Dark Matter and Baryon Asymmetry of the Universe in Large-Cutoff Supergravity

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Abstract

We propose a consistent scenario of the evolution of the universe based on the large cutoff supergravity (LCSUGRA) hypothesis of supersymmetry breaking, where the gravitino and sfermion become as heavy as \( \sim O(1 - 10 \text{ TeV}) \). With such a heavy gravitino, baryon asymmetry of the universe can be generated by the non-thermal leptogenesis via an inflaton decay without conflicting the serious gravitino problem. 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.
1 Introduction

In a recent paper [1] Izawa and two of us (M.I. and T.Y.) have proposed a large-cutoff hypothesis in supergravity. In this hypothesis all higher dimensional operators such as quartic terms in the Kähler potential at the GUT scale $M_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV (or at the reduced Planck scale $M_G \simeq 2.4 \times 10^{18}$ GeV) are suppressed by a large cutoff $M_* \sim O(4\pi M_G)$. Then, 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 [2] due to the fact that the universality of the scalar masses at the GUT scale is guaranteed in this scenario. In [1], it was shown that the 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 $M_*$ is a reflection of a more fundamental physics beyond the GUT scale.

In this letter we study the cosmology of the LCSUGRA scenario. In SUGRA, there is a serious cosmological problem, that is the gravitino problem. If the gravitino is unstable, it has a long lifetime and decays after the big-bang nucleosynthesis (BBN) for an interesting range of the gravitino mass, $m_{3/2} \sim 100 \text{GeV} - 10 \text{TeV}$. The decay products destroy light elements produced by the BBN and hence the primordial abundance of the gravitino is constrained from above to keep the success of the BBN. This leads to an upper bound on the reheating temperature $T_R$ after inflation, since the abundance of the gravitino is proportional to $T_R$. The recent detailed analysis derived a stringent upper bound $T_R < 10^{6-7}$ GeV when the gravitino decay has hadronic modes [3]. This upper bound is much lower than the temperatures required for most of thermal barygenesis including the leptogenesis [4]. There have been proposed various solutions to the above problem and, among them, we consider that the non-thermal leptogenesis via an inflaton decay [5] is the most interesting and plausible.\footnote{The Affleck-Dine (AD) baryogenesis [6] is an interesting scenario, but this does not work in the LCSUGRA, since the Kähler potential between inflatons and the AD fields is suppressed and the AD fields acquire the Hubble-induced masses which set the AD fields at the origin during the inflation [4].} Importantly, even if the right-handed neutrinos

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\footnote{The Affleck-Dine (AD) baryogenesis [6] is an interesting scenario, but this does not work in the LCSUGRA, since the Kähler potential between inflatons and the AD fields is suppressed and the AD fields acquire the Hubble-induced masses which set the AD fields at the origin during the inflation [4].}
are non-thermally produced from the decay of the inflaton, success of the baryogenesis is not automatic, as we will see.

The main conclusion of this letter is that the LCSUGRA scenario is also advantageous in cosmology. In the next section, we show that a general argument on the inflaton-decay scenario of leptogenesis gives us a lower bound on the reheating temperature as \( T_R \gtrsim 1 \times 10^6 \text{ GeV} \) \cite{8}. Together with the constraint from the BBN \cite{3} we find the lower bound on the gravitino mass \( m_{3/2} \gtrsim 4 \text{ TeV} \). Importantly, such heavy gravitino is a natural outcome of the LCSUGRA scenario. In section 3, we show that the LCSUGRA with such a large gravitino mass indeed explains the observed dark matter density. In particular, the result of the present analysis shows that the universal scalar mass of SUSY breaking, \( m_0 \), or equivalently the gravitino mass \( m_{3/2} \) is \( 4 - 10 \text{ TeV} \) which is in a special area of so-called focus point region. Remarkably, such a parameter region is also required for the success of the non-thermal leptogenesis via the inflaton decay as mentioned above. The last section is devoted to discussion.

2 Leptogenesis via inflaton decay

Let us first consider the leptogenesis in the framework of LCSUGRA. Here, for simplicity, we consider the case where only the lightest right-handed neutrino \( N_1 \) contributes to the lepton asymmetry, assuming mass hierarchy among the right-handed neutrinos. \( N_1 \) may decay into \( H_u + \ell \) or \( H_u^* + \ell^* \) where \( H_u \) is the up-type Higgs and \( \ell \) the lepton doublet. These two decay channels have different branching ratios when CP conservation is violated. Interference between tree-level and one-loop diagrams generates lepton-number asymmetry \cite{4,9,10} as

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

where \( m_{\nu_3} \) is the heaviest (active) neutrino mass. Using \( 3 \times 3 \) neutrino Yukawa matrix \( h \), the effective CP-violating phase \( \delta_{\text{eff}} \) is given by

\[
\delta_{\text{eff}} = \frac{\text{Im} \left[ h_{13}^2 + \frac{m_{\nu_2}}{m_{\nu_3}} h_{12}^2 + \frac{m_{\nu_1}}{m_{\nu_3}} h_{11}^2 \right]}{|h_{13}|^2 + |h_{12}|^2 + |h_{11}|^2}.
\]

In deriving above expressions, the seesaw mass formula \cite{11} has been used.
With non-vanishing $\epsilon$ parameter, lepton-number asymmetry is generated with the out-of-equilibrium decay of $N_1$. Since $\epsilon$ is proportional to $M_1$, $N_1$ is required to be heavy enough to generate sufficient amount of baryon number density. If we consider the case where the right-handed neutrino is thermally produced, for example, we obtain $M_1 \gtrsim 10^{9-10}$ GeV [12]. Since $T_R \gtrsim M_1$ in this case, we should conclude that the thermal leptogenesis requires too high reheating temperature to avoid the gravitino problem in the conventional SUGRA models.

If the right-handed neutrino is non-thermally produced by the decay of the inflaton $\Phi$ as $\Phi \rightarrow N_1 + N_1$ ($\Phi \rightarrow \tilde{N}_1 + \tilde{N}_1$) [5], the situation changes. In this case, Boltzmann equations for the $B-L$ number density $n_{B-L}$ and other quantities (i.e., the energy density of the inflaton $\rho_\Phi$ and that of radiation $\rho_{\text{rad}}$) are given by

\begin{align*}
\frac{dn_{B-L}}{dt} + 3Hn_{B-L} &= 2\epsilon\Gamma_\Phi m_\Phi^{-1} \rho_\Phi, \\
\frac{d\rho_\Phi}{dt} + 3H\rho_\Phi &= -\Gamma_\Phi \rho_\Phi, \\
\frac{d\rho_{\text{rad}}}{dt} + 4H\rho_{\text{rad}} &= \Gamma_\Phi \rho_\Phi,
\end{align*}

where $m_\Phi$ is the inflaton mass, $\Gamma_\Phi$ the decay rate of the inflaton. Notice that the factor 2 in the right-hand side of Eq. (3) is due to the fact that two right-handed neutrinos are produced by the decay of single $\Phi$. Here and hereafter, we adopt a mild assumption that the inflaton potential can be well approximated by the parabolic one at the last stage of the reheating.

The inflaton decays when the expansion rate of the universe becomes comparable to the decay rate $\Gamma_\Phi$; we define the reheating temperature as

$$T_R \equiv \left( \frac{10}{g_*\pi^2 M_G^2 \Gamma_\Phi^2} \right)^{1/4},$$

where $g_*$ is the effective number of the massless degrees of freedom. (In our numerical study, we use $g_* = 228.75$.) As well as the entropy production, generation of the $B-L$ asymmetry is most effective when $H \sim \Gamma_\Phi$. Consequently, we obtain the relation between $n_{B-L}$ and the entropy density $s$ as

$$\frac{n_{B-L}}{s} = \kappa \epsilon \frac{T_R}{m_\Phi} = 9.9 \times 10^{-11} \times \kappa \left( \frac{T_R}{10^6 \text{GeV}} \right) \left( \frac{2M_1}{m_\Phi} \right) \left( \frac{m_{\nu_3}}{0.05 \text{eV}} \right) \delta_{\text{eff}},$$
where $\kappa$ is a constant of $O(1)$. In the second equality, we have taken $\langle H_u \rangle \simeq 174$ GeV. Qualitative behavior of Eq. (7) can be easily understood using the instantaneous-decay approximation: $m_\phi n_{B-L}(T_R) \sim \epsilon [\rho_{\phi}]_{\text{decay}} \sim \epsilon g_* T_R^4$ and $s(T_R) \sim g_* T_R^3$. Using the order-of-magnitude relation $n_{B-L} \sim n_B$ after the spharelon transition, one can see that, in order to generate the enough baryon asymmetry of the universe suggested by the WMAP [13]

$$\frac{n_B}{s} \simeq 0.9 \times 10^{-10}, \quad (8)$$

$T_R$ should be higher than $\sim 10^6$ GeV.

Since this lower bound is very close to the upper bound on $T_R$ from the gravitino problem, we have performed a careful calculation of the resultant baryon number in this scenario. We have numerically solved Eqs. (3) – (5) from the cosmic time $t \ll \Gamma^{-1}_\Phi$ to $t \gg \Gamma^{-1}_\Phi$. Then, we have calculated the entropy density at $t \gg \Gamma^{-1}_\Phi$ using the relations

$$\rho_{\text{rad}} = \frac{\pi^2}{30} g_* T_R^4, \quad s = \frac{2\pi^2}{45} g_* T_R^3. \quad (9)$$

Taking the ratio of $n_{B-L}$ to $s$, we have obtained the $\kappa$ parameter in Eq. (7) as

$$\kappa \simeq 2.44. \quad (10)$$

Using $n_B = \frac{28}{79} n_{B-L}$ [14], baryon-to-entropy ratio is given by

$$\frac{n_B}{s} \simeq 8.2 \times 10^{-11} \times \left( \frac{T_R}{10^6\text{GeV}} \right) \left( \frac{2M_1}{m_\phi} \right) \left( \frac{m_\nu_3}{0.05\text{eV}} \right) \delta_{\text{eff}}. \quad (11)$$

Combining this with Eq. (8), we derive a constraint on the reheating temperature,

$$T_R > 1 \times 10^6 \text{GeV}, \quad (12)$$

for the neutrino mass suggested from the atmospheric neutrino oscillation experiments, $m_{\nu_3} \simeq \sqrt{\Delta m^2_{23}} \simeq 0.05$ eV. Here, we used $M_1 = \frac{1}{2} m_\phi$, which is the maximal possible value of $M_1$ so that the decay process $\Phi \rightarrow N_1 + N_1$ is kinematically allowed. From the point of view of the gravitino problem, such reheating temperature is quite dangerous.

It is notable that, in the LCSUGRA scenario where the gravitino is quite heavy, the present non-thermal leptogenesis scenario becomes viable in a wide parameter region. Indeed, the recent detailed analysis of the gravitino problem suggests that $T_R \sim 1 \times$
$10^6$GeV is consistent with the BBN constraints if the gravitino is heavier than 4 TeV \cite{3}.

In the LCSUGRA scenario, such a large gravitino mass can be realized without conflicting the naturalness of the electroweak symmetry breaking (the focus point mechanism). In addition, as we will discuss in the next section, when $m_{3/2} \gtrsim 4$ TeV, the predicted relic density of the LSP well agrees with the currently observed dark matter density.

Before closing this section, some discussion on the constraints from the gravitino problem may be relevant. For the gravitino mass we are interested, overproduction of D gives the most stringent upper bound on $T_R$. Ref. \cite{3} used averaged value of the observational values of D/H to set the bounds. If we adopt the largest value of observed D/H, upper bound on $T_R$ becomes larger by factor $2 - 3$. If so, the non-thermal leptogenesis may be consistent with the gravitino mass smaller than 4 TeV. However, notice that, in deriving Eq. (12), we considered the extreme case where $M_1 = \frac{1}{2}m_{\Phi}$ and $\delta_{\text{eff}} = 1$; in particular, if $M_1$ is smaller, higher reheating temperature is required to generate large enough baryon asymmetry.\footnote{In addition, the primary purpose of Ref. \cite{3} was to derive conservative constraints taking account of only the effects which are well understood. Thus, the hadrodissociations induced by the colored superparticles emitted from the gravitino decay were not included since such processes are very hard to estimate. Once they are included, the upper bound on $T_R$ will become severer.} Thus, the non-thermal leptogenesis scenario is severely constrained from the point of view of the gravitino problem and the large gravitino mass is preferred to solve the conflict.

\section{Dark matter in LCSUGRA}

In the previous section, we have seen that the LCSUGRA scenario is good for the viable scenario of the non-thermal leptogenesis. In this section, we discuss that the LCSUGRA scenario can provide, in the reasonable portion of the parameter space, the dark-matter density consistent with the WMAP result.

Before discussing the behavior of the dark-matter density, we first summarize the free parameters in our analysis. In the LCSUGRA scenario, all the higher-dimensional operators are assumed to be suppressed by the inverse powers of the cutoff scale $M_*$ much larger than the reduced Planck scale. Thus, as usual, all scalar particles obtain a universal SUSY breaking mass $m_0$ through the effectively minimal Kähler potential,\footnote{There may be non-diagonal corrections to scalar mass squared from higher dimensional terms in}
which is equal to the gravitino mass $m_{3/2}$. The universal gaugino mass $m_{1/2}$ is obtained through higher dimensional operators in the superpotential which is suppressed by $M_*$ (where we assume unification of the $SU(3) \times SU(2) \times U(1)$ gauge group). The SUSY-breaking (scalar) couplings, so called $A$ parameters, are also generated through higher dimensional operators in the superpotential and they are expected to be of order the gaugino mass.\footnote{If the vacuum expectation value (VEV) of the field $Z$ in the SUSY breaking sector is large as the reduced Planck scale $M_{G_{11}}$, we have the $A$ terms of the order the gravitino mass $m_{3/2}$. However, most of the dynamical SUSY breaking models suggest $\langle Z \rangle \ll M_{G_{11}}$. Therefore, we consider $|A| \ll m_{3/2}$.} In the present analysis, we simply take $A = 0$, since their contributions to the following discussions are sub-dominant as long as $|A| \ll m_0$. In addition, we can obtain $\mu$-parameter via the Giudice-Masiero mechanism \cite{17}. Then, the $\mu$-parameter is also suppressed by the large cutoff scale and becomes of the same order of the gaugino mass. Notice that the hierarchy $\mu \ll m_0$ is consistent with the radiative electroweak symmetry breaking due to the focus point mechanism \cite{2}. Notably, in this case, the so-called $B$-parameter is not a free parameter and it is equal to the gravitino mass (i.e., $B_0 = m_0$) for $m_0 \gg A$. We consider that all of the above parameters are given at the GUT scale in the present analysis. Then, we have three parameters $m_0, m_{1/2}, \mu$ ($m_0 \gg m_{1/2}, \mu$) at the GUT scale.

At the electroweak breaking scale, the gaugino masses $m_1, m_2$ and $m_3$ can be determined by solving the RG equations, and are roughly given by $m_1 \simeq 0.4 m_{1/2}, m_2 \simeq 0.8 m_{1/2}$ and $m_3 \simeq 2.8 m_{1/2}$, respectively. Notice that the RG evolution of $\mu$ from the electroweak scale to the GUT scale is negligible, and hence, in the followings we take $\mu$ as the value at the electroweak breaking scale. We see that all the scalar mass parameters other than the one of the up-type Higgs scalar are much larger than gaugino masses and $\mu$, thus, the gauginos and Higgsinos are much lighter than sfermions at the electroweak scale. Thus, the LSP is always lightest neutralino which is stable by $R$-parity conservation. Such lightest neutralino becomes a good candidate of the dark matter as we see below.

Postulating that the LSP becomes the dark matter, we calculate its relic density. Since heavier CP-even and CP-odd Higgs bosons as well as all the sfermions become much heavier than the neutralino, the leading processes in the neutralino pair annihilation are through $s$-channel exchanges of $Z$ and CP-even Higgs bosons and also through the $t$-
channel exchanges of the charginos and neutralinos \cite{18}. In this case, \( \tan \beta \) dependence of the relic density is small since the annihilation processes do not strongly depend on \( \tan \beta \) except for the CP-even Higgs exchanges.\(^5\) On the other hand, since we consider the situation \(|\mu| \sim m_{1/2}\), the lightest neutralino may have a sizable fraction of the Higgsino components \cite{19} and the relic density is sensitive to \( \mu \) and \( m_{1/2} \) through the Higgsino fraction of the lightest neutralino.

In Fig. 1, we show the parameter region in the \((\mu, m_2)\) plane where the neutralino density is consistent with WMAP result \cite{13}, \( \Omega_{\rm DM} h^2 = 0.1126^{+0.0161}_{-0.0181} \). In this computation, we have used micrOMEGAs 1.3.0 code \cite{20} which includes all the possible co-annihilation effects. In the figure, the red points are for \( \tan \beta = 5 \) and the blue ones for \( \tan \beta = 30 \). As we mentioned, the result is insensitive to \( \tan \beta \), and we find that the relic density is mainly determined by the parameters \( \mu \) and \( m_2 \). It should be noted that the result in \((\mu, m_2)\) plane is free from uncertainties arising from the analysis in the focus point region which we will discuss later. Especially, for \( m_2 \gtrsim 350 \) GeV, \( \tan \beta \) dependence of the relic density is negligible, since the neutralino becomes heavy enough so that the annihilation process is dominated by the channel into top-pair via Z-boson exchange.\(^6\) The WMAP value of the relic density is realized on the line \( \mu \simeq 0.6 m_2 \), and by comparing the bino mass parameter \( m_1 \simeq 0.5 m_2 \), we find that the neutralino has a significant Higgsino fraction which makes annihilation more efficient.

We next reinterpret the above results by using the LCSUGRA parameters. In the LCSUGRA, \( m_0 \) and \( \tan \beta \) are uniquely determined for a given value of \((\mu, m_2)\), since the SUSY-breaking Higgs mixing parameter \( B \) is determined as \( B_0 = m_0 \) at the GUT scale.

In order to understand how \( \tan \beta \) depends on the LCSUGRA parameters, it is instructive to see the tree-level condition for the electroweak symmetry breaking although important radiative corrections to the Higgs potential are taken into account in our numerical calculations. The tree-level minimization conditions of the effective Higgs potential

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\(^5\)\( \tan \beta \equiv v_u/v_d \) is the ratio of the two VEV’s of the neutral Higgs fields, \( v_{u,d} \equiv \langle H^0_{u,d} \rangle \).

\(^6\)The processes into other lighter fermion pairs are much suppressed, since the amplitudes are proportional to the fermion masses which are required for helicity flip in s-wave processes.
Figure 1: Dark matter constraints after WMAP ($\Omega_{DM} h^2 = 0.1126^{+0.0161}_{-0.0181}$) in the ($\mu, m_2$) plane for (a) $m_{t\text{op}} = 174$ GeV and (b) $m_{t\text{op}} = 178$ GeV. The reds and blue points correspond to $\tan \beta = 5$ and $\tan \beta = 30$, respectively. The black shaded regions are excluded by the chargino mass limit, $m_{\chi^\pm} \geq 104$ GeV [21].

are given by

$$\frac{1}{2}m_Z^2 = \frac{m_{H_u}^2 - m_{H_d}^2 \tan^2 \beta}{\tan^2 \beta - 1} - |\mu|^2, \quad (13)$$

$$B\mu = \frac{\sin 2\beta}{2}(m_{H_u}^2 + m_{H_d}^2 + 2|\mu|^2), \quad (14)$$

where $m_{H_u,d}^2$ denote the soft masses squared of up-type and down-type Higgs bosons, and all the parameters are evaluated at the electroweak scale. Since all the parameters in Eqs. (13) and (14) are determined from three input parameters ($m_0$, $m_{1/2}$, and $\mu_0$, since $B_0 = m_0$), we can determine $(m_0, \tan \beta)$ for a given ($\mu, m_2$) by solving above conditions.

In order to obtain explicit relations between $(m_0, \tan \beta)$ and ($\mu, m_2$), we express the parameters in Eqs. (13) and (14) in terms of the input parameters as

$$m_{H_u}^2 = a_u m_0^2 + b_u |m_{1/2}|^2, \quad (15)$$

$$m_{H_d}^2 = a_d m_0^2 + b_d |m_{1/2}|^2, \quad (16)$$

$$B \simeq B_0 = m_0, \quad (17)$$

\[7\text{In the actual numerical calculation, we use the running parameters at the typical stop mass scale, where the one-loop corrections to the Higgs potential tend to be small.}\]
where the coefficients $a_i$ and $b_i$ ($i = u, d$) are scale-dependent functions of dimensionless gauge and Yukawa coupling constants. Here, $B \simeq B_0$ is the consequence of the hierarchical spectrum ($m_0 \gg m_{1/2}, A$) of the LCSUGRA scenario. As discussed in Refs. [2], at the electroweak scale, $a_u$ is of order $10^{-2} - 10^{-1}$ while $a_d \simeq 1$ for $m_{\text{top}} \simeq 170 - 180$ GeV, which comes from universality of the scalar masses. The parameter $b_d$ becomes positive ($O(0.1)$) at the electroweak scale, which mainly comes from positive wino and bino contributions. On the other hand, $b_u$ becomes negative ($O(1)$) at the electroweak scale because of the positive gluino contributions to the stop mass squared which affects $m_{H_u}^2$ through top Yukawa interaction. Substituting Eqs. (15), (16) and (17) into the minimization conditions, we obtain explicit relation between $(m_0, \tan \beta)$ and $(\mu, m_2)$ as

$$
\mu \simeq \frac{a_d}{\tan \beta} m_0, 
$$

(18)

$$
m_2 \simeq 0.8 m_{1/2} \simeq 0.8 \sqrt{-\frac{a_u}{(b_u + b_d)}} m_0. 
$$

(19)

These relations show that, when $a_u \sim 10^{-2}$, the LCSUGRA scenario is consistent with the electroweak symmetry breaking for $\tan \beta \gtrsim 10$. By using Eqs. (18) and (19), we obtain the WMAP constraint in the $(\tan \beta, m_0)$ plane from the one in the $(\mu, m_2)$ plane.

With a more detailed numerical calculations, we relate the input parameters to the parameters at the electroweak scale, and derive the WMAP constraint in the $(\tan \beta, m_0)$ plane. The results are shown in Fig. 2 which is converted from the one in Fig. 1

8In Fig. 2, we reassure that the heavier CP-even, CP-odd Higgs bosons and all sfermions are much heavier than the neutralino, which is important to obtain the results in Fig. 1.

9It should be noted that, the lines in the figures show rough fitting of the results, since the code becomes somewhat unstable for $m_0 \gg m_{1/2}$.
for \( m_{\text{top}} = 174 \text{ GeV} \) and \( m_0 \gtrsim 4 \text{ TeV} \) for \( m_{\text{top}} = 178 \text{ GeV} \).

As we have seen in the previous section, we find that the non-thermal leptogenesis and the gravitino problem require \( m_{3/2} \gtrsim 4 \text{ TeV} \). Using the fact that \( m_0 = m_{3/2} \) in LCSUGRA, we can see that the LCSUGRA scenario with a relatively large gravitino mass provides a consistent cosmological scenario which explains both baryon and dark matter densities in the present universe.

Finally, we comment on the uncertainties in our results which comes from the technical difficulties in the precise calculation of the superparticle mass spectrum in LCSUGRA. Since some of the superparticles become much heavier than the electroweak-symmetry breaking scale in the focus point parameter region, an accurate determination of the mass spectrum requires careful treatments of various corrections. As a result, the numerical results are somewhat different from code to code \([24, 25]\) and our results may be affected by such uncertainty. To see its effect, we have repeated our analysis with SOFTSUSY 1.9 code \([26]\); we have find that the changes of \( m_0 \) and \( \tan \beta \) for a given set of \((\mu, m_2)\) are at most \( \sim 30\% \). Thus, we believe that our main conclusion, suggesting the LCSUGRA scenario with a relatively large gravitino mass, does not change although our quantitative estimations may be slightly changed.
4 Discussion

In this letter, we have proposed a consistent cosmological hypothesis based on the LC-SUGRA scenario of the supersymmetry breaking. In the LCSUGRA hypothesis, the gravitino (as well as the sfermions) becomes as heavy as a few TeV without conflicting the naturalness of the electroweak symmetry breaking. If so, baryon asymmetry of the universe can be explained by the non-thermal leptogenesis via an inflaton decay without spoiling the success of the big-bang nucleosynthesis. We have also shown that, in the parameter region suggested from the successful non-thermal leptogenesis, the relic density of the LSP explains naturally the dark matter density required from the WMAP results.

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