Leptogenesis via $LH_u$ flat direction
with a gauged $U(1)_{B-L}$

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Abstract

We study the supersymmetric leptogenesis via $LH_u$ flat direction with a gauged $U(1)_{B-L}$ symmetry. We find that the resultant baryon asymmetry is enhanced compared with the case without a gauged $U(1)_{B-L}$ symmetry. The baryon asymmetry is proportional to the reheating temperature of inflation, but it is independent of the gravitino mass. If high reheating temperatures of inflation $T_R \sim 10^{10}$ GeV are available, the mass of the lightest neutrino, $m_{\nu_1} \sim 10^{-4}$ eV, is small enough to explain the baryon asymmetry in the present universe. Furthermore, the gravitino mass independence of the produced baryon asymmetry allows us to explain the present baryon asymmetry even in low energy SUSY breaking scenarios such as gauge-mediation models. Our leptogenesis scenario is also very special in the sense that it is completely free from the serious Q-ball formation problem.
I. INTRODUCTION

Our universe is composed of baryon, but not anti-baryon. Accounting for this baryon asymmetry in the present universe is one of the most fundamental and challenging problems in particle physics as well as in cosmology. Various mechanisms have been proposed so far to solve this problem. In particular, there has been a growing interest in leptogenesis \[1\], since there are now strong evidences of neutrino oscillations, equivalently the existence of tiny but nonzero neutrino masses. It is, then, very natural to consider lepton-number violating interactions in nature, which is a crucial ingredient for leptogenesis to work. Especially, once we accept the supersymmetry (SUSY) and the existence of tiny Majorana neutrino masses, there automatically exist in the minimal supersymmetric standard model (MSSM) most of the necessary ingredients for the SUSY leptogenesis via $LH_u$ flat direction, which was originally suggested by Murayama and one of the authors (T.Y.) \[2\] based on the idea of the Affleck-Dine (AD) mechanism \[3\].

Recently, we performed a detailed analysis of this MY leptogenesis via $LH_u$ flat direction in the context of gravity-mediated SUSY breaking scenarios including all of the relevant thermal effects \[6,7\]. We found a very interesting aspect of this leptogenesis, that is, “reheating-temperature independence of cosmological baryon asymmetry” \[7\]. The resultant baryon asymmetry is almost determined by the mass of the lightest neutrino (and the gravitino mass), and hence we can determine the lightest neutrino mass from the baryon asymmetry in the present universe, quite independently of the reheating temperature of inflation. This results in a crucial prediction on the rate of the neutrinoless double-beta ($0$ν$\beta\beta$) decay \[7\], which may be testable in future $0$ν$\beta\beta$ decay experiments such as GENIUS \[8\].

In this article, we study the effects of gauging the $U(1)_{B-L}$ symmetry on the leptogenesis via the $LH_u$ flat direction. ($B$ and $L$ are baryon and lepton number, respectively.) As well known, the $U(1)_{B-L}$ symmetry is the unique global symmetry which can be gauged consistently with the MSSM. If we accept a gauged $U(1)_{B-L}$ symmetry, the existence of three families of the right-handed neutrinos are automatically required from the anomaly cancellation conditions. Furthermore, if we assume that this $U(1)_{B-L}$ symmetry is spontaneously broken at high energy scale, these right-handed neutrinos acquire large Majorana masses of the order of the $B-L$ breaking scale, which naturally explains the tiny neutrino masses suggested from the recent neutrino oscillation experiments via the so-called “seesaw
mechanism” [9].

Below the $B - L$ breaking scale, the scalar potential along the $LH_u$ direction is almost the same as in the case without a gauged $U(1)_{B-L}$ symmetry. However, above the $B - L$ breaking scale, the scalar potential is lifted by the effect of the $U(1)_{B-L}$ D-term in a certain region of parameter space, which provides us a new mechanism [10] to stop the $\phi$ field. ($\phi$ is the superfield parameterizing the $LH_u$ “effectively flat” direction.) We show that if the $\phi$ field is stopped by the D-term potential, the baryon asymmetry is enhanced compared with the “global” case without a gauged $U(1)_{B-L}$ symmetry. This is because the available A-term, which provides a phase rotational motion of the $\phi$ field generating the lepton asymmetry, is much bigger than that in the global case. On the other hand, if the $\phi$ field is stopped by the F-term potential coming from the operator responsible for generating the small neutrino mass, the obtained baryon asymmetry becomes the same as that in the global case [6,7]. (In the following discussion, we will call the former case “D-term stopping case”, while the latter one “F-term stopping case.”)

Surprisingly enough, as we will see in the following discussion, if the amplitude of the $\phi$ field is fixed by the D-term potential during inflation, the baryon asymmetry linearly depends on the reheating temperatures of inflation, but it is completely independent of the gravitino mass. This is totally an opposite situation to the global case, in which the final baryon asymmetry is almost independent of the reheating temperatures of inflation and linearly depends on the gravitino mass.

We find in the D-term stopping case that if high reheating temperatures, $T_R \simeq 10^{10}$ GeV, are available, the lightest neutrino of mass, $m_{\nu_1} \simeq 10^{-4}$ eV, can account for the present baryon asymmetry. Furthermore, the “gravitino-mass independence of the cosmological baryon asymmetry” provides us a great advantage in gauge-mediated SUSY breaking scenarios [11]. Gauge-mediated SUSY breaking is widely considered as one of the most interesting mediation mechanisms of SUSY breaking, since it can naturally explain the suppression of the flavor changing neutral currents (FCNC). However, in those scenarios, the reheating temperature of inflation is strongly restricted from above in order to avoid the overproduction of gravitinos [12]. This makes most of the baryo/leptogenesis scenarios impossible, especially when the gravitino is lighter than $\mathcal{O}(\text{MeV})$, since in this case the reheating temperature $T_R$ is restricted as $T_R \lesssim 10^5$ GeV [13]. It was considered that the AD baryogenesis [5]

\footnote{The electroweak baryogenesis seems very difficult to take place in gauge-mediated SUSY breaking models because it is hard to obtain the required light stop, $m_{\tilde{t}} < m_t$.}

\[3\]
(which is the baryogenesis based on the AD mechanism) is the unique candidate to produce an enough baryon asymmetry in gauge-mediated SUSY breaking scenarios. However, from a recent detailed analysis, it becomes clear that the AD baryogenesis is not successful in gauge-mediated SUSY breaking scenarios because of a serious Q-ball formation \[13\]. In the present work, we will show that the MY leptogenesis can naturally explain, in the D-term stopping case, the present baryon asymmetry even in gauge-mediated SUSY breaking models with the gravitino mass smaller than \(\mathcal{O}(100 \text{ keV})\), if the mass of the lightest neutrino is \(m_{\nu_1} \lesssim 10^{-10} \text{ eV}\). We should stress that such a light gravitino may be directly detected in future collider experiments \[14\]. We claim that our scenario is the first minimal model which explains the present baryon asymmetry in the gauge mediation model with a small gravitino mass \(m_{3/2} = \mathcal{O}(100 \text{ keV})\).

II. THE MODEL

First, let us discuss the potential for the relevant fields. To demonstrate our point, we use a simple superpotential as follows:

\[
W = \mu H_u H_d + h N L H_u + \frac{1}{2} \lambda S N N + \eta X (S \bar{S} - v^2).
\]

Here, \(h, \lambda, \eta\) are coupling constants and we assume \(\eta = \mathcal{O}(1)\). \(X, S, \bar{S}\) are singlets under the MSSM gauge group and \(S, \bar{S}\) carry \(-2, +2\) of \(B - L\) charges, respectively. \(N, L, H_u\) and \(H_d\) are a right-handed Majorana neutrino, an \(SU(2)_L\)-doublet lepton and Higgs fields which couple to up and down type quarks, respectively. \(v\) is the breaking scale of the \(U(1)_{B-L}\) symmetry\[4\]. As we will see later, the “effectively flat” direction relevant for the present baryon asymmetry is the flattest flat direction, and hence we will consider only one family which corresponds to the lightest left-handed neutrino. We adopt the following linear combination of \(L\) and \(H_u\) as the flat direction field \(\phi\) \[2\]:

\[
L = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi \\ 0 \end{pmatrix}, \quad H_u = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \phi \end{pmatrix}.
\]

The Kähler potential must have non-minimal couplings of the \(\phi\) field to the inflaton, otherwise the \(\phi\) field gets a large positive mass term of the order of the Hubble parameter

\[3\]

Here and hereafter, we take the couplings \(h, \lambda, \eta\) and \(\mu, v\) to be real. This can be done by field redefinitions in Eq.\((1)\).
(which we will call a Hubble mass term) and it is driven exponentially towards the origin during inflation, and hence the leptogenesis via $LH_u$ flat direction cannot take place \[3\]. Therefore, we assume that there are general non-minimal couplings of the $\phi$ field to the inflaton in the Kähler potential,

$$\delta K \ni \left( \frac{c_\phi}{M_*} I\phi^\dagger \phi + \text{h.c.} \right) + \frac{b_\phi}{M_*^2} \phi^\dagger \phi I^\dagger I + \ldots.$$ \hfill (3)

Here $|c_\phi|, b_\phi = \mathcal{O}(1)$ are coupling constants, and $M_* \simeq 2.4 \times 10^{18}$ GeV is the reduced Planck scale. $I$ is the inflaton superfield. From these non-minimal couplings to the inflaton, the $\phi$ field gets the following SUSY breaking effects during inflation and at the inflaton-oscillation dominated epoch:

$$\delta V \ni -\sqrt{3}H (c_\phi \phi W_\phi + \text{h.c.}) + 3(1 - b'_\phi)H^2 |\phi|^2 + \ldots ,$$ \hfill (4)

where $H$ is the Hubble parameter, $b'_\phi \equiv b_\phi - |c_\phi|^2$, $W_\phi \equiv \partial W/\partial \phi$, and we use the fact that $|F_I|^2 \simeq 3H^2M_*^2$ in these regimes. In the following discussion, we assume $3(1 - b'_\phi) \simeq -1$ for simplicity. We also assume that there are non-minimal couplings of other fields to the inflaton in the Kähler potential. Then, the full scalar potential relevant to the $\phi$ field is given by

$$V = \frac{1}{2} g^2 \left( -2|S|^2 + 2|\bar{S}|^2 + |N|^2 - |L|^2 \right)^2 + |\mu H_u|^2 + |hLH_u + \lambda SN|^2 + |\mu H_d + hNL|^2 + |hNH_u|^2 + \frac{1}{2} \lambda N^2 + \eta X \bar{S}^2 + |\eta(S\bar{S} - v^2)|^2 + |\eta XS|^2 + 3H^2 \left( \sum \lambda Y^2 |Y|^2 \right) - \sqrt{3}H \left( \sum c_Y Y W_Y + \text{h.c.} \right) + V_{SB},$$ \hfill (5)

where $b'_Y \equiv b_Y - |c_Y|^2$, $c_Y$ are non-minimal couplings of $Y (= L, H_u, H_d, X, S, \bar{S}, N)$ to the inflaton in the Kähler potential as coupling constants in Eq. \[3\], and $g = \mathcal{O}(1)$ is the gauge coupling constant of the $U(1)_{B-L}$. $V_{SB}$ represents soft SUSY breaking terms in the present true vacuum. Here, all fields denote the scalar components of the corresponding superfields.

\[4\]We redefined the phase of $c_\phi$.

\[5\]Here, we omit the potential coming from thermal effects, which will be discussed in Section \[V\] and \[\].
The scalar potential in Eq. (5) is so complicated that it seems very difficult to solve the dynamics of the relevant fields. Fortunately, however, we only need to know the shape in the neighborhood of the bottom of this potential, if the curvature around this bottom is as large as the Hubble parameter during inflation. This is because all the scalar fields, which have masses as large as the Hubble parameter, settle down at the bottom of the potential during inflation and trace this potential minimum throughout the history of the universe. This allows us to eliminate many terms in the potential by using the field relations required for minimizing the potential. Therefore, let us first find out the conditions to minimize the potential in Eq. (5). Though this is also a very hard task because of the complexity of the potential, one can find out the minimum of the potential at least when a condition $|\phi|, H \lesssim v$ is satisfied. In fact, we find the approximate minimum of F-terms for $|\phi|, H \lesssim v$, $\bar{S} \approx v^2/S$, $X \approx 0$, $N \approx -\frac{hLH_u}{\lambda S}$, $H_d \approx 0$. (6)

The curvatures around the first three minimums are of the order of the $B-L$ breaking scale $v$, and that around the last one is of the order of the Hubble parameter. As explained above, we simplify the potential by using the relations in Eq (6). Then, the potential in Eq. (5) is reduced to

$$V \approx \frac{1}{2} g^2 \left( -2|S|^2 + 2 \frac{v^4}{|S|^2} + \frac{h^2}{4 \lambda^2} |\phi|^4 - \frac{1}{2} |\phi|^2 \right)^2 + \frac{h^4}{4 \lambda^2} |\phi|^6 + \frac{h^4}{64 \lambda^2} |\phi|^8$$

$$-H^2 |\phi|^2 + 3(1 - b'_S)H^2|S|^2 + 3(1 - b'_S)H^2 \frac{v^4}{|S|^2} + 3(1 - b'_N)H^2 \frac{h^2}{4 \lambda^2} |\phi|^4$$

$$+ \frac{\sqrt{3}}{2} H \left(c'_\phi \frac{h^2 \phi^4}{\lambda S} + \text{h.c.} \right) + \frac{1}{2} \mu^2 |\phi|^2 + V_{SB}.$$  

(7)

Here, $c'_\phi \equiv c_\phi - 1/4 c_S$. One may wonder why the “flat” direction field $\phi$ can develop a large expectation value in spite of the presence of the $U(1)_{B-L}$ D-term. This is because the $S$ field shifts and absorbs the D-term potential. Here, we have

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6Here, we assume a positive Hubble-order mass term for the $H_d$ field. This is necessary in order to avoid contamination of $H_d$ to the $LH_u$ flat direction. We have checked both analytically and numerically that the contamination of the $H_d$ field becomes nonnegligible only after the $H \lesssim \mu$, and hence the following discussion is not affected much by the $H_d$ contamination.
Then, we get the following potential:

\[ V \simeq V_{SB} + \frac{1}{2} \mu^2 |\phi|^2 + \frac{\sqrt{3}}{2} H \left( c^\prime \frac{h^2}{\lambda} \phi^4 \right) + \frac{h^4}{4 \lambda^2 |S|^2} |\phi|^4 + \mathcal{O}\left( \frac{|\phi|^8}{v^2} \right) \]

\[ -H^2 |\phi|^2 + 3(1 - b'_S) H^2 |S|^2 + 3(1 - b'_S) H^2 \frac{v^4}{|S|^2} + 3(1 - b'_N) H^2 \frac{h^2}{4 \lambda^2 |S|^2} |\phi|^4. \]  \hfill (9)

The above procedure can be justified as long as the $B-L$ breaking scale $v$ is larger than the Hubble parameter during inflation, because the curvature around the minimum in Eq. (8) coming from the D-term potential is also of the order of the scale $v$. We assume that this is the case in the following discussion, which is, however, very plausible since the $B-L$ breaking scale $v$ and the Hubble parameter of inflation $H_I$ are most likely $v \gtrsim 10^{13}$ GeV and $H_I \lesssim 10^{13}$ GeV. As we will see later, the third term,\[ \frac{\sqrt{3}}{2} H \left( c^\prime \frac{h^2}{\lambda} \phi^4 \right), \]plays a crucial role in the D-term stopping case. We will call this term “Hubble A-term” in the following sections.

### III. EVOLUTION OF THE $\phi$ FIELD DURING INFLATION

In this section, we discuss the evolution of the $\phi$ field during the inflation. As we will see later, in the D-term stopping case, the $\phi$ field settles down at the $B-L$ breaking scale $v$ during the inflation. This allows us to use the Hubble A-term as a torque to rotate the $\phi$ field after the inflation ends. It will become clear in Section IV that the torque coming from this Hubble A-term is the crucial ingredient to enhance the baryon asymmetry compared with the global case as well as the F-term stopping case.

First, let us investigate the potential for the $\phi$ field during the inflation in the range $|\phi| \lesssim v$. In this range, the $S$ field in Eq. (8) is expanded into the following form:

\[ |S|^2 \simeq v^2 - \frac{1}{8} |\phi|^2 + \mathcal{O}\left( \frac{|\phi|^4}{v^2} \right), \quad \text{for } |\phi| \lesssim v. \]  \hfill (11)

Then, by substituting Eq. (11) into Eq. (8), we get the effective potential for the $\phi$ field in the range $|\phi| \lesssim v$ in the following form:
\[ V \simeq V_{SB} + \frac{1}{2} \mu^2 |\phi|^2 + \frac{\sqrt{3}}{2} \frac{H}{M} (c_\phi^2 \phi^4 + \text{h.c.}) + \frac{|\phi|^6}{4M^2} \]

\[-H^2 |\phi|^2 + \frac{3}{8} (b'_s - b'_{\bar{s}}) H^2 |\phi|^2 + \ldots , \quad (12)\]

where the ellipsis denotes higher order terms in $|\phi|^2/v^2$. Here, $M \equiv \lambda v/h^2$ corresponds to the same symbol used in Ref. [4,6,7] and the mass of the lightest neutrino is written as $m_{\nu_1} = \langle H_u \rangle^2 / M$.

From Eq. (12), we see that the effective potential is almost the same as in the global case with some changes of coefficients in the potential. However, there is an important difference in the Hubble-induced SUSY breaking mass terms. We should not neglect the contributions coming from the $S$ field. In fact, we need the following condition in order that the $\phi$ field develops a large expectation value during the inflation:

\[ b'_s - b'_{\bar{s}} \lesssim \frac{8}{3} . \quad (13) \]

If this is the case, the $\phi$ field has a negative Hubble mass term at least in the range $|\phi| \lesssim v$. Otherwise, the $\phi$ field is driven toward the origin during the inflation because of the positive Hubble mass term, and the MY leptogenesis does not take place.

In the following discussion, we assume that Eq. (13) is satisfied and discuss the scale where the flat direction is lifted. If the balance point between the negative Hubble mass term and the F-term potential $|\phi|^6/(4M^2)$ in Eq. (12) is below the $B-L$ breaking scale $v$, the $\phi$ field is stopped at this balance point,

\[ |\phi| \simeq \sqrt{M H_1} < v , \quad (14) \]

which corresponds to the F-term stopping case. On the other hand, if the condition

\[ \sqrt{M H_1} \gtrsim v , \quad (15) \]

is satisfied, the $\phi$ field can develop its expectation value as large as $v$.

Now, let us explain what happens when the expectation value of the $\phi$ field grows as large as the $B-L$ breaking scale $v$. At this scale (i.e. $|\phi| \simeq v$), the expansion of the $S$ field given in Eq. (11) becomes invalid and above this scale we must use another expansion of the $S$ field as follows:

\[ |S|^2 \simeq 4 \frac{v^4}{|\phi|^2} \phi^2 + \frac{h^2}{2\lambda^2} |\phi|^2 + \left[ \mathcal{O} \left( \frac{v^4}{|\phi|^2} \right) + \mathcal{O} \left( \frac{h^2}{\lambda^2} \right) \right]^2 |\phi|^2 \quad \text{for} \quad |\phi| \gtrsim v . \quad (16) \]
One might wonder whether the above expansion is reliable, since Eq. (8) is based on Eq. (6) that may not be applicable for $|\phi| \gg v$. Here, we first derive the conditions to fix the $\phi$ field at the scale $v$ during the inflation, assuming this expansion is effectively applicable at least for $|\phi| \sim v$. We will justify later the validity of the obtained conditions by numerical calculations.

By substituting the expansion Eq. (16) into Eq. (9), we obtain the following Hubble-induced mass term:\footnote{Here, we assume $h^2/\lambda^2 \ll 1$. If this is not the case, we must include the Hubble mass term coming from the coupling of the right-handed Majorana neutrino to the inflaton, which is the last term in Eq. (9).}

$$V \simeq \left(-1 + \frac{3}{4}(1 - b_S)\right) H^2|\phi|^2 + \mathcal{O}\left(\frac{v^4}{|\phi|^2}H^2\right) \quad \text{for} \quad |\phi| \gtrsim v.$$

Then, the $\phi$ field gets a positive Hubble mass term for $|\phi| \gtrsim v$ if the following condition is satisfied:

$$b_S' \lesssim -\frac{1}{3}.$$

Therefore, if this is the case, the $\phi$ field cannot develop its expectation value above the scale $v$ and hence it is fixed at the $B - L$ breaking scale $v$ during the inflation.

To summarize, the $\phi$ field, and hence $S$ and $\bar{S}$ fields also, are stopped at the $B - L$ breaking scale $v$ during the inflation, if the following conditions are satisfied:\footnote{The conditions for the region $h^2/\lambda^2 > 1$ are obtained by repeating the same procedure.}

$$b_S' - b_S'_{\text{c}} \lesssim \frac{8}{3}, \quad \sqrt{MH_I} \gtrsim v, \quad b_S^\prime \lesssim -\frac{1}{3}.$$

These are the conditions for the D-term stopping case.

To check the conditions in Eq. (19), (especially that of the last condition $b_S^\prime \lesssim -1/3$,) we have numerically solved the coupled equations of motions for the relevant fields using the full scalar potential in Eq. (5). We show the result in Fig. 1, where the amplitude of the $\phi$ field at the end of the inflation is plotted in $b_S' - b_S'_{\text{c}}$ plane. It is found that the conditions in Eq. (19) well explain the result of this numerical calculation. In fact, the amplitude of the $\phi$ field at the end of the inflation lies in the range $1 \lesssim |\phi|/v \lesssim 3$ in most of the parameter space where the conditions in Eq. (19) are satisfied.
As noted in the first paragraph of this section, we can use the Hubble A-term potential [Eq. (10)] as a torque to rotate the $\phi$ field after the inflation ends in the D-term stopping case. In the remaining part of this section, we discuss this point in detail and compare it with the F-term stopping case.

First, if the conditions in Eq. (19) are satisfied (i.e., the D-term stopping case), the curvature of the potential for the $\phi$ field along the phase direction (which is denoted by the symbol $m_{\text{phase}}^2$) is smaller than the Hubble parameter during the inflation. This can be seen from the following relation:

$$m_{\text{phase}}^2 \simeq \frac{H_I}{M} |\phi|^2 \simeq H_I^2 \left( \frac{|\phi|}{\sqrt{M H_I}} \right)^2 < H_I^2. \quad (20)$$

Then, there is no reason to expect that the $\phi$ field sits down at the bottom of the valley of the A-term potential in Eq. (10) during the inflation. Thus, unless there is an accidental fine tuning on the initial phase of the $\phi$ field, it is generally displaced from the bottom of the valley of this A-term potential when the inflation ends. Therefore, this A-term potential kicks the $\phi$ field along the phase direction when the Hubble parameter becomes comparable with the curvature along the phase direction $m_{\text{phase}}^2$. This is the reason why we can use the Hubble A-term potential as a torque to rotate the $\phi$ field and enhance the baryon asymmetry.

Next, we turn to the F-term stopping case. In this case, the amplitude of the $\phi$ field is fixed at $\sqrt{M H_I}$ during the inflation. [See Eq. (14).] An important point is that the curvature around the valley of the Hubble A-term potential, $m_{\text{phase}}^2$, is of the order of the Hubble parameter with the value of the $\phi$ field, $|\phi| \simeq \sqrt{M H_I}$:

$$m_{\text{phase}}^2 \simeq \frac{H_I}{M} |\phi|^2 \simeq H_I^2. \quad (21)$$

Therefore, the phase of the $\phi$ field settles down at the bottom of the valley of this A-term potential during the inflation, and hence the Hubble A-term cannot supply a torque to rotate the $\phi$ field. In this case, the relevant torque for the $\phi$ field only comes from the ordinary A-term potential proportional to the gravitino mass. In this case, the MY leptogenesis with a gauged $U(1)_{B-L}$ symmetry results in the same conclusions as in the global case [3,4], in which the resultant baryon asymmetry is proportional to the gravitino mass. We stress that

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9 The initial phase of the $\phi$ field is generally displaced from the bottom of the valley of this ordinary A-term potential unless the valley of this A-term potential accidentally coincides with that of the Hubble A-term potential.
this leads to a substantial suppression of the baryon asymmetry in gauge-mediation models with a small gravitino mass.\footnote{10}

Finally, we comment on the case where only the condition $b'_{\bar{S}} < -1/3$ in Eq. (19) is not satisfied. In this case, the $\phi$ field has the negative Hubble-mass term even above the $B - L$ breaking scale. We need the effective potential for the $\phi$ field above the scale $v$ in order to determine the scale at which the $\phi$ field stops during the inflation. However, there are so many coupling constants and terms in Eq. (5) which contribute to the potential with comparable importance and hence it is very hard to get a simple effective potential. In spite of this complexity of the potential, it is expected that the baryon asymmetry results in the range between the following two cases. If the curvature around the valley of the Hubble A-term potential is as large as the Hubble parameter during the inflation, $m^2_{\text{phase}} \approx H_I^2$, the baryon asymmetry approaches what is obtained in the global case \footnote{10}. On the other hand, if $m^2_{\text{phase}} < H_I^2$, the baryon asymmetry is enhanced and probably the same as that in the D-term stopping case with which we are mainly concerned in this paper. In any case, we need detailed numerical calculations using the full scalar potential in Eq. (5) in order to obtain the precise amount of the baryon asymmetry in this case, which is beyond the scope of this work.

IV. BARYON ASYMMETRY

In this section, we calculate the baryon asymmetry. First, let us discuss the D-term stopping case. In this case, as explained in the previous section, the $\phi$ field is fixed at the $B - L$ breaking scale $v$ during the inflation, where the curvature along the phase direction coming from the Hubble A-term is smaller than the Hubble parameter $H_I$. Thus, the phase of the $\phi$ field is generally displaced from the bottom of the valley of the Hubble A-term potential.

After the inflation ends, the universe becomes dominated by the oscillating inflaton, and the scale factor of the expanding universe $R$ increases as $R \propto H^{-2/3}$ like in the matter-dominated universe \footnote{11}. In the following discussion, we assume that the production of the lepton number takes place in this inflaton-oscillation dominated epoch, before the reheating process of the inflation completes. (We will justify this assumption in the next section.)

\footnote{10}{See the discussion in the next section.}
At first, the field value of the $\phi$ remains almost constant for a while. This is because the amplitude of the $\phi$ field which gives the potential minimum is mainly determined by the second line in Eq. (9) and it is given by $|\phi| \simeq v$ which is, of course, independent of the Hubble parameter $H$.

Then, the Hubble A-term potential in Eq. (10) kicks the $\phi$ field along the phase direction when the Hubble parameter becomes comparable with the curvature. The Hubble parameter of this time is given by

$$H_{ph} \simeq \frac{v^2}{M} \Leftrightarrow H_{ph}^2 \simeq m_{\text{phase}}^2 \left( \simeq \frac{H_{ph} v^2}{M} \right).$$

(22)

At this time ($H \simeq H_{ph}$), the phase of the $\phi$ field begins to oscillate around the bottom of the valley of the Hubble A-term potential. On the other hand, the amplitude of the $\phi$ field slowly decreases after this time, according to

$$|\phi| \simeq \sqrt{M H},$$

(23)

which is the balance point between the negative Hubble mass term and the operator $|\phi|^6/(4M^2)$.

Since the $\phi$ field starts its oscillation along the phase direction at $H = H_{ph}$, it has already had a large acceleration along the phase direction before it starts coherent oscillation along the radius direction around the origin when $H = H_{osc}$. This $H_{osc}$ is determined by thermal effects, soft SUSY breaking effects and $\mu$-term in the same way as in Ref. [7]. For example, if the $\phi$ field starts its oscillation due to the soft SUSY breaking mass in the true vacuum $m_\phi$, it is given by $H_{osc} \simeq m_\phi$. The evolution of the amplitude of the $\phi$ field and its phase obtained by numerical calculations are shown in Fig. 2 and 3. In these figures, $t$ is the cosmic time in the matter dominated universe $t = 2/(3H)$. We see from these figures that the evolution of the $\phi$ field is well explained by above arguments.

We are now at the point to estimate the baryon asymmetry. The lepton number density is related to the $\phi$ field as follows:

$$n_L = \frac{1}{2} i \left( \phi^* \dot{\phi} - \dot{\phi}^* \phi \right),$$

(24)

where the overdot denotes a derivative with time. The evolution of the $\phi$ field is described by the following equation of motion:

$$\ddot{\phi} + 3H \dot{\phi} + \frac{\partial V}{\partial \phi^*} = 0.$$

(25)
Then, we obtain, from Eq. (24) and (25), the equation of motion for the lepton number density as

\[ \dot{n}_L + 3Hn_L = 2\sqrt{3} \frac{M}{H} \text{Im} \left( c'_\phi \phi^4 \right) . \]  

(26)

Here, we have used \(|S| \simeq v\) for \(H \lesssim H_{ph}\). This can be easily integrated and we obtain the lepton number at time \(t\) as

\[ \left[ R^3 n_L \right] (t) = \int_0^t dt R^3 2 \sqrt{3} \frac{|c'_\phi|H}{M} |\phi^4| \sin \left( \arg(c'_\phi) + 4\arg(\phi) \right) . \]  

(27)

In the regime \(H_{osc} \lesssim H \lesssim H_{ph}\), the amplitude of the \(\phi\) field decreases as \(|\phi| \simeq \sqrt{MH}\) and the sign of \(\sin\) in Eq. (27) changes with \(1/H\) time scale. (See Fig. 2 and 3.) Therefore, from Eq. (27), one sees that the total lepton number oscillates with \(1/H\) time scale with almost constant amplitude, since \(R^3H|\phi^4| \propto H \propto t^{-1}\) in this regime. As soon as the \(\phi\) field starts its coherent oscillation around the origin at \(H = H_{osc}\), the total lepton number is fixed, since the amplitude of the \(\phi\) field decreases as fast as \(|\phi| \propto H\). The lepton number density at time \(t = 2/(3H)\) is written as

\[ n_L(t) \simeq \frac{4}{\sqrt{3}} |c'_\phi| MH^2 \delta_{eff} , \]  

(28)

where \(\delta_{eff} = \sin \left( \arg(c'_\phi) + 4\arg(\phi) \right)\) is an \(O(1)\) effective CP-phase. Note that the Eq. (28) is correct as long as \(H \lesssim H_{ph}\), and hence the resultant lepton-to-entropy ratio is totally independent of the time when the \(\phi\) field starts oscillations around the origin:

\[ \frac{n_L}{s} \simeq \frac{T_R M}{\sqrt{3}M_*^2} |c'_\phi| \delta_{eff} , \]  

(29)

where \(T_R\) is the reheating temperature of the inflation. This lepton asymmetry is partially converted into the baryon asymmetry \(\bar{n}_L\) due to the “sphaleron” effects [15], since it is

\[ \text{11}\text{Here, we neglect the contribution coming from the ordinary A-term, since its contribution is negligible.} \]

\[ \text{12}\text{We have also confirmed by numerical calculations that the final lepton asymmetry } n_L/s \text{ is independent of the oscillation time } H_{osc}. \text{ See also the result for the F-term stopping case in Eq. (34).} \]
produced before the electroweak phase transition. The present baryon asymmetry is given by \[16\]

\[
\frac{n_B}{s} = \frac{8}{23} \frac{n_L}{s}.
\]  

Thus, after all, the present baryon asymmetry is given by

\[
\frac{n_B}{s} \simeq \frac{8}{23} \frac{T_R M}{\sqrt{\delta}} |c'_\phi| |\delta_{\text{eff}}|
\]

\[
\simeq 1.1 \times 10^{-10} \left( \frac{T_R}{10^8 \text{ GeV}} \right) \left( \frac{10^{-6} \text{ eV}}{m_{\nu_1}} \right) |c'_\phi| |\delta_{\text{eff}}|.
\]  

(31)

Here, we have used the relation \( M = \langle H_u \rangle^2 / m_{\nu_1} \).

As stressed before, one sees that the resultant baryon asymmetry is independent of the gravitino mass \( m_{3/2} \) and the starting time of the oscillation of the \( \phi \) field, \( H_{\text{osc}} \). If the mass of the gravitino is large enough \( m_{3/2} \simeq O(1 \text{ TeV}) \), we can avoid the cosmological gravitino problem even if the reheating temperature is rather high \( T_R \simeq 10^{10} \text{ GeV} \) \[17\]. In such a case, we see from the above equation that the mass of the lightest neutrino \( m_{\nu_1} \simeq 10^{-4} \text{ eV} \) is small enough to explain the present baryon asymmetry. On the other hand, if the mass of the lightest neutrino is as small as \( m_{\nu_1} \simeq 10^{-10} \text{ eV} \), we can generate the required baryon asymmetry for low reheating temperatures as \( T_R \simeq 10^4 \text{ GeV} \). In such a low reheating temperature, we are free from the overproduction of gravitinos even in gauge-mediated SUSY breaking scenarios with the small gravitino mass \( m_{3/2} = O(100 \text{ keV}) \) \[12\].

We show the evolution of the baryon asymmetry \( \frac{n_B}{s} = \frac{8}{23} \frac{n_L}{s} \) obtained by a numerical calculation in Fig. 4 and compare it with the estimated value in Eq. (31). From this figure we see that the asymmetry begins to oscillate with almost constant amplitude at \( H = H_{\text{ph}} \) and it is fixed at the time when the \( \phi \) field starts to oscillate around the origin. We confirm from this figure that the resultant asymmetry is well explained by the arguments described in this section.

Next, we briefly discuss the F-term stopping case, in which both the evolution of the \( \phi \) field and the resultant baryon asymmetry are completely the same as those in the global case \[4\]. In this case, the \( \phi \) field starts to move as soon as the inflation ends, according to

13In the present analysis, we neglect the relative sign between the produced lepton and baryon asymmetries.
\(|\phi| \simeq \sqrt{MMH}\). As mentioned in the previous section, the Hubble-induced A-term cannot play a role to kick the phase of \(\phi\), since the flat direction field \(\phi\) is trapped in the valley of this Hubble-induced A-term in this case. The only available A-term that can cause the motion of \(\phi\) along the phase direction is the ordinary A-term, which is proportional to the gravitino mass \(m_{3/2}\):

\[
V_{SB} \ni \frac{m_{3/2}}{8M} \left( a_m \phi^4 + \text{h.c.} \right),
\]

(32)

where \(a_m\) is an \(O(1)\) constant. Thus, the equation of motion for the lepton number density in Eq. (26) is replaced by

\[
\dot{n}_L + 3Hn_L = \frac{m_{3/2}}{2M} \text{Im} \left( a_m \phi^4 \right).
\]

(33)

By repeating the same procedure as in the D-term stopping case, we obtain the following baryon asymmetry:

\[
\frac{n_B}{s} \bigg|_{\text{F-term}} \simeq \frac{2}{69} \frac{T_R M}{M^2} \left( \frac{m_{3/2}}{H_{osc}} \right) |a_m| \delta_{eff}.
\]

(34)

The factor \(m_{3/2}/H_{osc}\) gives rise to a strong suppression of the resultant baryon asymmetry for high reheating temperatures, since \(H_{osc}\) becomes much larger than the soft mass \(m_\varphi\) in that region \([6,7]\). We also stress here that this factor shows the strong suppression of the baryon asymmetry in the gauge-mediated SUSY breaking scenario, since \(m_{3/2} \ll m_\varphi < H_{osc}\).

**V. OTHER CONSTRAINTS**

In the previous section, we found that the resultant baryon asymmetry is enhanced in the D-term stopping case compared with the F-term stopping case and also the global case. However, in the D-term stopping case, there are some conditions we must check, other than those in Eq. (19). These are the conditions for avoiding the early oscillation of the \(\phi\) field. Although the baryon asymmetry in the D-term stopping case is independent of the \(H_{osc}\) as long as \(H_{osc} < H_{ph}\), if the \(\phi\) field starts its oscillation around the origin before \(H = H_{ph}\),

\[14\] The “reheating temperature independence of the baryon asymmetry” in the global case comes from the fact that \(H_{osc}\) in Eq. (34) is proportional to \(T_R\) or \(T_R^{3/2}\) in a large part of the parameter space, which leads to \(n_B/s \propto T_R^{1/3}\) \([6]\).
i.e. $H_{\text{osc}} > H_{ph}$, the resultant baryon asymmetry is strongly suppressed.\footnote{15} The soft SUSY breaking mass term of the $\phi$ field and the $\mu$-term cause the early oscillation of the $\phi$ field when they satisfy\footnote{16}

$$m_{\phi}, \mu > H_{ph} \simeq \frac{v^2}{M}.$$  \hspace{1cm} (35)

Furthermore, thermal effects may cause the early oscillations. A field which couples to the $\phi$ field gets an effective mass $f_k|\phi|$, where $f_k$ is a coupling constant to the $\phi$ field. If the cosmic temperature $T$ is larger than this effective mass, thermal fluctuations of that field produce the thermal mass term for the $\phi$ field, $c_k f_k^2 T^2 |\phi|^2$. The list of $f_k$, $c_k$ for the LH$_u$ flat direction is given in Ref. \cite{6}. Then, if one of these thermal mass terms exceeds the Hubble parameter in the regime $H > H_{ph}$, it causes the early oscillations of the $\phi$ field, and then suppresses the baryon asymmetry substantially. Namely, the following condition is needed to avoid the early oscillation of the $\phi$ field by the thermal mass terms:

$$H^2 > \sum_{f_k, v<T} c_k f_k^2 T^2 \text{ for } H > H_{ph}.$$  \hspace{1cm} (36)

There is another thermal effect which was pointed out in Ref. \cite{19}. Along the LH$_u$ flat direction, $SU(3)$ gauge symmetry remains unbroken, and hence, gluons and gluinos are massless. Furthermore, the down type (s)quarks also remain massless since they have no coupling to the $\phi$ field. These light fields produce the free energy which depends on the $SU(3)$ gauge coupling constant. We obtain the effective potential $V \propto g_s^2 T^4$ at two loop level \cite{20}, where $g_s$ is the gauge coupling constant of the $SU(3)$. At first sight, there seems no dependence on the $\phi$ field in this free energy. However, the up type (s)quarks get large masses from the couplings to the $\phi$ field $y_u$, and if $y_u |\phi| > T$, they decouple and change the trajectory of the running coupling constant of the $SU(3)$. This effect produces the effective potential for the $\phi$ field.\footnote{17}

\footnote{15} In this case, the resultant baryon asymmetry is inversely proportional to $H_{osc}^2$. 

\footnote{16} Note that, in gauge-mediated SUSY breaking scenarios, the condition for $m_{\phi}$ is absent if the messenger scale is lower than the breaking scale of the $U(1)_{B-L}$. The existence of the $\mu$-term depends on the energy scale of the dynamics which produces the $\mu$-term.

\footnote{17} For non-abelian gauge symmetries, decouplings of matter particles give always positive contri-
\[ \delta V \simeq a_g \alpha_S^2 T^4 \log \left( \frac{|\phi|^2}{T^2} \right), \]

(37)

where \( a_g \) is a constant a bit larger than unity and \( \alpha_S \equiv g_S^2/4\pi \). Then, the following condition is required to avoid the early oscillation caused by this potential:

\[ H^2 > \frac{a_g \alpha_S^2 T^4}{v^2} \text{ for } H > H_{ph}. \]

(38)

From Eq. (36) and (38), we obtain the following condition to avoid the early oscillation by thermal effects:

\[ T_R < \min \left[ \min_k \left\{ \max \left( \frac{f_k v^{3/2}}{c_k^{1/4} M_*^{1/2}}, \frac{v^3}{c_k f_k^2 (M_* M^3)^{1/2}} \right) \right\}, \frac{v^2}{a_g^{1/2} \alpha_s (M_* M)^{1/2}} \right], \]

(39)

where, we have used the fact that the cosmic temperature behaves as \( T = (HT_R^2 M_*)^{1/4} \) before the reheating process of the inflation ends. Taking all of these effects into account, we get the allowed regions which are free from the early oscillations as in Fig. 5. The regions below four solid lines are free from the early oscillations. These four lines correspond to the breaking scale of the \( U(1)_{B-L} \), \( v = 10^{16}, 10^{15}, 10^{14}, 10^{13} \) GeV from left to right, respectively. The vertical parts of these lines come from the condition in Eq. (35). Here, we assume the existence of \( m_\phi \) or \( \mu \)-term when the amplitude of the \( \phi \) field is of the order of the \( B-L \) breaking scale. The shaded region denotes the present baryon asymmetry, \( n_B/s \simeq (0.4 - 1) \times 10^{-10} \). From the Fig. 5, we see that the early oscillations can be easily avoided, especially when the \( B-L \) breaking scale satisfies \( v \gtrsim 10^{14} \) GeV.

Finally, we here check the assumption that the production of the lepton asymmetry takes place during the inflaton-oscillation dominated era, \( i.e., \) before the reheating process of the inflation completes. Then, we have the following constraint:

\[ H_{osc} > \Gamma_I \simeq \left( \frac{\pi^2 g_*}{90} \right)^{1/2} \frac{T_R^2}{M_*}, \]

(40)

\[ \text{butions, and the net contribution from decouplings of gauge bosons and gauginos is always negative.} \]

\[ \text{The } LH_u \text{ flat direction has only positive contributions for } SU(3) \text{ gauge symmetry. On the other hand, other flat directions such as } \bar{u} \bar{d} \bar{d}, LL\bar{e}, \ldots \text{ receive large negative contributions and probably have negative thermal-log term.} \]

\[ \text{In the regions above the four solid lines, the estimation of } H_{osc} \text{ is different from that in Ref. [1].} \]
where $\Gamma_I$ is the decay rate of the inflaton $I$ and $g_*$ denotes the number of relativistic degrees of freedom, which is $g_*(T) \approx 200$ in the MSSM for $T \gg 1$ TeV. As long as there exists a soft mass $m_\phi$ or the $\mu$-term at high energy scale, $H_{osc}$ is at least larger than that. Thus, the above constraint is indeed satisfied for $T_R \lesssim 10^{10}$ GeV $\times (m_\phi$ or $\mu/1$ TeV)$^{1/2}$. Furthermore, the Hubble parameter of the oscillation time $H_{osc}$ becomes much larger for higher reheating temperatures, because of the thermal effects. For example, the $H_{osc}$ always satisfies the following relation \[7\]:

$$H_{osc} \geq \alpha_S T_R \left( a_g M_\ast \right)^{1/2}.$$  

Thus, the constraint given in Eq. \[40\] is satisfied as long as

$$T_R \lesssim 10^{15} \text{ GeV} \times \left( m_{\nu_1}/10^{-8} \text{ eV} \right)^{1/2},$$

which is the case in all the relevant parameter space in the present analysis. (See Fig. \[5\].) Therefore, the assumption $\Gamma_I < H_{osc}$ is justified.

**VI. Q-BALL PROBLEM**

One of the physical reasons to make the leptogenesis via $LH_u$ flat direction so special among other baryo/leptogenesis scenarios using the AD mechanism is that it is free from the Q(L)-ball \[21\] problem. The coherent oscillation of the $\phi$ field is unstable with spatial perturbations if the potential of the $\phi$ field is flatter than the quadratic potential. If this is the case, the coherent oscillation of the $\phi$ field fragments into non-topological solitons, Q-balls \[22\]. By the recent detailed analysis by lattice simulations, it becomes clear that almost all of the charges carried by the $\phi$ field are absorbed into these Q-balls \[23\], and the present baryon asymmetry must be provided by the decay of the Q-balls, not by the direct decay of the $\phi$ field. The formation of Q-balls is a generic feature of the AD baryo/leptogenesis regardless of the mediation mechanism of SUSY breaking.

In gravity-mediated SUSY breaking scenarios, the potential of the $\phi$ field is slightly flatter than the quadratic potential due to the running of the soft mass of the $\phi$ field coming from

**19** In this section, we will use the same symbol $\phi$ to denote the field which parameterizes a general flat direction.
gaugino loops [24,25]. In this case, the Q-ball is generally unstable. A difficulty arises from its long life time. If the decay temperature of the Q-balls is well below the freeze out temperature of the lightest supersymmetric particle (LSP), the amount of the LSP cold dark matter can be written as

\[ \Omega \chi = 3 \left( \frac{N_\chi}{3} \right) f_B \left( \frac{m_\chi}{m_n} \right) \Omega_B, \]

where \( N_\chi \) is the number of LSP’s produced per baryon number, which is at least 3, and \( f_B \approx 1 \) is the fraction of baryon number stored in the form of Q-balls, and \( m_n \) and \( m_\chi \) are the nucleon mass and the neutralino (\( \chi \)) LSP mass, respectively. \( \Omega_X \) denotes the ratio of the energy density of \( X \) to the critical density of the present universe. From this relation, it is clear that the late time Q-ball decay leads to the over production of the LSP cold dark matter. In fact, this late time decay of Q-balls invalidates most of the AD baryo/leptogenesis scenarios in gravity-mediated SUSY breaking models.

On the other hand, in gauge-mediated SUSY breaking scenarios, the soft SUSY breaking mass of the \( \phi \) field vanishes above the messenger scale, and the \( \phi \) field has only logarithmic potential. This fact leads to the formation of Q-balls. The Q-balls in gauge-mediated SUSY breaking scenarios have very different characters compared with those in gravity-mediated SUSY breaking models. The effective mass of the Q-ball per baryon number is proportional to \( Q^{-1/4} \) [26], where \( Q \) is the charge of the Q-ball. Therefore, if the charge of the produced Q-ball is large, the Q-balls become completely stable and they remain as a cold dark matter in the universe [22]. Once we fix the flat direction for the AD baryo/leptogenesis, we can estimate the size of the produced Q-ball and its number density by using the result of the recent numerical calculations in Ref. [13]. We find that the produced Q-balls overclose the universe in almost all of the regions of parameter space. Though there remains some tiny regions in parameter space in which the produced Q-balls do not overclose the universe, most part of such regions are already experimentally excluded [13]. If we use the leptonic flat directions, such as \( LL\bar{e} \), the produced Q(L)-balls can decay into neutrinos, and hence they do not overclose the universe. Unfortunately, however, we can only use the leptonic charges which evaporates before the electroweak phase transition, because we must convert

\[ ^{20} \text{The present status of AD baryo/leptogenesis in gravity-mediated SUSY breaking models is summarized in Ref. [11]. There, we also proposed an interesting AD baryo/leptogenesis model with a gauged } U(1)_{B-L} \text{ symmetry, which is an almost unique solution to solve the Q-ball problem in gravity-mediation models except the present leptogenesis via } LH_u \text{ flat direction.} \]
the lepton asymmetry to the baryon asymmetry by “sphaleron effects” [15]. As a result, we find that the present baryon asymmetry can be explained only in extremely small regions of the parameter space.

How about the $LH_u$ flat direction? The soft SUSY breaking mass of the $\phi$ field has no contribution from gluino loops, but, on the other hand, it has a big opposite contribution from large top Yukawa coupling, which makes the potential for the $\phi$ field steeper than quadratic potential. Therefore, the gravity-mediation type Q-balls are not formed in the $LH_u$ flat direction [25]. Furthermore, the $LH_u$ flat direction is the unique flat direction available for the AD baryo/leptogenesis which has a SUSY mass term, $\mu$-term. Thus, once we assume the existence of the $\mu$-term even in the high energy scale relevant for the MY leptogenesis, there always exists the quadratic potential for the $\phi$ field and the formation of Q-balls is avoided even in the gauge-mediated SUSY breaking scenarios. (This is possible if the $\mu$-term is generated by an expectation value of a nonrenormalizable operator.)

Therefore, the MY leptogenesis is free from the Q-ball problem not only in gravity-mediated SUSY breaking scenarios but also in gauge-mediated SUSY breaking scenarios.

VII. DISCUSSION AND CONCLUSIONS

In this paper, we have investigated the supersymmetric leptogenesis via $LH_u$ flat direction originally suggested by Murayama and Yanagida (MY) [2] in the presence of a gauged $U(1)_{B-L}$ symmetry. We have shown that if the flat direction field $\phi$ is stopped by the F-term potential, the situation is almost the same as that in the global case where the $U(1)_{B-L}$ is not gauged.

Interestingly, however, we have found that the resultant baryon asymmetry in the D-term stopping case is enhanced compared with the global case. In this case, the $\phi$ field is stopped at the $B - L$ breaking scale, where the curvature along the phase direction coming from the Hubble-induced A-term is smaller than the Hubble parameter. This allows us to use a much

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21 The authors thank S. Kasuya and M. Kawasaki for useful discussion.

22 If the $\phi$ field starts to oscillate around the origin due to the thermal-log term $\delta V = a_g \alpha_s^2 T^4 \log \left( |\phi|^2 / T^2 \right)$, gauge-mediation type Q-balls are formed. Fortunately, however, in such a case, the Q-balls produced by the thermal-log term are small enough to evaporate well above the electroweak phase transition. Therefore, the thermal-log term does not harm the leptogenesis via $LH_u$ flat direction.
bigger torque to rotate the $\phi$ field, since the $\phi$ field is initially displaced from the bottom of the valley of the Hubble A-term potential unless we tune the initial phase of the $\phi$ field. We have shown that the baryon asymmetry is proportional to the reheating temperature of the inflation, but it is completely independent of the size of the gravitino mass. As a result, if high reheating temperatures $T_R \sim 10^{10}$ GeV are available, the mass of the lightest neutrino, $m_{\nu_1} \sim 10^{-4}$ GeV, is sufficiently small to explain the present baryon asymmetry. Furthermore, the gravitino mass independence of the resultant baryon asymmetry gives us a great advantage in gauge-mediated SUSY breaking scenarios. If the mass of the lightest neutrino is as small as $m_{\nu_1} \approx 10^{-10}$ eV, we can explain the present baryon asymmetry even in the case where the gravitino mass is $O(100\text{keV})$, which may be directly tested in future experiments [14]. We stress that our scenario is the first minimal framework which can explain the present baryon asymmetry in gauge-mediated SUSY breaking scenarios with such a small gravitino mass.

In addition, the MY leptogenesis via $LH_u$ flat direction is very special in the sense that it is completely free from the Q-ball problem. In gauge-mediation scenarios, this requires the existence of the $\mu$-term (SUSY-invariant mass for Higgs multiplets) in the high energy scale relevant for the leptogenesis, which may imply the high energy origin of the $\mu$-term in gauge-mediated SUSY breaking scenarios. The generation of the correct size of the $\mu$-term has been one of the most difficult problems in gauge-mediated SUSY breaking models. The presence of the baryon asymmetry in our universe can be regarded as an interesting implication for the origin of the $\mu$-term.

Finally, we briefly comment on the prediction on the rate of the $0\nu\beta\beta$ decay. As in the global case, the present scenario also requires a hierarchical neutrino mass spectrum to generate the present baryon asymmetry even in the D-term stopping case. This results in almost the same prediction on the rate of the $0\nu\beta\beta$ decay as that in the global case [1]. This consistency check of the present leptogenesis scenario will be given in future $0\nu\beta\beta$ decay experiments [8].

ACKNOWLEDGMENTS

The authors are grateful to S. Kasuya and M. Kawasaki for useful discussion. M.F. and K.H. thank the Japan Society for the Promotion of Science for financial support. This work was partially supported by “Priority Area: Supersymmetry and Unified Theory of Elementary Particles (# 707)” (T.Y.).
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FIG. 1. The amplitudes of the $\phi$ field fixed during inflation which are determined by numerical calculations. In this calculation, we have used the full scalar potential in Eq. (5) to follow the evolution of the relevant fields. Here, we have assumed that $H_I/v = 0.1$, $3(1 - b_\phi') = -1$, $h = 10^{-4}$, $g, \lambda, \eta = 1$, and have randomly generated other coupling constants $b_Y, |c_Y|$ in the range $-2.5 \leq b_Y \leq 2.5$ and $0 \leq |c_Y| \leq 2.5$, respectively. Various symbols denote the $|\phi|$ to $v$ ratio at the end of inflation. We see that the $\phi$ field is, in fact, stopped at the $B - L$ breaking scale $v$ if the conditions in Eq. (13) are satisfied.
FIG. 2. The evolution of the $\phi$ field estimated by a numerical calculation. We assume that $M = 3 \times 10^{23}$ GeV (i.e., $m_{\nu_1} = 10^{-10}$ eV), $v = 10^{14}$ GeV, $H_I = 10^{12}$ GeV, $m_{\phi} = 10^3$ GeV. We see that the $\phi$ field starts to move at $H = H_{ph} \simeq 3 \times 10^4$ GeV according to the equation, $|\phi| \simeq \sqrt{MH}$. Here, $t = 2/(3H)$ is the cosmic time. The $\phi$ field starts to oscillate around the origin at $H_{osc} \simeq m_{\phi}$ since we neglect the thermal effects in this calculation.

FIG. 3. The evolution of the phase of the $\phi$ field estimated by a numerical calculation. The parameters used in this figure are the same as in Fig. 2. Here, we also defined that $\phi'$ is real, and hence the valleys coming from the Hubble A-term lie along $(\arg(\phi) = \pi/4, 3\pi/4, 5\pi/4, 7\pi/4)$. We see that the phase of the $\phi$ field starts to oscillate around the bottom of the valley at $H \simeq H_{ph}$.  

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FIG. 4. The evolution of the baryon asymmetry $\frac{n_B}{s} = \frac{8}{23} \frac{n_L}{s}$ estimated by a numerical calculation. Parameters used in this figure are the same as in Fig. 2 and 3. Here, we also take $T_R = 10^4$ GeV. The thick line denotes the estimated value using Eq. (31). We can see that the asymmetry begins to oscillate with constant amplitude at $H \simeq H_{ph}$ and it is fixed when the $\phi$ field starts the oscillations around the origin. The analytical estimation of the baryon asymmetry agrees well with the obtained value by the numerical calculation.
FIG. 5. The plot of the parameter regions which are free from the early oscillation in $m_{\nu_1} - T_R$ plane. The regions below the four solid lines are free from the early oscillation. These four lines correspond to the breaking scale of the $U(1)_{B-L}$, $v = 10^{16}$, $10^{15}$, $10^{14}$, $10^{13}$ (GeV) from left to right, respectively. The shaded region corresponds to the present baryon asymmetry, $n_B/s \simeq (0.4 - 1) \times 10^{-10}$. We have taken $\delta_{\text{eff}}, |c'_\phi| = 1$ in this figure. Here, we assume the existence of $m_\phi$ or the $\mu$-term when the amplitude of the $\phi$ field is of the order of the $U(1)_{B-L}$ breaking scale $v$. From this figure, we see that if we take the breaking scale of the $U(1)_{B-L}$, $v \gtrsim 10^{14}$ GeV, the early oscillation of the $\phi$ field $H_{\text{osc}} > H_{\text{ph}}$ can be avoided in most of the regions of parameter space.