Affleck-Dine Leptogenesis and Low Scale Inflation

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Abstract

We study Affleck-Dine leptogenesis via the $\tilde{L}H_u$ flat direction in supersymmetric theories. We find that the baryon asymmetry is enhanced when the energy scale of the inflation is sufficiently low. Especially, we consider models of low scale inflation in which the Hubble parameter during inflation is comparable to (but slightly larger than) the gravitino mass $m_{3/2} \sim 1$ TeV. The observed cosmological baryon asymmetry is obtained with the lightest neutrino mass $m_{\nu_1} \sim 10^{-4}$ eV, if the reheating process is suddenly terminated after inflation.
The origin of cosmic baryon asymmetry is one of the most fundamental problems in particle physics and cosmology. Although various mechanisms have been proposed to solve it so far, the mechanism proposed by Affleck and Dine [1] is particularly attractive if supersymmetry (SUSY) is the physics beyond the standard model. In the SUSY standard model there appear various flat directions in the vacuum configuration which carry $B$ and/or $L$ charges in the SUSY limit. Their non-trivial evolution in the early universe could generate sufficiently large $B$ and/or $L$ densities to explain the observed baryon asymmetry.

Especially, Affleck-Dine (AD) leptogenesis via the $\tilde{L}H_u$ flat direction [2] ($\tilde{L}$ is a scalar component of the lepton-doublet superfield) has attracted the attention [3]–[9], since there is now convincing evidence of neutrino oscillations and the suggested tiny neutrino masses indicate the lepton-number violation in nature, which is essential to leptogenesis [10]. It is also noteworthy that in leptogenesis scenarios the present baryon asymmetry is closely related to masses and mixings of neutrinos.

Furthermore, the $\tilde{L}H_u$ direction is very special among various flat directions relevant for the AD mechanism, since we can avoid the serious problem associated with $Q$-balls. The spatial instability in the coherent oscillation of the AD field might lead to the formation of $Q$-balls and spoil the simple description of the AD baryogenesis [11]. In the $\tilde{L}H_u$ direction, however, the $Q$-balls are not formed since its potential is steeper than quadratic one (and hence there is no instability) due to the absence of radiative correction from gluino loops and also the large contribution from the top Yukawa coupling [12].

Recently, detailed analyses of the AD leptogenesis were performed in Refs. [6, 8] including the effects of surrounding thermal plasma pointed out in Refs.[3, 5, 7]. It was shown in Ref. [8] that the resultant baryon asymmetry $n_B/s$ (the ratio of the baryon number density $n_B$ to the entropy density $s$ in the present universe) is determined mainly by the mass of the lightest neutrino $m_{\nu_1}$ and is almost independent on the reheating temperature of inflation $T_R$ in a wide region of $T_R \simeq 10^{5} - 10^{12}$ GeV, and hence the observed value $n_B/s \simeq (0.4 - 1) \times 10^{-10}$ [13] predicts $m_{\nu_1} \simeq (0.1 - 3) \times 10^{-9}$ eV. They also pointed out that such a ultralight neutrino can be tested in the future experiments of neutrinoless double beta decay.

In these analyses, parameters of the inflation model are considered as free parameters. Although the dependence on $T_R$ was discussed, the energy scale of inflation was just assumed to be sufficiently high so that the inflation ends well before the AD field starts to oscillate, which is a crucial time since the net lepton asymmetry produced by the $\tilde{L}H_u$ direction is fixed at this time. However, if one considers inflation models which take place at relatively low energy scale, the above assumption might break down. In this letter, therefore, we investigate the AD leptogenesis in the presence of such a low scale inflation and show that the resultant baryon asymmetry is enhanced.

1In this analysis we assume gravity-mediated models of SUSY breaking with a gravitino mass $m_3/2 \sim 1$ TeV.

2Although the AD leptogenesis introducing $U(1)_{B-L}$ gauge symmetry was discussed in Ref. [8], we do not consider this possibility here.
Let us start by explaining the AD leptogenesis via the flat directions $H_u = \tilde{L}_i$. Here we follow the discussion in Refs. [3, 8]. In the minimal SUSY standard model we can incorporate neutrino masses by introducing the effective operators in the superpotential,

$$W = \frac{1}{2M_i} (L_i H_u)(L_i H_u).$$

Through the seesaw mechanism [14] neutrinos obtain masses

$$m_{\nu_i} = \frac{\langle H_u \rangle^2}{M_i} = \sin^2 \beta \times (3 \times 10^{-7}) \text{ eV} \left(\frac{10^{20} \text{ GeV}}{M_i}\right),$$

where $\langle H_u \rangle = \sin \beta \times 174 \text{ GeV}$ and we take $\sin \beta \simeq 1$ since the final result does not change much. Notice that, although we do not specify here the origin of the operators in Eq. (1), the scales $M_i$ in the presence of heavy Majorana neutrinos correspond to roughly their masses divided by squared of the neutrino Yukawa couplings and hence can be larger than the reduced Planck scale $M_* = 2.4 \times 10^{18} \text{ GeV}$. Since the leptogenesis works most effectively for the flat direction of the first family, we suppress the family index $i$ and consider only the flat direction $\phi/\sqrt{2} \equiv H_u = \tilde{L}_1$. The flat direction $\phi$ obtains its potential from SUSY breaking effects and also from the non-renormalizable operator in Eq. (1) as

$$V_0 = m_\phi^2 |\phi|^2 + \frac{m_{3/2}^2}{8M} (a_m \phi^4 + h.c.) + \frac{1}{4M^2} |\phi|^6,$$

where $m_\phi$ denotes the soft SUSY breaking mass and $a_m$ is a coupling of order one. We take $m_\phi \simeq m_{3/2} \simeq 1 \text{ TeV}$ and $|a_m| \simeq 1$.

In the early universe, the potential (3) is modified as follows. During the inflation and also the inflaton-oscillation period after the inflation ends, the energy of the universe is dominated by the inflaton. This non-zero energy induces an additional SUSY breaking which gives corrections to the potential (3) [3]. Although the explicit form of these terms highly depends on details of the Kähler potential and the inflation model [3, 15], we introduce here the additional terms

$$\delta V_{\text{inf}} = -c_H H^2 |\phi|^2 + \frac{H}{8M} (a_H \phi^4 + h.c.),$$

where $H$ denotes the Hubble parameter and $c_H$ and $a_H$ are real and complex constants. This is because certain values of $c_H$ and $a_H$ can explain the desirable initial condition for the AD mechanism.

Furthermore, $\phi$ receives an additional potential from the thermal effects of the surrounding plasma [3, 5, 7]. It should be noted that even in the period of the inflaton oscillation there is a dilute plasma as a result of scatterings with the thermalized decay products of the inflaton, which temperature is given by $T \simeq (T^2 R H M_*)^{1/4}$ [14]. Here we do not explain these thermal effects in detail but only give the induced terms

$$\delta V_{\text{th}} = \sum_{f_k, |\phi| < T} c_k f_k^2 T^2 |\phi|^2 + a_g \alpha_s^2 T^4 \log \left(\frac{|\phi|^2}{T^2}\right),$$

$^3\tan \beta = \langle H_u \rangle/\langle H_d \rangle$, where $H_u$ ($H_d$) are Higgs fields which couple to up (down) type quarks, respectively.
where \( f_k \) correspond to Yukawa or gauge coupling constants of the field \( \psi_k \) which couple to \( \phi \) and \( c_k \) are real positive constants of order one. \( \alpha_s \) is a strong coupling constant and \( a_g \) is a constant which is a bit larger than unity. (Details can be found in Refs. [6, 8].)

The effective total potential \( V_{\text{tot}} \), which is relevant for the following discussion, is given by

\[
V_{\text{tot}} = \left( m^2 - c_H H^2 + \sum_{f_k|\phi|<T} c_k f_k^2 T^2 \right) |\phi|^2 + a_g \alpha_s^2 T^4 \log \left( \frac{|\phi|^2}{T^2} \right) + \frac{m^{3/2}}{8M} (a_m \phi^4 + h.c.) + \frac{H}{8M} (a_H \phi^4 + h.c.) + \frac{|\phi|^6}{4M^2}.
\]

With this potential we can describe the evolution of \( \phi \) by the equation of motion

\[
\ddot{\phi} + 3H \dot{\phi} + \frac{\partial V_{\text{tot}}}{\partial \phi^*} = 0,
\]

where the dot denotes a derivative with time.

During inflation the energy of the universe is dominated by the vacuum energy of the inflaton and there is no thermal plasma. The Hubble parameter takes an almost constant value \( H_{\text{inf}} \). If \( H_{\text{inf}} \) is larger than \( m_\phi \) and also \( c_H \sim |a_H| \sim 1 \), it is found from Eq. (6) that there is an instability of \( \phi \) at origin and \( \phi \) is trapped at one of the minima of the potential

\[
|\phi| \simeq \sqrt{MH_{\text{inf}}},
\]

\[
\arg \phi \simeq -\frac{\arg a_H + (2n + 1)\pi}{4} \quad (n = 0, 1, 2, 3).
\]

This is because the curvature of the potential around the minimum along both radius and phase directions is of the order of \( H_{\text{inf}} \), \( \phi \) moves towards one of the above minima from any given initial value and settles there. This gives the desirable initial condition for the AD mechanism. Hereafter, we assume \( c_H \simeq |a_H| \simeq 1 \) and consider only the inflation models with \( H_{\text{inf}} \gtrsim m_\phi \simeq m^{3/2} \sim 1 \) TeV.

After inflation ends, the energy of the universe is dominated by the coherent oscillation of the inflaton until the reheating process completes. In this period, although there exists a dilute plasma, as long as the potential for \( \phi \) is dominated by \( \delta V_{\text{inf}} \) [4] and also \( |\phi|^6 \) term in Eq. (3), the flat direction \( \phi \) tracks the instantaneous minimum of the potential

\[
|\phi| \simeq \sqrt{MH},
\]

\[
\arg \phi \simeq -\frac{\arg a_H + (2n + 1)\pi}{4}.
\]

Therefore, the amplitude \( |\phi| \) decreases as \( |\phi| \propto H^{1/2} \propto t^{-1/2} \).

As the universe evolves the negative mass term (i.e., \( -H^2|\phi|^2 \)) is eventually exceeded by another term in the full potential (6). At this time the evolution of \( \phi \) is drastically changed and \( \phi \) begins to oscillate and to rotate around the origin \( \phi = 0 \). The Hubble parameter at this
time $H_{\text{osc}}$, which is crucial to estimate the lepton asymmetry produced by $\phi$ (see below), is given by [3] [8]

$$H_{\text{osc}} \simeq \max \left[ m_\phi, H_k, \alpha_s T_R \left( \frac{a_g M_*}{M} \right)^{1/2} \right],$$

(12)

where $H_k$ are

$$H_k \simeq \min \left[ \frac{M_* T_R^2}{f_k^4 M^2}, \left( c_{k, R} T_R^2 M_*^2 \right)^{1/3} \right].$$

(13)

It should be noted that $H_{\text{osc}}$ should be smaller than $H_{\text{inf}}$. We shall assume this nontrivial fact for a while (see, however, the later discussion).

The evolution of $\phi$ for $H < H_{\text{osc}}$ is fixed depending on which term is dominated the total potential [3]. There are three possibilities, i.e., the dominant term is (i) $m_\phi^2 |\phi|^2$ term, (ii) $T^2 |\phi|^2$ term, or (iii) $T^4 \log(|\phi|^2)$ term. In each case, the damping rate of the amplitude is estimated as (i) $|\phi| \propto t^{-1}$, (ii) $|\phi| \propto t^{-7/8}$ [3], or (iii) $|\phi| \propto t^{-\alpha}$ with $\alpha \simeq 1.5$ [8], respectively. Note that the damping rate in all the above cases is faster than the rate before $\phi$ starts to oscillate.

Now, we are at the point to estimate the lepton asymmetry produced in the considering AD leptogenesis. The lepton number density is given by

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

(14)

and its evolution is described by the equation

$$\dot{n}_L + 3H n_L = \frac{m_{3/2}}{2M} \text{Im} \left( a_m \phi^4 \right) + \frac{H}{2M} \text{Im} \left( a_H \phi^4 \right).$$

(15)

Notice that inflation sets $n_L = 0$ initially. The phase of $\phi$ is kicked by the relative phase between $a_m$ and $a_H$ and the rotational motion of $\phi$ generates the lepton asymmetry [3]. It was shown in Refs. [3] [8] that, comparing two source terms of RHS in Eq. (15), the first term (i.e., the original $A$-term in $V_0$ [3]) gives the dominant contribution in generating $n_L$. Thus, by integrating Eq. (15) we obtain the produced lepton asymmetry at time $t > t_{\text{osc}} \sim H_{\text{osc}}^{-1}$ as

$$\left[ R^3 n_L \right] (t) \simeq \int^t dt' R^3 \frac{m_{3/2}}{2M} \text{Im} \left( a_m \phi^4 \right) + \int^t_{t_{\text{osc}}} dt' R^3 \frac{m_{3/2}}{2M} \text{Im} \left( a_m \phi^4 \right) + \int^t_{t_{\text{osc}}} dt' R^3 \frac{m_{3/2}}{2M} \text{Im} \left( a_m \phi^4 \right),$$

(16)

where $R$ is the scale factor of the universe. The second term of RHS in Eq. (16) gives only a small contribution to the total lepton asymmetry, since $\text{Im} (a_m \phi^4)$ changes its sign rapidly due to the $\phi$ oscillation and also the damping rate of $|\phi|^4$ is faster than the rate of $R^{-3}$. On the other hand, the integrand of the first term in Eq. (16) is almost constant ($\propto t^0$) since $|\phi| \propto t^{-1/2}$. Therefore, the resulting lepton asymmetry for $t > t_{\text{osc}}$ is dominated by the contribution at $H \simeq H_{\text{osc}}$ and we have

$$n_L(t) \simeq n_L(t_{\text{osc}}) \times \frac{R(t_{\text{osc}})}{R(t)^3} \simeq \frac{1}{3} \delta_{\text{eff}} |a_m| m_{3/2} M H_{\text{osc}} \times \frac{R(t_{\text{osc}})}{R(t)^3},$$

(17)
where $\delta_{\text{ef}} \simeq \sin(4\text{arg}\phi + \text{arg}a_m)$ denotes an effective CP-violating phase.

The lepton-to-entropy ratio when the reheating process of inflation completes at $T = T_R$ is estimated as

$$\frac{n_L}{s} \simeq \delta_{\text{eff}} \frac{M_{T_R}}{12M_s^2} \left(\frac{m_3/2}{H_{\text{osc}}}\right).$$

(18)

This ratio takes a constant value as long as no dilution exists in the later epoch. The lepton asymmetry, since it is produced well before the electroweak phase transition (at $T \sim 10^2$ GeV), is partially converted into the baryon asymmetry from the chemical equilibrium between lepton and baryon number through the sphaleron effects. The baryon asymmetry in the present universe is

$$\frac{n_B}{s} = |\delta_{\text{eff}}| \frac{2MT_{R}}{69M_s^2} \left(\frac{m_3/2}{H_{\text{osc}}}\right).$$

(19)

Here we neglected the sign of the produced baryon asymmetry. One sees that the AD leptogenesis is most effective for the flat direction of the first family corresponding to the largest
Figure 2: Contour plot of the Hubble parameter $H_{\text{osc}}$ in the $m_{\nu_1} - T_R$ plane. Corresponding values of $H_{\text{osc}}$ are represented in unit of GeV. In the region below the thick line $H_{\text{osc}}$ takes a constant value $H_{\text{osc}} = m_\phi = 1 \text{ TeV}$.

scale $M$ (or the lightest neutrino mass $m_{\nu_1}$). If $H_{\text{osc}} \gg m_\phi \simeq m_{3/2}$ due to the early oscillation by the thermal effects, the baryon asymmetry is suppressed. In Fig. 1 we show the contour plot of $n_B/s$ by the dashed lines. It is found that the present baryon asymmetry is determined almost independently on the reheating temperature for a wide range of $T_R \simeq 10^5 - 10^{12}$ GeV and the observed value $n_B/s \simeq (0.4 - 1) \times 10^{-10}$ suggests an ultralight neutrino with a mass $\sim 10^{-9}$ eV [8].

In deriving Eq. (19) we made two assumptions on parameters of the inflation model. First, we assumed that the reheating temperature is sufficiently low and its process completes after the lepton asymmetry is fixed at $H \simeq H_{\text{osc}}$. In the parameter space shown in Fig. 1 this assumption is justified. Moreover, such a low reheating temperature is preferred to avoid the cosmological gravitino problem (see the discussion below).

The second assumption is that the scale $v$ of the inflation is sufficiently high and the AD field starts to oscillate well after the inflation ends, i.e., $H_{\text{inf}} = v^2/(\sqrt{3}M_*) > H_{\text{osc}}$. We show in Fig. 2 the contour plot of $H_{\text{osc}}$. It is seen that the scale of $H_{\text{osc}}$ becomes higher in the heavier $m_{\nu_1}$ region due to the large thermal effects. Therefore, for a model of low scale inflation with a fixed $H_{\text{inf}}$ this assumption breaks down in some region of parameter space and the final expression of the present baryon asymmetry in Eq. (19) should be modified.

Thus, we discuss the AD leptogenesis for the case $H_{\text{osc}} > H_{\text{inf}}$. As mentioned before, we
consider only the models with $H_{\text{inf}} \gtrsim m_\phi \simeq m_{3/2}$ to explain the initial condition for the AD mechanism. Since the energy of the radiation after the reheating process completes is smaller than the inflaton vacuum energy $v^4$, one has

$$T_R < 0.5 \sqrt{H_{\text{inf}} M_*} = 2 \times 10^{11} \text{ GeV} \left( \frac{H_{\text{inf}}}{10^5 \text{ GeV}} \right)^{1/2}.$$  \hspace{1cm} (20)

Considering the evolution of the dilute plasma in the period of the inflaton oscillation, the maximum temperature $T_{\text{MAX}}$ after the inflation is achieved when $H = H_{\text{MAX}} = 0.6 H_{\text{inf}}$ and is given by [16]

$$T_{\text{MAX}} \simeq \left( T_R^2 H_{\text{inf}} M_* \right)^{1/4}.$$  \hspace{1cm} (21)

The temperature for $H < H_{\text{MAX}}$ is given by $T \simeq (T_R^2 H M_*)^{1/4}$ as long as $T > T_R$. Therefore, the previous discussion can be applied for $H \leq H_{\text{MAX}}$. Neglecting the small time difference between $H_{\text{inf}}$ and $H_{\text{MAX}}$, it is found that the AD field starts to oscillate and hence the produced lepton asymmetry is fixed just after the inflation ends at $H \simeq H_{\text{inf}}$, since $H_{\text{osc}} > H_{\text{inf}}$. The present baryon asymmetry in the considering case is obtained by replacing $H_{\text{osc}}$ by $H_{\text{inf}}$ in Eq. (19) as

$$\frac{n_B}{s} \simeq |\delta_{\text{eff}}| \frac{2 MT_R}{69 M_*^2} \left( \frac{m_{3/2}}{H_{\text{inf}}} \right).$$  \hspace{1cm} (22)

In Fig. we also show the contour plot of $n_B/s$ by the solid lines for the case $H_{\text{inf}} = 10^5 \text{ GeV}$ (i.e., $v = 6 \times 10^{11} \text{ GeV}$). It is seen that the present baryon asymmetry is enhanced by the rate $H_{\text{osc}}/H_{\text{inf}}$ in the region $H_{\text{osc}} > H_{\text{inf}}$. In this case, with relatively high reheating temperatures, the lightest neutrino mass $m_{\nu_1} \sim 10^{-6} \text{ eV} (\gg 10^{-9} \text{ eV})$ is sufficient to explain the observed baryon asymmetry. Notice that such low values of $H_{\text{inf}}$ are available in a class of SUSY inflation models [19, 20].

Finally, we consider the extreme case that the Hubble parameter is comparable to (but slightly larger than) the SUSY breaking mass of the AD field $\phi$ ($H_{\text{inf}} \sim m_\phi \simeq m_{3/2}$). In this case $\phi$ starts to oscillate just after the end of inflation at $H \simeq H_{\text{inf}} \sim m_\phi$ and hence the produced lepton asymmetry is determined independently on details of the additional potential [5] induced by the thermal plasma. Then, the resultant baryon asymmetry is enhanced since there is no suppression coming from the early oscillation by the thermal effects and its expression is given by dropping off the factor $m_{3/2}/H_{\text{osc}}$ in Eq. (19), which is the one obtained without including the thermal effects in the earlier works.

Furthermore, if the reheating process completes just after the end of the inflation, we expect to have a larger baryon asymmetry since the lepton asymmetry produced at $H \simeq H_{\text{inf}} \sim m_\phi$ does not receive the entropy dilution by the inflaton decay. In this sudden reheating case,

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5Here we do not specify the value of the reheating temperature predicted by the inflation models, but take $T_R$ as a free parameter in the region.

6See, for example, Ref. [4].
the inflationary epoch is just followed by the radiation dominated universe and the reheating temperature is given by

\[ T_R = 0.5 \sqrt{H_{\text{inf}} M_*} = 2 \times 10^{10} \text{ GeV} \left( \frac{H_{\text{inf}}}{m_{3/2}} \right)^{1/2} \left( \frac{m_{3/2}}{1 \text{ TeV}} \right)^{1/2}. \]  

(23)

At this time the lepton asymmetry produced by \( \phi \) is fixed as

\[ \frac{n_L}{s} \simeq \frac{1}{4} \delta_{\text{eff}} |a_m| \left( \frac{m_{3/2}}{H_{\text{inf}}} \right)^{1/2} \left( \frac{m_{3/2}}{1 \text{ TeV}} \right)^{1/2} \left( \frac{10^{-4} \text{ eV}}{m_{\nu_1}} \right). \]  

(24)

where we used the fact \( H = 1/(2t) \) in the radiation dominated universe. We obtain, then, the present baryon asymmetry

\[ \frac{n_B}{s} \simeq 3 \times 10^{-11} \delta_{\text{eff}} |a_m| \left( \frac{m_{3/2}}{H_{\text{inf}}} \right)^{1/2} \left( \frac{m_{3/2}}{1 \text{ TeV}} \right)^{1/2} \left( \frac{10^{-4} \text{ eV}}{m_{\nu_1}} \right). \]  

(25)

Therefore, for models of low scale inflation with \( H_{\text{inf}} \sim m_{3/2} \), if the reheating process is suddenly terminated, the lightest neutrino mass of \( m_{\nu_1} \sim 10^{-4} \text{ eV} \) is small enough to account for the observed baryon asymmetry in the present universe.

Before closing this letter, we should mention the constraint on \( T_R \) from the cosmological gravitino problem. Recent analysis [21], using the gravitino density given in Ref. [22], shows that the nucleosynthesis puts the the upper bound on \( T_R \) as \( T_R \lesssim 10^6, 10^9 \) and \( 10^{12} \) GeV for the gravitino mass \( m_{3/2} = 100 \text{ GeV}, 1 \text{ TeV} \) and \( 3 \text{ TeV} \), respectively. Therefore, a gravitino mass of a few TeV is sufficient to have the relatively high reheating temperatures in the present analysis. Further, considerably higher reheating temperatures are acceptable in some cases [24].

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7 The upper bound on \( T_R \) depends on a gluino mass \( m_{\tilde{g}} \). We take \( m_{\tilde{g}} = 1 \text{ TeV} \).

8 The gravitinos produced nonthermally in the preheating epoch [23] are ignored.
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