \frac{\frac{m_S}{2}\left(\frac{a_D}{a(t)}\right)}{\sqrt{m_{3/2}^2 + \frac{m_S^2}{4}\left(\frac{a_D}{a(t)}\right)^2}}, \tag{31}$$

where we have approximated $m_S \gg m_{3/2}$, and a_D is the scale factor at the decay of S. Integrating (30) yields

$$\begin{aligned} \lambda_{FS} &\simeq \frac{1}{H_0\sqrt{1+z_{\rm eq}}} X^{-1} \sinh^{-1} X, \\ &\sim 1\,{\rm kpc} \left(\frac{g_*}{100}\right)^{\frac{1}{2}} \left(\frac{m_{3/2}}{1{\rm GeV}}\right)^{-\frac{3}{2}} \left(\frac{\Lambda}{10^{14}{\rm GeV}}\right)^{\frac{3}{2}} \end{aligned} \tag{32}$$

with

$$\begin{aligned} X &\equiv \frac{2m_{3/2}}{m_S}\frac{a_{\rm eq}}{a_D}, \\ &\sim 10^6 \left(\frac{g_*}{100}\right)^{-\frac{1}{2}} \left(\frac{m_{3/2}}{1{\rm GeV}}\right)^{\frac{3}{2}} \left(\frac{\Lambda}{10^{14}{\rm GeV}}\right)^{-\frac{3}{2}}, \end{aligned} \tag{33}$$

where H_0 is the Hubble parameter at present, and $z_{\rm eq}$ and $a_{\rm eq}$ are the red-shift and the scale factor at the matter-radiation equality. In the second equation of (32), we have used $H_0^{-1} \sim 4 \times 10^3\,{\rm Mpc}$ and $z_{\rm eq} \sim 3000$. Thus, the free streaming length λ_{FS} is expressed in terms of the gravitino mass and the scale Λ, and it can be as small as 1 kpc. Interestingly, the recent observations on the dSph galaxies seem to exhibit a sharp cut-off around 100 pc [16] in the smallest size of the galaxies, which may be explained by DM with free streaming length of $O(100\,{\rm pc})$ [e]. Our scenario may be supported by further observations in the near future.

[e] It is also possible to produce a right amount of the gravitino DM from the inflaton decay [18], and the gravitino can have right free streaming length to explain the cut-off in the smallest size of the galaxies.

IV. CONCLUSIONS AND DISCUSSION

The metastable SUSY breaking vacua provide a drastically simplified scheme of the gauge mediation. The models possess an accidental $U(1)_R$ symmetry, which is broken in the messenger sector. The breaking induces a linear term of the SUSY breaking field in the Kähler potential. We have pointed out that the SUSY breaking field is forced away from its minimum due to this linear term while the inflaton field dominates the energy of the universe. Then the gravitinos are produced when it decays, and we have shown that a right abundance of the gravitino DM can be realized for certain parameters. Further, the free streaming length of the gravitino may explain the recent observations on the smallest size of the dSph galaxies.

Acknowledgment

We thank W. Buchmüller for reading the manuscript and comments.

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