Cosmological Moduli Problem and Oscillating Inflation in Gauge-Mediated Supersymmetry Breaking

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Abstract

We investigate cosmological moduli problem in the gauge-mediated supersymmetry breaking (GMSB). A mini-inflation (oscillating inflation) takes place when a scalar field corresponding to the flat direction in GMSB oscillates along the logarithmic potential induced by the gauge-mediation mechanism. It is shown that this oscillating inflation can sufficiently dilute the relic abundance of the string moduli for some ranges of parameters in the GMSB models.
1 Introduction

Supersymmetry (SUSY) is an attractive candidate beyond the standard model because it stabilizes the electroweak scale against the radiative corrections. However, SUSY is not exact symmetry and should be broken since we have not detected degenerate superpartners of ordinary particles. Dynamics of the SUSY breaking has not been clear and been one of the most important open questions left to us. So far many ideas have been proposed as the SUSY breaking mechanism. Among them the gauge-mediated SUSY breaking (GMSB) mechanism is very attractive since it gives a solution of the SUSY flavor problem.

However, the models of GMSB are faced with a serious cosmological problem. One of the consequences of GMSB is the existence of a light stable gravitino. Such a light gravitino is produced after the primordial inflation and if its mass is \( m_{3/2} \gtrsim 1 \text{ keV} \), the reheating temperature of the inflation should be low enough [2].

Furthermore, when one considers GMSB in the framework of the superstring theories, more severe cosmological difficulty, “cosmological moduli problem”, arises. As a general consequence of the superstring theories, light moduli particles \( \phi \) appear and their masses \( m_\phi \) are expected of the order of the gravitino mass [3]. With such small masses \( m_\phi \simeq m_{3/2} \), since they have only the gravitationally suppressed interaction, they are stable within the age of our universe (\( \sim 10^{17} \text{ sec} \)) and easily overclose the present critical density of the universe.

The thermal inflation model proposed by Lyth and Stewart [4] gives a solution of this problem. In fact it was shown in Ref. [5] that the above overclosure problem could be solved if one assumed the thermal inflation. However, it was pointed out [6] that if \( m_\phi \gtrsim 100 \text{ keV} \), a more stringent constraint comes from the observation of the present cosmic x(\( \gamma \))-ray backgrounds, which excludes the moduli with mass \( 100 \text{ keV} < m_\phi < 1 \text{ GeV} \) even if the universe experienced the thermal inflation [7, 8].

Recently, Moroi [9] reported that the SUSY breaking field, which has a very flat potential, gives another dilution mechanism of the relic abundance of the moduli. This is because an accelerating universe (say an oscillating inflation [10]) takes place when the scalar field corresponding to the flat direction oscillates along the logarithmic potential induced by the GMSB mechanism. In Ref. [9] the specific model of GMSB was analyzed and was shown that the moduli is significantly diluted. However, during the oscillating inflation, the minimum of the moduli potential deviates from the true minimum through the additional SUSY breaking effect due to the large vacuum energy of the inflaton field. After the inflation ends, the moduli field begins to move toward the true minimum and hence the (secondary) oscillation starts even if the amplitude of the original coherent oscillation of the moduli sufficiently decreases during the inflation. However, this deviation of the minimum was not properly taken into account and the moduli density coming from the secondary oscillation was totally neglected in Ref. [9].

In this letter we consider the cosmological moduli problem assuming that the oscillating inflation originates in the minimal SUSY standard model (MSSM) flat direction [11] and taking account of the secondary oscillation. In the following, we discuss the oscillating inflation which is a general consequence of the flat directions in the GMSB models and examine how dilution mechanism works. Then we study whether this oscillating inflation could solve the cosmological difficulties of the moduli or not. Finally, we summarize our conclusions.

\[ ^1 \text{Our discussion applies to a general flat direction of the GMSB models.} \]
2 Oscillating Inflation and Entropy Production

In the generic SUSY models, there appears the flat direction, say $X$, in the scalar potential. However, the effects of SUSY breaking lift the flat direction. In most models of GMSB, SUSY is broken in the dynamical SUSY-breaking sector and its effects are transferred into the messenger sector. The messenger sector contains a gauge singlet multiplet $S$ which is supposed to have an $A$-component vacuum expectation value (vev) $\langle S \rangle$ and also have a $F$-component vev $\langle F_S \rangle$. Then SUSY is broken in the messenger sector and its effects are mediated to the ordinary sector through the standard model (SM) gauge interaction by integrating out the heavy messenger fields (for a review see Ref. [11]). For example, the sfermions $\tilde{f}$ acquire masses of the order of

$$m_{\tilde{f}}^2 \sim \left( \frac{\alpha_{SM}}{4\pi} \right)^2 \Lambda^2,$$

where $\alpha_{SM}$ denotes the appropriate coupling of the SM gauge interaction, and $\Lambda$ is a ratio between $\langle F_S \rangle$ and $\langle S \rangle$ and should be $\Lambda = \langle F_S \rangle / \langle S \rangle \gtrsim 10$ TeV from the present experimental limits on the masses of the superparticles.

The flat direction $X$ obtains a potential through the gauge mediation mechanism. For a large $X (|X| \gg \langle S \rangle)$, the potential is expressed as (e.g. see Ref.[5])

$$V(X) \simeq \left( \frac{\alpha_{SM}}{4\pi} \right)^2 \langle F_S \rangle^2 \left( \ln \frac{|X|^2}{\langle S \rangle^2} \right)^2 + \cdots.$$  (2)

Here we have assumed the positive sign of overall factor. However, for an extremely large $X$, we cannot neglect the contribution from the supergravity. The SUSY breaking effects communicated by the gravity are expected to give the soft mass to the flat direction $X$ which is of the order of the gravitino mass $m_{3/2}$. Then the potential of $X$ due to the gravity is expressed as

$$V(X) \simeq m_{3/2}^2 |X|^2 + \cdots.$$  (3)

Comparing Eq.(2) with Eq.(3), the contribution from the gravity-mediation effect becomes more important for the region $|X| \gtrsim X_{eq}$, where $X_{eq}$ is estimated as

$$X_{eq} \simeq \frac{\langle F_S \rangle}{m_{3/2}} \simeq 10^{13} \text{ GeV} \left( \frac{\langle S \rangle}{10^5 \text{GeV}} \right) \left( \frac{\Lambda}{10^4 \text{GeV}} \right) \left( \frac{m_{3/2}}{100 \text{ keV}} \right)^{-1}.$$  (4)

Here note that in Eq.(2) the logarithmic factor and the gauge coupling constants are almost cancelled. In order that $X_{eq}$ be smaller than about the Planck scale $M_G = 2.4 \times 10^{18}$ GeV, $\langle S \rangle$ should be constrained as

$$\langle S \rangle \lesssim \frac{M_G m_{3/2}}{\Lambda} \simeq 2.4 \times 10^{10} \text{ GeV} \left( \frac{m_{3/2}}{100 \text{ keV}} \right) \left( \frac{\Lambda}{10^4 \text{ GeV}} \right)^{-1}.$$  (5)

\footnote{In the following analysis, we will drop off such an order one constant to make a more conservative analysis.}

\footnote{This upper bound on $\langle S \rangle$ ensures that the $F$-vev in the dynamical SUSY breaking sector $\langle F_{DSB} \rangle$ is equal to or larger than $\langle F_S \rangle$, i.e., $\langle F_{DSB} \rangle \sim m_{3/2} M_G \gtrsim \langle F_S \rangle \sim \langle S \rangle \Lambda$.}
Then we turn to discuss the cosmological evolution of the flat direction $X$ in the inflationary universe. The initial value of $X$ when the primordial inflation ends ($|X| = X_0$) is crucial in this discussion. Our claim here is that the $X_0$ is much larger than $\langle S \rangle$ because of the chaotic condition of the early universe.\textsuperscript{4}

First, we consider the case that $X_0 \gtrsim X_{eq} (\gg \langle S \rangle)$. Since the effective mass of $X$ after the primordial inflation is of the order of the gravitino mass, $X$ starts to roll down along the potential (3) when the Hubble parameter, $H$, of the universe becomes nearly equal to $m_{3/2}$. At this time, the cosmic temperature of the universe $T_{3/2}$ is estimated as

$$T_{3/2} \simeq 1.7 g_*^{-1/4} \sqrt{M_G m_{3/2}} \simeq 7.2 \times 10^6 \text{GeV} \left( \frac{g_*}{200} \right)^{-1} \left( \frac{m_{3/2}}{100 \text{keV}} \right)^{3/2},$$

(6)

where $g_*$ is the effective degree of the relativistic freedom. Here we have assumed that the reheating process of the primordial inflation had been completed at $T > T_{3/2}$. For $T < \sim T_{3/2}$, the flat direction $X$ causes a coherent oscillation with the initial amplitude $X_0$. As the universe expands, the amplitude decreases as $|X| \propto R^{-3/2}$ ($R$: the scale factor) due to the parabolic form of the potential (3).

When the amplitude of the oscillation becomes smaller than $X_{eq}$ [Eq.(4)], the cosmological evolution of $X$ is drastically changed since the logarithmic potential (2) governs the dynamics in this region. The temperature at $|X| \approx X_{eq}$ is given by

$$T_{eq} \simeq T_{3/2} \left( \frac{X_{eq}}{X_0} \right)^{2/3}. \quad (7)$$

The evolution of $X$ is understood as follows; from the virial theorem, we have a relation,

$$2 \langle K \rangle = \left\langle \frac{\partial V}{\partial X} X + \frac{\partial V}{\partial X} X^* \right\rangle, \quad (8)$$

where $K$ is the kinetic energy of $X$ and the bracket represents time average over a cycle. This leads to $\langle K \rangle = \left\langle 2V/\ln(|X|^2/\langle S \rangle^2) \right\rangle$ for $|X| \lesssim X_{eq}$ and $\langle K \rangle = \langle V \rangle$ for $|X| \gtrsim X_{eq}$. With the help of the equation of motion of $X$, the total energy density, $\mathcal{E} \equiv K + V$, behaves as $d\mathcal{E}/dt = -6HK$ and the amplitude $X$ decreases as $|X| \propto R^{-3}$ for $|X| \lesssim X_{eq}$ and $|X|^2 \propto R^{-3}$ for $|X| \gtrsim X_{eq}$. Thus, the energy density of a flat direction rapidly dominates the universe under the logarithmic potential.\textsuperscript{5} It should be noted that an accelerating universe does occur in spite of the oscillation of the field.\textsuperscript{10}

To confirm this, we define an averaged adiabatic index, $\gamma$, as

$$\gamma \equiv \frac{2K}{K + V} = \frac{4}{\ln \frac{|X|^2}{\langle S \rangle^2} + 2}, \quad (9)$$

For a successful inflation we need $\gamma < 2/3$, that is, $\ln \frac{|X|^2}{\langle S \rangle^2} > 4$. Thus, till the amplitude of $X$ becomes of the order of $\langle S \rangle$, the inflation due to the logarithmic potential (2) takes place. We call this inflation “oscillating inflation”.

\textsuperscript{4}The additional SUSY breaking effects during the primordial inflation may naturally explain the initial condition $X_0$.

\textsuperscript{5}This situation does not change even if there exists a coherent oscillation of the moduli field.
When the oscillating inflation ends at $|X| \sim \langle S \rangle$, the temperature of the universe is estimated as

$$T_S \simeq T_{eq} \left( \frac{\langle S \rangle}{X_{eq}} \right)^{\frac{4}{3}}$$

$$\simeq T_{3/2} \left( \frac{X_{eq}}{X_0} \right)^{\frac{4}{3}} \left( \frac{\langle S \rangle}{X_{eq}} \right)^{\frac{4}{3}}$$

$$\simeq 3.9 \text{ GeV} \left( \frac{g_*}{200} \right)^{\frac{4}{3}} \left( \frac{m_{3/2}}{100 \text{ keV}} \right)^{\frac{4}{3}} \left( \frac{\langle S \rangle}{10^6 \text{ GeV}} \right)^{\frac{4}{3}} \left( \frac{\Lambda}{10^4 \text{ GeV}} \right)^{\frac{4}{3}} \left( \frac{X_0}{M_G} \right)^{-\frac{4}{3}}. \quad (10)$$

After the oscillating inflation it is expected that the flat direction $X$ causes an oscillation around its true minimum $\langle X \rangle \lesssim \langle S \rangle$ and the $X$ decay occurs when the Hubble parameter becomes of the order of its decay width. At this epoch, the vacuum energy $V_0$ of the oscillating inflation is transferred into the thermal bath of the universe and reheats the universe to $T = T_{RX}$. Here we do not specify the explicit dynamics during this reheating epoch, but simply estimate the entropy production rate considering the reheating temperature, $T_{RX}$, as a free parameter.\(^6\)

If the vacuum energy $V_0$ is completely transferred into the thermal bath, the oscillating inflation can increase the entropy by a factor $\Delta$:

$$\Delta \simeq \frac{4}{3} \frac{V_0}{2^{\frac{2}{27}} g_* T_{3/2}^3 T_{RX}}$$

$$\simeq 0.15 \frac{\Lambda M_G^2}{m_{3/2}^2 T_{RX}} \left( \frac{g_*}{200} \right)^{-\frac{4}{3}} \left( \frac{X_0}{M_G} \right)^2$$

$$\simeq 2.4 \times 10^{16} \left( \frac{g_*}{200} \right)^{-\frac{4}{3}} \left( \frac{\Lambda}{10^4 \text{ GeV}} \right) \left( \frac{m_{3/2}}{100 \text{ keV}} \right)^{-\frac{4}{3}} \left( \frac{T_{RX}}{10 \text{ MeV}} \right)^{-1} \left( \frac{X_0}{M_G} \right)^2. \quad (11)$$

Here we have used the fact that $V_0 \sim \langle F_S \rangle^2$.\(^7\) From Eq. (11) we find that the maximum entropy production rate is achieved when we take the lowest value of the reheating temperature and the largest value of the initial displacement of $X$ ($X_0 \simeq M_G$). To keep the success of the big-bang nucleosynthesis, we require that the reheating temperature should be larger than $10 \text{MeV}$. Therefore, the oscillating inflation can dilute unwanted particles which survive until late times of the thermal history of the universe.

To end this section, we briefly mention about the initial amplitude $X_0$. We have assumed that $X_0$ might take an arbitrary value ($X_0 \gg \langle S \rangle$) and considered the case $X_0 \lesssim X_{eq}$. As shown in Eq. (11) we find that the maximum entropy production is achieved when $X_0$ takes its maximum value of the order of the Planck scale. Thus we obtain less dilution factor for the case $X_0 \lesssim X_{eq}$. In the following we assume $X_0 \simeq M_G$.

### 3 Cosmological Moduli Problem with Oscillating Inflation

In this section, we assume the dilution mechanism of the oscillating inflation and examine whether it could solve the cosmological moduli problem or not.

\(^6\)In fact, one example of such a low reheating process after the oscillating inflation is explained in Ref. [9].

\(^7\)Here we have dropped off the gauge coupling in Eq. (2) to make a conservative analysis.
The moduli $\phi$ starts to oscillate with the initial amplitude $\phi_0 \sim M_G$ when its mass $m_\phi$ becomes comparable to the Hubble parameter at $T = T_\phi$. Since the moduli mass is $m_\phi \simeq m_{3/2}$, $T_\phi$ is almost same as $T_{3/2}$ [Eq. (1)]. We call this moduli “big-bang moduli”. At $T = T_\phi$, the ratio of the energy density of this oscillating moduli to the entropy density is given by [4]

$$
\left( \frac{\rho_\phi}{s} \right)_{BB} \simeq \frac{1}{2} \frac{m_\phi^2 \phi_0^2}{2 \pi^2 g_* T_\phi^3} \simeq 9.0 \times 10^5 \text{ GeV} \left( \frac{g_*}{200} \right)^{\frac{1}{2}} \left( \frac{m_\phi}{100 \text{ keV}} \right)^{\frac{1}{2}} \left( \frac{\phi_0}{M_G} \right)^2.
$$

Since $\rho_\phi \propto R^{-3}$, this ratio takes a constant value if no entropy is produced. Here it should be noted that the flat direction $X$ and the moduli $\phi$ starts to oscillate at almost the same time. Therefore the oscillating inflation always occurs after the big-bang moduli oscillation begins and can dilute its abundance. Using Eq. (11), the relic abundance of the big-bang moduli after the oscillating inflation is given by

$$
\left( \frac{\rho_\phi}{s} \right)_{BB,0} = \left( \frac{\rho_\phi}{s} \right)_{BB} \times \frac{1}{\Delta} 
\simeq 0.38 \frac{m_{3/2} T_{RX}}{\Lambda} \left( \frac{m_\phi}{m_{3/2}} \right)^{\frac{1}{2}} \left( \frac{X_0}{M_G} \right)^{-2} \left( \frac{\phi_0}{M_G} \right)^2 
\simeq 3.4 \times 10^{-11} \text{ GeV} \left( \frac{m_{3/2}}{100 \text{ keV}} \right) \left( \frac{\Lambda}{10^4 \text{ GeV}} \right)^{-1} \left( \frac{T_{RX}}{10 \text{ MeV}} \right) \times \left( \frac{m_\phi}{m_{3/2}} \right)^{\frac{1}{2}} \left( \frac{X_0}{M_G} \right)^{-2} \left( \frac{\phi_0}{M_G} \right)^2.
$$

You should notice that this abundance is independent of $\langle S \rangle$.

There is another contribution to the moduli density other than the big-bang moduli. During the oscillating inflation the moduli is displaced from its true minimum due to the additional SUSY breaking effects. This displacement $\delta \phi$ is estimated as [4]

$$
\delta \phi \sim \frac{3H_{OI}^2}{m_\phi^2 + 3H_{OI}^2} \phi_0 \sim \frac{3H_{OI}^2}{m_\phi^2} \phi_0,
$$

where $H_{OI}$ is the Hubble parameter during the oscillating inflation and $H_{OI} \ll m_\phi$. When the oscillating inflation ends at $|X| \sim \langle S \rangle$, the moduli displacement from the true minimum is $\delta \phi_0 \sim \langle V_0 \phi_0 \rangle/(m_\phi^2 M_G^2)$. After this inflationary epoch ($|X| \ll \langle S \rangle$), since the logarithmic potential is not effective, the moduli displacement goes to zero at the rate $\delta \phi \propto R^{-3}$ as the universe expands. On the other hand, the amplitude of the big-bang moduli decreases at the rate $\phi \propto R^{-3/2}$ and can not catch up with the motion of the $\delta \phi$. This causes another coherent oscillation of the moduli [4], which was totally neglected in Ref. [3].

We call this moduli “oscillating inflation moduli”. The ratio of the energy density of this oscillation to the entropy density is estimated as

$$
\left( \frac{\rho_\phi}{s} \right)_{OI} \simeq \frac{1}{2} \frac{m_\phi^2 \delta \phi_0^2}{2 \pi^2 g_* T_S^3},
$$

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8 The occurrence of the secondary oscillation of the moduli can be easily understood from the point of view of the energy conservation.
Then the present abundance of the oscillating inflation moduli is given by

\[
\left( \frac{\rho_{\phi}}{s} \right)_{OI,0} \simeq 0.38 \frac{(S)^2 \Lambda^2 T_{RX}}{m_{3/2}^2 M_G^2} \left( \frac{m_{\phi}}{m_{3/2}} \right)^{-2} \left( \frac{\phi_0}{M_G} \right)^2
\]

\[
\simeq 6.5 \times 10^{-14} \text{ GeV} \left( \frac{m_{3/2}}{100 \text{ keV}} \right)^{-2} \left( \frac{\Lambda}{10^4 \text{ GeV}} \right)^2 \left( \frac{(S)}{10^5 \text{ GeV}} \right)^2 \left( \frac{T_{RX}}{10 \text{ MeV}} \right)
\times \left( \frac{m_{\phi}}{m_{3/2}} \right)^{-2} \left( \frac{\phi_0}{M_G} \right)^2 .
\]

(16)

Comparing with the abundance of the big-bang moduli [Eq.(13)], this abundance does depend on the value of \((S)\) and becomes important for the lighter gravitino. Furthermore, both abundances take minimum values for the lowest reheating temperature. We show in Fig.1 and Fig.2 the present abundances of these two moduli with \(T_{RX} = 10 \text{ MeV}\) in term of the density parameter defined by

\[
\Omega_{\phi} h^2 \equiv \frac{(\rho_{\phi}/s)_0}{(\rho_c/s_0)} ,
\]

(17)

where \((\rho_c/s_0)\) is the ratio of the critical density to the present entropy density:

\[
\frac{\rho_c}{s_0} \simeq 3.6 \times 10^{-9} h^2 \text{ GeV} ,
\]

(18)

with \(h\) the present Hubble constant in units of 100 km/sec/Mpc.

Next we see the cosmological constraints on the moduli abundance. Since in the GMSB models the mass of the gravitino is predicted as \(m_{3/2} (\simeq m_{\phi}) \lesssim 1 \text{ GeV}\), the moduli decays most likely into two photons through the nonrenormalizable interaction suppressed by the gravitational scale. Thus, the typical lifetime is estimated as

\[
\tau_\phi \simeq \frac{64\pi M_G^2}{b^2 m_{\phi}^3} \left( \frac{1 \text{ MeV}}{m_{\phi}} \right)^3 \text{ sec} .
\]

(19)

where \(b\) is a constant of order unity depending on the models of the superstring theory. Thus, the moduli becomes stable within the age of the universe if \(m_{\phi} \lesssim 100 \text{ MeV}\) and it should not overclose the universe \((\Omega_{\phi} \lesssim 1)\). We show this limit in Figs.1 and 2.

Moreover, as pointed out in Ref.[6] the moduli with mass \(m_{\phi} \gtrsim 100 \text{ keV}\) is more stringently constrained from the observation of the present x(\(\gamma\))-ray background spectrum since photons produced by the moduli decay directly contribute to the spectrum. This leads to the upper bound on \(\Omega_{\phi}\) as shown in Figs.1 and 2. (Details are found in Ref.[6, 7, 8].)

From the figures it is seen that the abundance of the big-bang moduli [Eq.(13)] puts the upper bound on \(m_{\phi} \simeq m_{3/2}\). This abundance is inversely proportional to \(\Lambda\) and the upper bound on \(\Lambda \lesssim 100 \text{ TeV}\) from the naturalness leads to \(m_{\phi} \lesssim 1 \text{ MeV}\). On the other hand, the abundance of the oscillating inflation moduli puts the lower bound on \(m_{\phi}\). Since this abundance is proportional to \((S)\) and \(\Lambda\), the minimum values of them \((\langle S \rangle \gtrsim \Lambda \gtrsim 10 \text{ TeV})\) put the lower bound as \(m_{\phi} \gtrsim 100 \text{ eV}\). Therefore the allowed region for the moduli mass is obtained
as $m_\phi \sim 100$ eV–1 MeV. We also find that if we take $\langle S \rangle \gtrsim 10^8$ GeV, no allowed region exists because the oscillating inflation moduli exceeds the cosmological constraints (See Fig.2). Thus the oscillating inflation could give a solution to the cosmological moduli problem if $\langle S \rangle \lesssim 10^8$ GeV, and the gravitino mass lies in the region $m_{3/2} \simeq m_\phi \sim 100$ eV–1 MeV.

4 Conclusion

In this letter we have discussed the cosmological moduli problem in the presence of the oscillating inflation caused by the flat direction in the GMSB models. We have found that there are two types of the coherent oscillation of the moduli, i.e. the big-bang moduli and the oscillating inflation moduli and both coherent oscillations are important. The former (latter) leads to the upper (lower) bound of $m_{3/2}$. It has been shown that the oscillating inflation solves the cosmological moduli problem if $\langle S \rangle \lesssim 10^8$ GeV, and the gravitino mass lies in the region $m_{3/2} \simeq m_\phi \sim 100$ eV–1 MeV.

The allowed mass range of the gravitino comes from the requirement that the present cosmic density of the moduli should not overclose the universe for $m_\phi \lesssim 100$ keV and that the photon flux produced by the moduli decay should not exceed the observed x($\gamma$)-ray backgrounds for $m_\phi \gtrsim 100$ keV. Thus, in other words, the moduli with mass $m_\phi \sim 100$ eV-100 keV can be dark matter of the universe. In particular, if the mass of the moduli is $\sim 100$ keV, the x-ray flux from the moduli dark matter will be detected in experiments with high energy resolution as discussed in Ref. [12].

In the present scenario, the primordial baryon asymmetry is also diluted by the oscillating inflation. Since the reheating temperature is quite low ($\sim 10$ MeV), the electroweak or GUT baryogenesis does not work at all. However, as shown in Ref. [3], the Affleck-Dine baryogenesis [13] can produce sufficient baryon asymmetry in the present model if the moduli mass is small ($m_\phi \lesssim 1$ MeV). We also show the constraint on $\Omega_\phi$ from the baryon asymmetry in Fig.1 and Fig.2.

Acknowledgments

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Figure 1: Abundances of the big-bang (BB) moduli (the solid line) and the oscillating inflation (OI) moduli (the long dashed line) for the case that $\Lambda = 10$ TeV, $\langle S \rangle = 100$ TeV, and $T_{RX} = 10$ MeV. We take $m_\phi = m_{3/2}$. The upper bounds for the moduli abundance $\Omega_{\phi}$ from the present critical density (the dotted line) and the cosmic x(\gamma) ray backgrounds (the dot-dashed line) are also shown. We also show the lower bound from the present baryon-entropy ratio $Y_B = 10^{-11}$ (the thick dotted line).
$\Lambda = 10 \text{ TeV}, \langle S \rangle = 10^8 \text{ GeV}, T_{RX} = 10 \text{ MeV}, X_0 = M_G, \phi_0 = M_G$

Figure 2: Same as Fig.1, except for $\langle S \rangle = 10^8 \text{ GeV}$. In this figure we neglect moduli abundances for $m_{3/2} \lesssim 100 \text{ eV}$, since the condition (5) is broken.