Gravitino Production in the Early Universe and Its Implications to Particle Cosmology

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Abstract

Effects of the unstable gravitino on the big-bang nucleosynthesis (BBN) and its implications to particle cosmology are discussed. If the gravitino mass is smaller than $\sim 20$ TeV, lifetime of the gravitino becomes longer than $\sim 1$ sec and its decay may spoil the success of the standard BBN. In order to avoid such a problem, upper bound on the reheating temperature after the inflation is obtained, which may be as low as $\sim 10^5 - 10^6$ GeV. For a successful baryogenesis with such low reheating temperature, a consistent scenario based on the large cutoff supergravity (LCSUGRA) hypothesis of supersymmetry breaking, where the gravitino and sfermion become as heavy as $\sim O(1-10)$ TeV, is proposed. In the LCSUGRA, non-thermal leptogenesis can produce large enough baryon asymmetry. We also see that, in the LCSUGRA scenario, relic density of the lightest superparticle becomes consistent with the WMAP value of the dark matter density in the parameter region required for the successful non-thermal leptogenesis. In this case, the dark matter density may be reconstructed with the future $e^+e^-$ linear collider.

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1 Introduction

Supersymmetry (SUSY) has been attracted many attentions not only from particle-physics point of view but also from cosmological point of view. Indeed, supersymmetry may give new insights into cosmological problems to which standard model cannot answer, like the origin of the cold dark matter, dynamics of the scalar field responsible to the inflation, mechanism to generate the baryon asymmetry, and so on. Thus, even in cosmology, it is expected that SUSY will play important roles.

In order to construct viable and natural scenario of the evolution of the universe in the framework of the supersymmetric models, there is one serious problem caused by the gravitino, which is the superpartner of the graviton. Since the gravitino may have lifetime longer than $\sim 1$ sec if its mass is lighter than $\sim O(10$ TeV), thermally produced gravitino in the early universe may decay after the big-bang nucleosynthesis (BBN). If this is the case, the decay products of the gravitino induce electromagnetic and hadronic shower and the high energy particles in the shower cause dissociation of the light elements produced by the BBN. Since the standard BBN scenario more or less predicts light element abundances consistent with the observations, primordial gravitino may spoil the success of the BBN if its abundance is too large \cite{1}. Usually, for unstable gravitino with relatively small mass, this problem (called “gravitino problem”) is avoided by putting upper bound on the reheating temperature after inflation. (For details, see \cite{2,3,4,5} and references therein.).\footnote{In fact, the reheating temperature here should be understood as the maximal temperature when the (last) radiation dominated epoch is realized. If some scalar field other than the inflaton dominates the universe after the inflation, reheating temperature here is given by the temperature at the time of the decay of such scalar particle. Examples of such scenarios are given in, for example, \cite{6}.} As we will see, the bound is quite stringent and hence it is required to construct a scenario of cosmological evolution consistent with such low reheating temperature.

Here, we would like to discuss three subjects which are closely related. First, we review the current situation of the calculation of the upper bound on the reheating temperature from the gravitino problem. Then, we propose a cosmological scenario based on large cutoff supergravity scenario, which is consistent with the constraints on the reheating temperature. In this scenario, the baryon asymmetry of the universe is explained by the non-thermal leptogenesis while the LSP becomes a good candidate of the dark matter of the universe. Finally, we consider a possible test of such a scenario; we point out that the precise reconstruction of the dark matter density may be possible in this scenario once the superparticles are produced at the linear collider.

2 Gravitino Problem

We first briefly review the gravitino production in the early universe and its effects on the BBN \cite{2,3,4,5}. Even though the gravitino is a very weakly interacting particle, it can be produced in the early universe by the scattering processes of the particles in the thermal
bath. Using the thermally averaged gravitino production cross section given in [7], the “yield variable” of the gravitino, which is defined as \( Y_{3/2} \equiv \frac{n_{3/2}}{s} \), is given by [3]

\[
Y_{3/2} \approx 1.9 \times 10^{-12} \times \left( \frac{T_R}{10^{10} \text{ GeV}} \right) \left[ 1 + 0.045 \ln \left( \frac{T_R}{10^{10} \text{ GeV}} \right) \right] \left[ 1 - 0.028 \ln \left( \frac{T_R}{10^{10} \text{ GeV}} \right) \right],
\]

where \( n_{3/2} \) is the number density of the gravitino while \( s = \frac{2\pi^2}{45} g_s S(T) T^3 \) is the entropy density with \( g_s S(T) \) being the effective number of the massless degrees of freedom at the temperature \( T \), and the reheating temperature is defined as

\[
T_R \equiv \left( \frac{10}{g_s \pi^2} M_*^2 \Gamma_{\text{inf}}^2 \right)^{1/4},
\]

with \( \Gamma_{\text{inf}} \) being the decay rate of the inflaton. (Here, \( M_* \approx 2.4 \times 10^{18} \text{ GeV} \) is the reduced Planck scale.) As one can see, the number density of the gravitino increases as the reheating temperature becomes larger.

Since the gravitino decays with very long lifetime (if it is unstable), it may decay after the BBN starts. Indeed, if the gravitino mass \( m_{3/2} \) is smaller than \( \sim 20 \text{ TeV} \), its lifetime becomes longer than 1 sec and its decay may affect the abundances of the light elements which are synthesised by the standard BBN processes. In particular, since the prediction of the standard BBN is in a reasonable agreement with the observations, light-element abundances become inconsistent with the observations if the abundance of the gravitino is too large. Consequently, we obtain the upper bound on the reheating temperature.

In [2, 3, 4], with a detailed analysis of the non-standard processes induced by the gravitino decay (in particular, the hadro- and photo-dissociations as well as the \( p \leftrightarrow \bar{n} \) conversion), light-element abundances are calculated as a function of the reheating temperature and the gravitino mass. In addition, in the analysis given in [4], decay processes of the gravitino are studied in detail. Then, comparing the resultant light-element abundances with the observations, upper bound on the reheating temperature is obtained.

One of the results is shown in Fig. [4] Here, the mSUGRA-type parameterisation of the minimal supersymmetric standard model (MSSM) parameters is used and, for Fig. [4] the following choice of the model parameters is adopted: unified gaugino mass \( m_{1/2} = 300 \text{ GeV} \), universal scalar mass \( m_0 = 2397 \text{ GeV} \), SUSY invariant Higgs mass \( \mu_H = 231 \text{ GeV} \), and \( \tan \beta = 30 \). With this choice of parameters, the mass of the LSP becomes 116 GeV. (The parameter region where the gravitino mass becomes lighter than the LSP mass is shaded.) In the figure, each line shows the upper bound on the reheating temperature from the considerations of different light-element abundances. Here, the following observational constraints are used. For D/H [8],

\[
D/H(\text{Low}) = (2.78^{+0.44}_{-0.38}) \times 10^{-5},
\]
for $^3\text{He}/D$ [9]

$$^3\text{He}/D < 0.59 \pm 0.54 \ (2\sigma); \quad (2.5)$$

for $^4\text{He}$ mass fraction $Y$, one by Fields and Olive [10]

$$Y(\text{FO}) = 0.238 \pm (0.002)_{\text{stat}} \pm (0.005)_{\text{syst}}, \quad (2.6)$$

one by Izotov and Thuan [11]

$$Y(\text{IT}) = 0.242 \pm (0.002)_{\text{stat}} \pm (0.005)_{\text{syst}}, \quad (2.7)$$

and one by Olive and Skillman [12]

$$Y(\text{OS}) = 0.249 \pm 0.009; \quad (2.8)$$

for $^7\text{Li}$ [13]

$$\log_{10} \left( ^7\text{Li}/H \right) = -9.66 \pm (0.056)_{\text{stat}} \pm (0.300)_{\text{add}}; \quad (2.9)$$

and for $^6\text{Li}$ [14]

$$^6\text{Li}/H < (1.10^{+5.14}_{-0.94}) \times 10^{-11} \ (2\sigma). \quad (2.10)$$

The upper bound on the reheating temperature is from various effects. When the gravitino mass is larger than a few TeV, most of the primordial gravitinos decay at very early stage of the BBN. In this case, in addition, photo- and hadro-dissociations are ineffective. Then, overproduction of $^4\text{He}$ due to the $p \leftrightarrow n$ conversion becomes the most important. We note here that we consider three different observational constraints on $^4\text{He}$, which are given by Fields and Olive (FO) [10], Izotov and Thuan (IT) [11] and Olive and Skillman (OS) [12]. As one can see, the upper bound on $T_R$ in this case is sensitive to the observational constraint on the primordial abundance of $^4\text{He}$; for the case of $m_{3/2} = 10$ TeV, for example, $T_R$ is required to be lower than $3 \times 10^7$ GeV if we use the lowest value of $Y$ given by Fields and Olive, while, with the highest value given by Olive and Skillman, the upper bound on the reheating temperature becomes as large as $4 \times 10^9$ GeV.

When $400$ GeV $\lesssim m_{3/2} \lesssim 5$ TeV, gravitinos decay when the cosmic temperature is $1$ keV $- 100$ keV. In this case, hadro-dissociation gives the most stringent constraints; in particular, the overproduction of $D$ and $^6\text{Li}$ become important. Furthermore, when the gravitino mass is relatively light ($m_{3/2} \lesssim 400$ GeV), the most stringent constraint is from the ratio $^3\text{He}/D$ which may be significantly changed by the photo-dissociation processes of $^4\text{He}$.

To see how the upper bound depends on the mass spectrum of the MSSM particles, in Fig. 2 result with $m_{1/2} = 1200$ GeV, $m_0 = 800$ GeV, $\mu_H = 1215$ GeV, and $\tan \beta = 45$ is shown. (In this case, the lightest neutralino mass is given by 509 GeV). As one can see, the behaviour qualitatively similar, although the detailed bound depends on the mass spectrum.
3 Non-Thermal Leptogenesis and LCSUGRA

So far, we have seen that, in order to suppress the gravitino production in the early universe, it is required that the reheating temperature after the inflation be lower than $10^5 - 10^7$ GeV for the gravitino mass 100 GeV – 10 TeV, if the gravitino is unstable. This result imposes a serious constraint on one of the well-motivated mechanisms of baryogenesis, the leptogenesis scenario \[15\] where the present baryon asymmetry of the universe is generated from the decay of thermally produced right-handed neutrinos.

In order to generate large enough baryon number asymmetry by the thermal leptogenesis scenario, it is necessary to raise the reheating temperature up to $10^9 - 10^{10}$ GeV or higher \[16\]. We can see that such high reheating temperature conflicts with the upper bound on $T_R$ given in Figs. 1 and 2 for large range of the gravitino mass.

The thermal leptogenesis may become viable if the gravitino mass is extremely large ($m_{3/2} \gtrsim 100$ TeV) or if the gravitino is stable. However, in this article, we would like to pursue another direction, the non-thermal leptogenesis \[17\]. In the non-thermal leptogenesis scenario, the right-handed neutrinos are assumed to be produced by the decay of the inflaton. Then, the rest of the scenario is almost the same as the thermal leptogenesis; decay of the non-thermally produced right-handed neutrinos produce lepton number.
asymmetry, which is converted to the baryon number asymmetry by the spharelon effect.

One can easily see that, in the non-thermal leptogenesis scenario, larger amount of the baryon number asymmetry can be generated with the same reheating temperature compared to the thermal leptogenesis case. In the non-thermal leptogenesis, the inflaton $\Phi$ decays into the right-handed neutrinos $\Phi \rightarrow \nu_R\nu_R$. Thus, at the time of the reheating when the energy density of the inflaton $\rho_{\text{inf}}$ gives the rough estimate of $T_R^4$, baryon asymmetry is estimated by

$$n_B(T_R) \sim \frac{\epsilon \rho_{\text{inf}}}{m_\Phi}.$$  \hspace{1cm} (3.1)

Here, $\epsilon$ is the baryon asymmetry from the single decay of $\nu_R$, and is estimated as \[18\]

$$\epsilon \equiv \frac{\Gamma(N_1 \rightarrow H_u + \ell) - \Gamma(N_1 \rightarrow H_u^* + \ell^*)}{\Gamma_{N_1}} \simeq -\frac{3}{8\pi} \frac{M_1}{\langle H_u \rangle^2} m_{\nu_3} \delta_{\text{eff}},$$  \hspace{1cm} (3.2)

where $m_{\nu_3}$ is the heaviest (active) neutrino mass, $M_1$ the mass of the lightest right-handed neutrino, and $\delta_{\text{eff}}$ the effective CP-violating phase which is assumed to be $\sim 1$. Using the
relation $\rho_{\text{inf}} \sim T_R^4$, we obtain the following relation for the baryon-to-entropy ratio

$$\frac{n_B}{s} = 3.5 \times 10^{-11} \times \kappa \left( \frac{T_R}{10^6 \text{GeV}} \right) \left( \frac{2M_1}{m_{\Phi}} \right) \left( \frac{m_{\nu_3}}{0.05 \text{eV}} \right) \delta_{\text{eff}},$$

(3.3)

where $\kappa$-parameter is a constant which is expected to be of $O(1)$. We evaluated this constant by numerically solving the Boltzmann equations and found that $\kappa = 2.44$. We can see that, in order to generate observed baryon asymmetry, reheating temperature of $O(10^6 \text{ GeV})$ is enough in the case of non-thermal leptogenesis.

Of course, even if the reheating temperature is as low as $O(10^6 \text{ GeV})$, gravitino may be still overproduced in some case, as seen in Figs. 1 and 2 if the hadronic branching ratio is close to 1, the gravitino mass should be larger than a few TeV in this case. Although small hierarchy is possible between the gravitino mass and scalar masses in the observable sector even in the gravity-mediated SUSY breaking scenario, the gravitino mass larger than $\sim$ TeV seems quite high from the point of view of the naturalness of the electroweak symmetry breaking since, in the gravity mediated SUSY breaking scenario, gravitino mass provides rough estimate of the soft SUSY breaking scalar masses.

Recently, however, one interesting scenario has been proposed where relatively large gravitino mass is realized. The scenario is based on the hypothesis such that all higher dimensional operators such as quartic terms in the Kähler potential are suppressed by a cut-off scale much higher than the Planck scale $M_*$. In this scenario, the sfermions and gravitino become order-of-magnitude heavier than the gauginos and, consequently, masses of the sfermions and gravitino are required to be significantly larger than the electroweak scale. Even so, naturalness of the electroweak symmetry breaking can be maintained by the focus-point mechanism due to the fact that the universality of the scalar masses at the GUT scale is guaranteed in this scenario. In [19], it was shown that this scenario, called “large cutoff supergravity (LCSUGRA) scenario,” is well consistent with low-energy phenomenology. In particular, heaviness and universality of the sfermion masses are good for suppressing dangerous supersymmetric effects on the flavor violating processes, proton decay, and so on. We consider that the presence of the large cutoff is a reflection of a more fundamental physics beyond the GUT scale.

In the LCSUGRA scenario, it is notable that we can construct natural and consistent scenario of cosmology. In particular, LCSUGRA predicts relatively large gravitino mass $\sim$ a few TeV, so the serious gravitino problem may be evaded with the reheating temperature of $O(10^6 \text{ GeV})$. As we mentioned, even with such a low reheating temperature, baryon asymmetry can be produced by the non-thermal leptogenesis.

In addition, it is also notable that the lightest neutralino $\chi_0^0$ becomes the LSP in this scenario, and it can be a good candidate of the dark matter. Although all the sfermions acquire multi-TeV masses in this scenario, the pair annihilation of the lightest neutralino (which is dominantly the Bino) can be enhanced via the sizable Higgsino component.

The low-energy effective theory from the LCSUGRA is the same as mSUGRA-type
models, and the low-energy MSSM parameters are parameterised by the unified gaugino mass $m_{1/2}$, the universal scalar mass $m_0$, and $\tan \beta$. So, we can calculate the relic density of the LSP as a function of these parameters. With a detailed numerical calculations, we calculate the relic density and obtained the parameter region where the relic density becomes consistent with the WMAP value \[23\]

\[\Omega_{\text{DM}}^{(\text{WMAP})} h^2 = 0.1126^{+0.0161}_{-0.0181}. \tag{3.4}\]

The result is shown in Fig. 3 on $\tan \beta$ vs. $m_0$ plane.\[^3\] In our numerical calculations, we have used the ISAJET 7.69 code \[24\] which takes into account the one-loop corrections to the effective Higgs potential and the two-loop RG evolutions of parameters.\[^4\] In the figure, there is an upper bound on $\tan \beta$, which is from the lower bounds on the $\mu$ and $m_2$. The upper bound corresponds to parameters $(\tan \beta, m_0, \mu, m_2) \simeq (25, 4 \text{ TeV}, 140 \text{ GeV}, 160 \text{ GeV})$. (The lower bound on the $\tan \beta$ comes from the upper bound on $m_2 \lesssim 1 \text{ TeV}$ where we confine our attention.) From the figure, we find that the relic density is consistent for the WMAP result for $m_0 \gtrsim 2 \text{ TeV}$.

\[^3\]Here, we reassure that the heavier CP-even, CP-odd Higgs bosons and all sfermions are much heavier than the neutralino.

\[^4\]It should be noted that, the lines in the figures show rough fitting of the results, since the code becomes somewhat unstable for $m_0 \gg m_{1/2}$. 

Figure 3: Cosmologically allowed regions of the relic density on the $\tan \beta$ vs. $m_0$ plane for $m_{\text{top}} = 174\text{GeV}$. 

\[^3\]This diagram is published under a Creative Commons Attribution 4.0 International License. 

\[^4\]The ISAJET code is available at http://www.isa.infn.it/isajet/.
4 Reconstructing $\Omega_{\text{LSP}}$

Finally, we discuss a possible test of this scenario at the future $e^+e^-$ collider (which is recently called as the International Linear Collider, or ILC) [25, 26, 27]. If the cold dark matter consists of the LSP from the focus-point supersymmetry, the dark matter density may be reconstructed once the charginos and neutralinos become kinematically accessible the ILC [28].

The important point is that, in the focus point scenario, all the sfermions (as well as the heavier Higgses) acquire masses of $O(1 \text{ TeV})$, so the pair annihilation of the lightest neutralino is dominated by the processes $\chi_1^0 \chi_1^0 \to t\bar{t}$, $\chi_1^0 \chi_1^0 \to W^+W^-$ and $\chi_1^0 \chi_1^0 \to ZZ$. These processes are via the Higgsino component in the lightest neutralino, as mentioned before. Thus, in this case, the dark-matter density is primarily determined by four parameters $m_{G1}$, $m_{G2}$, $\mu_H$, and $\tan \beta$. (Here, $m_{G1}$ and $m_{G2}$ are gaugino masses for the $U(1)_Y$ and $SU(2)_L$ gauge groups, respectively.) Since some of the superparticles (in particular, the charginos and the neutralinos) are relatively light, we may be able to obtain information about these parameters from the study of these superparticles.

If the relic density of the LSP is close to the WMAP value, the LSP is the Bino-like neutralino. In addition, in order for the sizable contamination of the Higgsino component into the lightest neutralino, the $\mu_H$ parameter is required to be relatively small. Consequently, the lightest chargino as well as the second and third lightest neutralinos become Higgsino like. Study of their properties will give important information for the calculation of the thermal relic density of the LSP.

Once the Higgsino-like charginos and neutralinos are produced at the ILC, their properties as well as the mass of the LSP can be precisely determined. For example, from the threshold scan at $\sqrt{s} \sim 2m_{\chi_1^\pm}$, we can determine the chargino mass from the process $e^+e^- \to \chi_1^+\chi_1^-$. In the linear collider, neutralinos can be also produced. Since the neutralinos are Majorana particle, pair productions of the identical neutralinos are suppressed at the threshold region. The process $e^+e^- \to \chi_2^0\chi_2^0$ can have, however, sizable cross section. From the threshold scan of this process, we can determine the combination $m_{\chi_2^0} + m_{\chi_3^0} \equiv 2\tilde{m}_{\chi_2^0}$. At the ILC, errors in the measurements of the masses are expected to be mostly from the detector resolutions [27]. For example, it was pointed out that, for some choice of the SUSY parameters, masses of the charginos can be determined using $e^+e^-$ colliders with the errors of $\sim 50$ MeV by the threshold scan. In addition, from the energy distribution of the decay products of the chargino and neutralinos, the mass of the LSP is also determined with the uncertainty of $\sim 50$ MeV. Although these results are for the case of Wino-like chargino and neutralino, we expect that three mass parameters (i.e., $m_{\chi_1^\pm}$, $m_{\chi_1^0}$, and $\tilde{m}_{\chi_{23}^0}$) are accurately measured once $\chi_1^\pm$, $\chi_2^0$, and $\chi_3^0$ become kinematically accessible at the ILC. Since $\chi_1^0$ is Bino-like while $\chi_2^0$ (as well as $\chi_3^0$ and $\chi_0^0$) are Higgsino-like, we can constrain $m_{G1}$ and $\mu_H$ from the measurements of $m_{\chi_1^0}$ and $m_{\chi_1^\pm}$ (or from the masses of other Higgsino-like neutralinos).

Thus, if the Higgsino-like chargino ($\chi_1^\pm$) and neutralinos ($\chi_2^0$ and $\chi_3^0$) are produced at the ILC, three constraints will be obtained on the MSSM parameters $m_{G1}$, $m_{G2}$, $\mu_H$, and
\[ \tan \beta. \] Of course, if, for example, the heavier chargino can become kinematically accessible at the ILC, we can impose four constraints on the MSSM parameters so all the parameters relevant for the calculation of the \( \Omega_{\text{LSP}} \) can be in principle reconstructed. Even without producing the Wino-like chargino and neutralino, however, interesting bound on \( \Omega_{\text{LSP}} \) can be obtained. To see this, we can perform the following analysis. Let us imagine a situation where \( m_{\tilde{\chi}_1^\pm}, m_{\tilde{\chi}_1^0}, \) and \( \bar{m}_{\tilde{\chi}_2^0} \) are well measured at the ILC. Using these quantities, we impose three constraints on the four underlying parameters and determine \( m_{G1}, \mu_H, \) and \( \tan \beta \) as functions of \( m_{G2} \). In the determination of \( m_{G1} \) and \( \mu_H \), in fact, there are four possible choices of their signs: \((\text{sign}(m_{G1}), \text{sign}(\mu_H)) = (+, +), (+, -), (-, +), \) and \((-, -).\) \#5 Effects of the signs of \( \mu_H \) and \( m_{G1} \) are quite different. In order to see how the reconstructed relic density depends on \( m_{G2} \) and \( \mu_H, \) here we consider the case where the sign of the reconstructed \( m_{G1} \) is the same as that of the underlying one; effects of \( \text{sign}(m_{G1}) \) will be discussed later. Once we reconstruct \( m_{G1}, \mu_H, \) and \( \tan \beta, \) we calculate the relic density of the LSP as a function of \( m_{G2}, \) which we call

\[ \hat{\Omega}_{\text{LSP}}(m_{G2}; m_{\tilde{\chi}_1^\pm}, m_{\tilde{\chi}_1^0}, \bar{m}_{\tilde{\chi}_2^0}). \]

In Figs. 4 and 5 we plot \( \hat{\Omega}_{\text{LSP}}(m_{G2}; m_{\tilde{\chi}_1^\pm}, m_{\tilde{\chi}_1^0}, \bar{m}_{\tilde{\chi}_2^0}) \) as a function of \( m_{G2} \) for two different choices of the underlying Points: Point 1 with \( m_{G1} = 144 \text{ GeV}, m_{G2} = 300 \text{ GeV}, \) \( \mu_H = 200 \text{ GeV}, \) and \( \tan \beta = 10, \) and Point 2 with \( m_{G1} = 240 \text{ GeV}, m_{G2} = 500 \text{ GeV}, \) \( \mu_H = 307 \text{ GeV}, \) and \( \tan \beta = 10. \) The lines have endpoints; this is due to the fact that, when \( m_{G2} \) becomes too large or too small, there is no value of \( \tan \beta \) which consistently reproduces the observed mass spectrum. To demonstrate this, we also showed the points where \( \tan \beta \) takes several specific values. In the figures, the results for the cases with positive and negative \( \mu_H \) are shown.

As one can see, the case with negative \( \mu_H \) may give large uncertainty in the reconstructed \( \Omega_{\text{LSP}}. \) If \( \mu_H < 0, \) however, smaller \( m_{G2} \) is required than in the case of positive-\( \mu_H \) in order to reproduce the observed mass spectrum, as shown in the figures. Then, \( m_{\tilde{\chi}_2^\pm} \), for example, may become smaller than the experimental bound from the negative search for the \( \chi_1^0 \chi_2^\pm \) production process. For Point 1 (Point 2), \( \sqrt{s} \gtrsim 480 \text{ GeV} \) (750 GeV) is enough to exclude the \( \mu_H < 0 \) case. Thus, in the following, we assume that this is the case and neglect the uncertainty from the \( \mu_H < 0 \) case. Then, even without any further constraint on \( m_{G2}, \) relic density of the LSP can be determined within a factor of \( \sim 2 \) or smaller. Of course, if the Wino-like chargino and neutralino become kinematically accessible at the ILC, we can determine \( m_{G2} \) and hence the relic density of the LSP can be determined with a great accuracy. \#6

So far, we have not considered effects of the sign of \( m_{G1}. \) Unfortunately, \( \Omega_{\text{LSP}} \) depends on the sign of \( m_{G1} \) although the determination of \( \text{sign}(m_{G1}) \) seems challenging. For the

\#5 To be more precise, these signs are the relative signs between \( m_{G1} \) and \( m_{G2} \) or \( \mu_H \) and \( m_{G2}. \) We assume that the gaugino masses and \( \mu_H \) are real in order to avoid constraints from CP violations.

\#6 In such a case, radiative corrections to the masses and cross sections should be also calculated to reduce theoretical uncertainties [29].
Figure 4: $\Omega_{\text{LSP}}(m_{G2}, m_{\chi^\pm}, m_{\chi^0}, \bar{m}_{\chi^0})$ as a function of $m_{G2}$, where $m_{\chi^\pm}$, $m_{\chi^0}$, and $\bar{m}_{\chi^0}$ are fixed by the underlying values for Point 1 with positive $\mu_H$ (solid) and negative $\mu_H$ (dashed). Marks on the figure indicate the points with $\tan \beta = 2, 3, 5, 10, 50$.

case with negative $m_{G1}$, $\Omega_{\text{LSP}} h^2$ varies from 0.9 to 7.4 (Point 1) and from 0.2 to 6.7 (Point 2). If sign$(m_{G1})$ is undetermined, thus, two-fold ambiguity will remain. However, experimental determination of sign$(m_{G1})$ may be possible [30]. In addition, if the GUT relation among the (absolute values of) gaugino masses is experimentally confirmed, it will give another hint of the signs of the gaugino masses.

5 SUMMARY

Here, it is discussed that the LCSUGRA framework can provide an interesting cosmological scenario which explains the origin of the baryon asymmetry of the universe as well as the identity of the dark matter without conflicting the gravitino problem. In this scenario, the gravitino acquires the mass of a few TeV and hence the reheating temperature after the inflation can be as high as $O(10^6 \text{ GeV})$. With such a reheating temperature, the baryon asymmetry of the universe can be generated by the non-thermal leptogenesis. In addition, the lightest neutralino becomes the LSP in this scenario and its relic density can be consistent with the dark matter density determined by the WMAP. Importantly, the relic density of the LSP can be reconstructed once the superparticles are produced at the ILC.
Figure 5: Same as Fig. 4 except for Point 2.

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References

[1] S. Weinberg, Phys. Rev. Lett. 48 (1982) 1303.

[2] M. Kawasaki, K. Kohri and T. Moroi, arXiv:astro-ph/0402490.

[3] M. Kawasaki, K. Kohri and T. Moroi, arXiv:astro-ph/0402490; Phys. Rev. D 71 (2005) 083502.

[4] K. Kohri, T. Moroi and A. Yotsuyanagi, arXiv:hep-ph/0507245.

[5] For recent studies on the effects of the unstable gravitino on the BBN other than 2 3 4, see, for example, R. H. Cyburt, J. R. Ellis, B. D. Fields and K. A. Olive, Phys. Rev. D 67 (2003) 103521; K. Jedamzik, Phys. Rev. D 70 (2004) 063524;

[6] K. Enqvist and M. S. Sloth, Nucl. Phys. B 626 (2002) 395; D. H. Lyth and D. Wands, Phys. Lett. B 524 (2002) 5; T. Moroi and T. Takahashi, Phys. Lett. B 522 (2001)
[21] M. Ibe, T. Moroi and T. Yanagida, arXiv:hep-ph/0502074.

[22] J. L. Feng, K. T. Matchev and F. Wilczek, Phys. Lett. B 482 (2000) 388
arXiv:hep-ph/0004043.

[23] D. N. Spergel et al., Astrophys. J. Suppl. 148 (2003) 175.
[24] F.E. Paige, S.D. Protopescu, H. Baer, and X. Tata, \texttt{arXiv:hep-ph/0312045}.

[25] S. Matsumoto \textit{et al.} [JLC Group], “JLC-1,” KEK Report 92-16 (1992).

[26] S. Kuhlman \textit{et al.} [The NLC ZDR Design Group and The NLC Physics Working Group], “Physics and Technology of the Next Linear Collider,” BNL 52-502 (1996).

[27] J. A. Aguilar-Saavedra \textit{et al.} [ECFA/DESY LC Physics Working Group], \texttt{arXiv:hep-ph/0106315}.

[28] T. Moroi, Y. Shimizu and A. Yotsuyanagi, \texttt{arXiv:hep-ph/0505252}.

[29] B. C. Allanach, G. Belanger, F. Boudjema and A. Pukhov, JHEP \textbf{0412} (2004) 020.

[30] S. Y. Choi, J. Kalinowski, G. Moortgat-Pick and P. M. Zerwas, Eur. Phys. J. C \textbf{22} (2001) 563 [Addendum-ibid. C \textbf{23} (2002) 769].