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Affleck-Dine leptogenesis and its backreaction to inflaton dynamics

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We investigate the backreaction of the Affleck-Dine leptogenesis to inflaton dynamics in the F-term hybrid and chaotic inflation models in supergravity. We determine the lightest neutrino mass in both models so that the predictions of spectral index, tensor-to-scalar ratio, and baryon abundance are consistent with observations.

I. INTRODUCTION

The success of the Big Bang nucleosynthesis theory requires a large amount of baryon asymmetry at least at the temperature of 1 MeV in the early Universe. However, the baryon asymmetry must be washed out due to the primordial inflation, so that there has to be some mechanism to generate the baryon asymmetry after inflation. The Affleck-Dine baryo/leptogenesis is a promising candidate of baryogenesis in supersymmetric (SUSY) theories [1–3]. A $B-L$ charged scalar field with a flat potential, called an Affleck-Dine (AD) field, can obtain a large tachyonic effective mass and have a large vacuum expectation value (VEV) during and after inflation. As the energy of the Universe decreases, the effective mass becomes inefficient and the AD field starts to oscillate coherently around the origin of its potential. At the same time, the phase direction of the AD field is kicked by its A-term potential. Since the $B-L$ number density is proportional to the phase velocity of the AD field, the $B-L$ asymmetry is generated through this dynamics. Finally, the coherent oscillation decays and dissipates into thermal plasma and then the $B-L$ asymmetry is converted to the desired baryon asymmetry through the sphaleron effects [4, 5].

Since the AD field obtains a large VEV during inflation, we should take into account its effect on inflaton dynamics via supergravity effects. In fact, there are many works revealing that a constant term in superpotential and a scalar field with a large VEV may affect inflaton dynamics [6–9]. These effects may rescue the F-term hybrid and chaotic inflation models, which themselves are somewhat inconsistent with the observations of CMB temperature fluctuations. In this letter, we apply their calculation to the scenario of the Affleck-Dine leptogenesis, focusing on the LH_u flat direction in the minimal SUSY standard model sector. We show that the backreaction of the AD field on the inflaton dynamics can rescue the F-term hybrid and chaotic inflation models and the baryon asymmetry can be consistent with the observation at the same time. We predict extremely small mass for the lightest neutrino, which allows us to calculate the effective Majorana mass for the $0\nu\beta\beta$ decay process.

II. AFFLECK-DINE LEPTOGENESIS

A. Potential of the AD field

Let us focus on the dynamics of the LH_u flat direction: $\phi^2 = (L_i H_u)/\sqrt{2}$, where L_i and H_u are left-handed slepton with a family index i and up-type Higgs, respectively. Since the observations of neutrino oscillation imply nonzero masses of neutrinos, we introduce a superpotential of

$$W^{(\mathrm{AD})} = \frac{m_{\nu_i}}{2\langle H_u\rangle^2}\left(L_i H_u\right)^2, \tag{1}$$

$$\equiv \frac{\lambda}{4M_{\mathrm{Pl}}}\phi^4, \tag{2}$$

where $\langle H_u\rangle = \sin\beta \times 174$ GeV and $\tan\beta \equiv \langle H_u\rangle / \langle H_d\rangle$. We take the mass basis where the mass matrix for the neutrinos is diagonal. Here, the flat direction corresponding to the lightest neutrino is most important for the purpose of the Affleck-Dine leptogenesis, so that we identify that direction as the AD field and denote it as ϕ. The lightest left-handed neutrino mass is then given by

$$m_\nu \simeq 5.1\times 10^{-10}\ \mathrm{eV}\left(\frac{\lambda}{8.2\times 10^{-5}}\right). \tag{3}$$

The coupling constant λ is determined to account for the observations of baryon asymmetry and CMB temperature fluctuations.

The relevant potential of the AD field ϕ is written as

$$V_\phi = V_F + V_{\mathrm{soft}} + V_H + V_T, \tag{4}$$

where

$$V_F = \left|\frac{\partial W^{(\mathrm{AD})}}{\partial\phi}\right|^2, \tag{5}$$

is the F-term potential. The potential V_{soft} represents the Higgs μ term and soft SUSY breaking terms in low energy:

$$V_{\mathrm{soft}} = m_\phi^2|\phi|^2 + \left(a m_{3/2}\lambda\frac{\phi^4}{4M_{\mathrm{Pl}}} + \mathrm{c.c.}\right), \tag{6}$$

where m_ϕ ($=\mathcal{O}(1)$ TeV) is the mass of the LH_u flat direction and $m_{3/2}$ is gravitino mass. We expect that the coefficient of A-term a is of order unity.

The potential of V_H is a so-called Hubble-induced mass term, which comes from supergravity effects [3]. In supergravity, the potential of scalar fields is given by

$$V_{\rm SUGRA} = e^{K/M_{\rm Pl}^2}\left[(D_iW)\, K^{i\bar{j}}\, (D_jW)^* - \frac{3}{M_{\rm Pl}^2} |W|^2 \right], \tag{7}$$

where K is a Kähler potential and $D_iW \equiv W_i + K_iW/M_{\rm Pl}^2$. The subscripts represent derivatives with respect to corresponding fields and $K^{i\bar{j}} \equiv (K_{i\bar{j}})^{-1}$. In order to realize the Affleck-Dine leptogenesis, we assume

$$K = |\phi|^2 + |S|^2 + c\frac{1}{M_{\rm Pl}^2} |\phi|^2 |S|^2 , \tag{8}$$

where S is an inflaton superfield and c is an $O(1)$ constant. When we consider F-term inflation models, where the F-term of inflaton satisfies $|F_S|^2 \simeq 3H_{\rm inf}^2 M_{\rm Pl}^2$ with $H_{\rm inf}$ being the Hubble parameter during inflation, we obtain a Hubble-induced mass of the AD field:

$$V_H = c_H H^2 |\phi|^2 \tag{9}$$

$$c_H = 3(1-c). \tag{10}$$

To realize the Affleck-Dine leptogenesis, we assume $c_H < 0$.

After inflation ends, the inflaton gradually decays into radiation and a background plasma develops with a temperature of

$$T = \left(\frac{36 H(t) \Gamma_I M_{\rm Pl}^2}{g_*(T)\pi^2} \right)^{1/4} \quad \text{for} \;\; T \gtrsim T_{\rm RH}, \tag{11}$$

where $g_*(T)$ is the effective number of relativistic degrees of freedom in the thermal plasma. The decay rate of inflaton Γ_I is related with the reheating temperature $T_{\rm RH}$ as

$$T_{\rm RH} \simeq \left(\frac{90}{g_*(T_{\rm RH})\pi^2} \right)^{1/4} \sqrt{\Gamma_I M_{\rm Pl}}. \tag{12}$$

The AD field acquires a thermal potential via 2-loop effect when its VEV is larger than the temperature:

$$V_T = c_T \alpha_s^2 T^4 \log\left(\frac{|\phi|^2}{T^2} \right), \tag{13}$$

where $c_T = 45/32$ and $\alpha_s \equiv g_s^2/4\pi$ is the strong coupling constant [10, 11].

B. Dynamics of the AD field

When we consider F-term inflation models, the Hubble induced mass term of Eq. (9) arises during inflation. Since we assume $c_H < 0$, the AD field stays at the following potential minimum:

$$\langle |\phi| \rangle_{\rm inf} \simeq \left(\sqrt{\frac{|c_H|}{3}} \frac{H_{\rm inf} M_{\rm Pl}}{\lambda} \right)^{1/2}. \tag{14}$$

After inflation ends, its VEV is determined as Eq. (14) with the replacement of $H_{\rm inf} \to H(t)$ during the inflaton oscillation dominated era. Note that the phase direction of the AD field stays at a certain phase due to the Hubble friction effect. When the Hubble parameter decreases to m_ϕ or $(\phi^{-1}V_T')^{1/2}$, the AD field starts to oscillate around the origin of the potential. We denote the Hubble parameter as $H_{\rm osc}$:

$$H_{\rm osc} \simeq {\rm Max}\left[m_\phi,\; \sqrt{\phi_{\rm osc}^{-1} V_T'} \right], \tag{15}$$

where $\phi_{\rm osc}$ is the VEV of the AD field at the beginning of oscillation. At the same time, the AD field starts to rotate in the complex plane because its phase direction is kicked by the A-term of Eq. (6). This is the dynamics that generates $B-L$ asymmetry. The amplitude of the flat direction decreases as time evolves due to the Hubble friction effect and the $B-L$ breaking effect of Eq. (6) becomes irrelevant soon after the beginning of oscillation. Thus, the generated $B-L$ number is conserved soon after the AD field is kicked in the complex plane. We numerically solve the equation of motion and find that the $B-L$ number density at the beginning of oscillation is given by

$$n_{B-L}(t_{\rm osc}) \equiv \epsilon H_{\rm osc} \phi_{\rm osc}^2 \tag{16}$$

$$\epsilon = (0.2 - 1.7) \times a \sin(n\theta_0) \frac{m_{3/2}}{H_{\rm osc}} \tag{17}$$

$$\equiv \tilde{\epsilon} \frac{m_{3/2}}{H_{\rm osc}}, \tag{18}$$

where θ_0 is an initial phase of the AD field. Here, we define the ellipticity parameter ϵ (≤ 1), which represents the efficiency of the Affleck-Dine mechanism.

Finally, the coherently oscillating AD field decays and dissipates into thermal plasma [12] and the sphaleron effect converts the $B-L$ asymmetry to baryon asymmetry [4, 5]. The resulting baryon-to-entropy ratio Y_b is given by

$$Y_b \equiv \frac{n_b}{s} \simeq \frac{8}{23} \frac{\epsilon T_{\rm RH}}{4 H_{osc}} \left(\frac{\phi_{osc}}{M_{\rm Pl}} \right)^2, \tag{19}$$

$$\simeq 6.5 \times 10^{-11} \tilde{\epsilon} \left(\frac{\lambda}{10^{-4}} \right)^{-3/2} \left(\frac{m_{3/2}}{100\ {\rm GeV}} \right), \tag{20}$$

where $8/23$ in the first line is the sphaleron factor [13]. In the second line, we assume $\alpha_s \sqrt{\lambda} T_{\rm RH} \gtrsim m_\phi$ to use $H_{\rm osc} \simeq \sqrt{\phi_{\rm osc}^{-1} V_T'}$ in Eq. (15). Note that the result is independent of the reheating temperature [11]. The observed baryon asymmetry of $Y_b^{\rm obs} \simeq 8.6 \times 10^{-11}$ [14] can be explained when the coupling λ satisfies

$$\lambda \simeq 8.2 \times 10^{-5} \tilde{\epsilon}^{2/3} \left(\frac{m_{3/2}}{100\ {\rm GeV}} \right)^{2/3}. \tag{21}$$

III. BACKREACTION TO INFLATON DYNAMICS

A. F-term hybrid inflation

In this subsection, we consider the simplest model of F-term hybrid inflation [15, 16] taking into account the effect of the AD field on the dynamics of inflaton. The superpotential of the inflaton sector is given by

$$W^{(\text{inf})} = \kappa S \left(\psi \bar{\psi} - \frac{v^2}{2} \right), \tag{22}$$

where κ is a coupling constant, S is inflaton, and ψ and $\bar{\psi}$ are waterfall fields. When the inflaton S has a sufficiently large VEV, the waterfall fields obtain large effective masses of $\kappa \langle S \rangle$ and thus stay at the origin of the potential. The inflaton S obtains the Coleman-Weinberg potential of

$$V_{\text{CW}} \simeq \frac{\kappa^4 v^4}{32\pi^2} \ln \left(\frac{|S|}{S_{\text{cr}}} \right), \tag{23}$$

where $S_{\text{cr}} \equiv v/\sqrt{2}$. The inflaton S slowly rolls down to the origin of the potential until its VEV reaches the critical value of S_{cr}. During this slow roll, the energy density is dominated by the F-term potential energy of $\kappa^2 v^4/4$, so that inflation occurs. When the inflaton S reaches a critical VEV of S_{cr}, the waterfall fields and inflaton start to oscillate about their global minimum and inflation ends. In this simplest model, the spectral index is predicted as $n_s \simeq 1 - 1/\mathcal{N}_* \simeq 0.98$, where $\mathcal{N}_*$ (≈ 55) is the e-folding number at the horizon exit of the CMB scale. This prediction is inconsistent with the observed value more than 2 sigma level: $n_s^{(\text{obs})} = 0.963 \pm 0.008$ [17].

Now we take into account the backreaction of the AD field to the dynamics of the inflaton. In supergravity, the potential of scalar fields is determined by Eq. (7). When we consider the total superpotential $W = W^{(\text{AD})} + W^{(\text{inf})}$, the terms of $W_S K^{\bar{S}\phi} W_{\bar{\phi}} + \text{c.c.} - 3|W|^2$ give a linear potential of inflaton such as [7]

$$V_{\text{BR}} \simeq a' \frac{\kappa v^2}{M_{\text{Pl}}^2} \left\langle W^{(\text{AD})} \right\rangle S + \text{c.c.}, \tag{24}$$

where a' is an $O(1)$ constant determined by higher dimensional Kahler potentials and $\langle W^{(\text{AD})} \rangle$ is determined by Eqs. (2) and (14). Hereafter we assume $a' = 1$.

The effect of the linear term in the F-term hybrid inflation model has been studied in Ref. [7]. They have found that the linear term affect the inflaton dynamics when the slope of the linear term is the same order with that of the Coleman-Weinberg potential. They introduce a parameter to describe the relative importance of the two contributions to the slope:

$$\xi \equiv \frac{2^{9/2}\pi^2}{\kappa^3 \ln 2} \frac{\langle W^{(\text{AD})} \rangle}{v M_{\text{Pl}}^2}, \tag{25}$$

which should be smaller than unity so that the inflaton can rolls towards the critical value. When ξ is of order but below unity, the linear term is efficient for the inflaton dynamics. We define a critical value of coupling constant for the AD field such as

$$\lambda_c \equiv 2.2 \frac{v^3}{\kappa}, \tag{26}$$

where we use $H_{\text{inf}}^2 = \kappa^2 v^4 / 12 M_{\text{Pl}}^2$. When λ is near the critical value, ξ is close to unity and the backreaction of the AD field to inflaton dynamics is efficient. Note that λ should not larger than λ_c so that the inflaton can rolls towards the critical value and inflation can ends.

Since the linear term breaks R-symmetry, under which the inflaton S is charged, we need to investigate the inflaton dynamics in its complex plane as done in Ref. [7].[1] We read their result of Fig. 9, where desired values of $\langle W^{(\text{AD})} \rangle$ can be read from the contours of gravitino masses by the relation of $m_{3/2} M_{\text{Pl}}^2 \leftrightarrow \langle W^{(\text{AD})} \rangle$.[2] The result is shown in Fig. 1, where the spectral index as well as the baryon asymmetry can be consistent with the observed values in the colored region. Here, we assume that the final phase of the inflaton is larger than $\pi/32$ to avoid a fine-tuning of initial condition. Above the red-dashed curves for each case of gravitino mass, we can neglect the effect of a linear term arising from low energy SUSY breaking, which is investigated in the original work of Ref. [7]. If there is only the effect of a linear term arising from low energy SUSY breaking, the spectral index can be consistent with the observation on the red-dashed curve for each case of gravitino mass. Thus we can explain the observation on and above the red-dashed curve for each case of gravitino mass in our model.[3] The right region is excluded by the cosmic string bound [17], while the upper-right region is excluded by the overproduction of gravitinos via inflaton decay [18] and/or thermal production [19, 20].

Since the value of superpotential of the AD field is determined at each point in Fig. 1, we can determine its coupling constant λ. Then we can use Eq. (20) to calculate the baryon abundance. For the case of $m_{3/2} = 100$ GeV, we can explain the baryon abundance by taking $\tilde{\epsilon}$ properly. On the other hand, for the case of $m_{3/2} = 100$ MeV, the baryon asymmetry cannot be produced efficiently below the blue-doted curve even if $\tilde{\epsilon}$ is as large as unity.[4] We predict lightest neutrino mass m_ν as given in the contour plot. Since the coupling constant in the

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[1] A CP-odd component of inflaton is excited via this dynamics, which also provide another scenario of baryogenesis [21].

[2] The dynamics of the phase direction of the AD field can be neglected for the case of $\lambda \ll \kappa$ [22], which is actually satisfied in our case, so that the dynamics of inflaton is basically equivalent to the one in Ref. [7].

[3] We neglect an $\mathcal{O}(1)$ uncertainty arising near the red-dashed curve that comes from the phase difference between two linear terms.

[4] When the coefficient of A-term a in Eq. (6) is much larger than unity, $\tilde{\epsilon}$ can be larger than unity and the bound of the blue curve disappear.

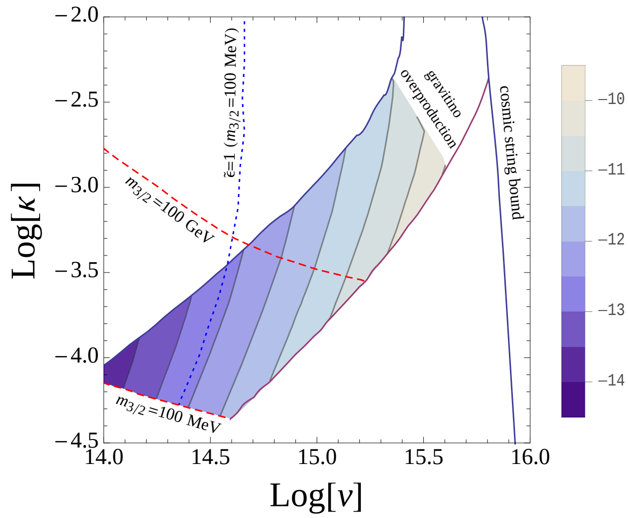

FIG. 1. Contour plot of neutrino mass in the unit of eV in Log[v]-Log[κ] plane. For the case of $m_{3/2} = 100$ GeV, the spectral index as well as the baryon abundance can be consistent with the observations above the corresponding red-dashed curve in the colored region, while for the case of $m_{3/2} = 100$ MeV, they can above the corresponding red-dashed curve and blue-dotted curve in the colored region.

superpotential of the AD field is roughly determined by Eq. (26) to affect the inflaton dynamics, m_ν is larger for larger v and smaller κ. From the figure, we can see that m_ν can be as large as 10^{-10} eV for the case of $m_{3/2} =$ 100 GeV, while it is at most 10^{-13} eV for the case of $m_{3/2} = 100$ MeV.

B. Chaotic inflation

In this subsection, we consider a chaotic inflation model in supergravity where an inflaton superfield I has Z_2 and shift symmetries in the Kahler potential [23]:

$$K^{(\mathrm{inf})} = \frac{1}{2}\left(I + I^*\right)^2. \tag{27}$$

The imaginary part of its scalar component $\eta \equiv (I - I^*)/\sqrt{2}$ is identified with inflaton. The shift symmetry is explicitly broken by a superpotential of

$$W^{(\mathrm{inf})} = mIS, \tag{28}$$

where S is a stabiliser field with a Kahler potential of Eq. (8). When the inflaton has a large VEV, the stabiliser field obtains a large effective mass and stays at the origin. The inflaton potential is then given by the quadratic potential via the F-term of the stabiliser field. Thanks to the shift symmetry in the Kahler potential, the VEV of inflaton can be larger than the Planck scale and quadratic chaotic inflation can be realized in this model.

Here, we take into account the backreaction of the AD field. The full supergravity potential is given by

$$V = e^{|\phi|^2/M_{\mathrm{Pl}}^2}\left[\frac{1}{2}m^2\eta^2\frac{1}{1 + c\,|\phi|^2/M_{\mathrm{Pl}}^2} + \lambda^2\left(\frac{|\phi|^6}{M_{\mathrm{Pl}}^2} + \frac{5}{16}\frac{|\phi|^8}{M_{\mathrm{Pl}}^4} + \frac{1}{16}\frac{|\phi|^{10}}{M_{\mathrm{Pl}}^6}\right)\right], \tag{29}$$

where c is the parameter in the Kahler potential [see Eq. (8)]. This potential implies that the effect of the AD field is relevant when its VEV is as large as the Planck scale. Since $H_{\mathrm{inf}} \sim 10m$ in the chaotic inflation model, the VEV of the AD field is as large as the Planck scale for the case of

$$\lambda \sim \lambda_c \equiv 10\sqrt{c-1}\frac{m}{M_{\mathrm{Pl}}}, \tag{30}$$

[see Eq. (14)].

We numerically solve the equations of motion of the inflaton η and the AD field ϕ and calculate the tensor-to-scalar ratio and spectral index. We show the result in Fig. 2, where we take the parameters c and λ randomly within the intervals of $[1, 10]$ and $[0, 100m/M_{\mathrm{Pl}}]$, respectively. The red, green, and blue dots represent the results at e-folding numbers of 50, 55, and 60, respectively. As a result, the tensor-to-scalar ratio can be as small as 0.14, 0.13, and 0.12 at the e-folding number of 50, 55, and 60, respectively, which is marginally consistent with the present upper bound within 2σ. We plot the results as the light dots for the case of $\lambda/\lambda_c < 0.5$, $5 < \lambda/\lambda_c$, or $c < 5$, which clarifies that the tensor-to-scalar ratio can be smaller only for the case of $0.5 < \lambda/\lambda_c < 5$ and $c > 5$. This requires that the coupling constant in the superpotential is of order $10m/M_{\mathrm{Pl}} \sim 10^{-4}$, so that the lightest neutrino mass is predicted to be of order 10^{-9} eV. Note that the resulting baryon asymmetry of Eq. (20) is naturally consistent with the observation when gravitino mass is of order 100 GeV $-$ 1 TeV.

Finally, we also perform numerical calculations including higher dimensional Kahler potentials of

$$K \supset d\frac{1}{M_{\mathrm{Pl}}^2}|\phi|^4 + d'\frac{1}{M_{\mathrm{Pl}}^4}|\phi|^6 + c'\frac{1}{M_{\mathrm{Pl}}^4}|S|^2|\phi|^4, \tag{31}$$

and find that the tensor-to-scalar ratio can not be smaller than about 0.11 at the e-folding number of 60 even in this case.[5] This is in contrast with the result of Ref. [9], where they have investigated the effect of an additional scalar field to chaotic inflation in a non-SUSY model and found that the tensor-to-scalar ratio can be much smaller than 0.1. This is because the exponential factor in the supergravity potential of Eq. (29) makes the VEV of the AD field smaller and its backreaction to the inflaton dynamics smaller in supergravity.

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[5] We also take into account kinetic couplings between the inflaton and AD field due to the higher-dimensional Kähler potential of $c''|\phi|^2(I + I^*)^2/2M_{\mathrm{Pl}}^2$. However, we find that their effect is also very limited.

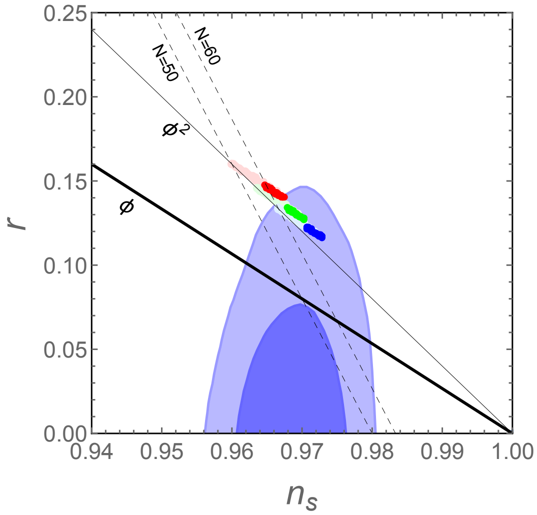

FIG. 2. Spectral index and tensor-to-scalar ratio in the chaotic inflation model with the backreaction of the AD field. The red, green, and blue dots represent our results at e-folding numbers of 50, 55, and 60, respectively. We randomly take 100 points for the parameters c and λ within the intervals of $[1, 10]$ and $[0, 100m/M_{\rm Pl}]$, respectively. We plot the results as the light dots for the case of $\lambda/\lambda_c < 0.5$, $5 < \lambda/\lambda_c$, or $c < 5$. The blue regions are the 1σ (deep colored regions) and 2σ (pale colored regions) constraints of the Planck experiment [24]. For comparison with standard results, we plot the predictions in the chaotic inflation models with linear and quadratic potentials without the backreaction as the black thin and thick lines, respectively, where the results are given as intersection points of black lines and dashed lines for corresponding e-folding numbers.

IV. DISCUSSION AND CONCLUSIONS

We have investigated the backreaction of the AD field to inflaton dynamics, focusing on the LH_u flat direction in the minimal SUSY standard model. In the F-term hybrid inflation model, a linear term of inflaton potential is induced by the nonzero superpotential of the AD field. As a result, the spectral index as well as baryon abundance can be consistent with the observed values. In the chaotic inflation model with a shift symmetry in the inflaton Kahler potential, we have found that the tensor-to-scalar ratio can be as small as 0.12 due to the backreaction of the AD field.

All of the above scenarios require a large VEV of the AD field during inflation. This is also favoured in light of avoiding the baryonic isocurvature constraint, which is particularly important in the chaotic inflation model [25–28]. To realize a large VEV during inflation for the LH_u flat direction, the mass of the lightest neutrino has to be extremely small. Thus the total neutrino mass is given by

$$\sum m_\nu \simeq \begin{cases} 0.06\ \text{eV} & \text{for NH} \\ 0.1\ \text{eV} & \text{for IH,} \end{cases} \tag{32}$$

for the cases of normal hierarchy (NH) and inverted hierarchy (IH), respectively. We can also calculate the upper and lower bounds on the effective Majorana mass for the $0\nu\beta\beta$ decay process such as [11, 29]

$$0.001\ \text{eV} \lesssim |m_{\beta\beta}| \lesssim 0.004\ \text{eV} \quad \text{for NH} \tag{33}$$

$$0.01\ \text{eV} \lesssim |m_{\beta\beta}| \lesssim 0.04\ \text{eV} \quad \text{for IH,} \tag{34}$$

where we take the values for the experimentally measured parameters from Ref. [30]. These results of total neutrino mass and effective Majorana mass are too small to measure in the near future at least for the case of NH. Therefore, if we would measure the total neutrino mass or the effective Majorana mass in the near future, we can falsify our scenario of the AD leptogenesis. On the other hand, if we would experimentally obtain only their upper bound and if the tensor-to-scalar ratio would be measured as $0.12 - 0.15$, our scenario of the AD leptogenesis after the chaotic inflation would be more attractive.

Finally, let us comment on other scenarios of Affleck-Dine baryogenesis using other flat directions, such as the $u^c d^c d^c$ flat direction, where u^c and d^c are up-type and down-type right-handed squarks, respectively. In this case, there are some corrections in calculations of baryon abundance. The most important difference from our scenario is the possibility of the formation of non-topological solitons called Q-balls [31–35]. In particular, as we have shown in this letter, the backreaction of the AD field is relevant when its VEV is sufficiently large during inflation, which implies that large Q-balls may form after the AD baryogenesis. In this case, Q-balls may decay after dark matter (the lightest SUSY particle) freezes out, so that their decay can be a non-thermal source of dark matter. There are interesting scenarios that the non-thermal production of dark matter from Q-ball decay can naturally explain the coincidence between the energy density of baryon and dark matter [36–39] (see, e.g., Ref. [27] in the case of chaotic inflation).

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[1] I. Affleck and M. Dine, Nucl. Phys. B **249**, 361 (1985).
[2] H. Murayama and T. Yanagida, Phys. Lett. B **322**, 349 (1994) [hep-ph/9310297].
[3] M. Dine, L. Randall and S. D. Thomas, Nucl. Phys. B **458**, 291 (1996). [hep-ph/9507453].
[4] V. A. Kuzmin, V. A. Rubakov and M. E. Shaposhnikov, Phys. Lett. B **155**, 36 (1985).
[5] M. Fukugita and T. Yanagida, Phys. Lett. B **174**, 45 (1986).
[6] C. M. Lin and J. McDonald, Phys. Rev. D **74**, 063510 (2006) [hep-ph/0604245].
[7] W. Buchmller, V. Domcke, K. Kamada and K. Schmitz, JCAP **1407**, 054 (2014) [arXiv:1404.1832 [hep-ph]].
[8] W. Buchmuller, C. Wieck and M. W. Winkler, Phys. Lett. B **736**, 237 (2014) [arXiv:1404.2275 [hep-th]]; W. Buchmuller, E. Dudas, L. Heurtier and C. Wieck, JHEP **1409**, 053 (2014) [arXiv:1407.0253 [hep-th]]; W. Buchmuller and K. Ishiwata, Phys. Rev. D **91**, no. 8, 081302 (2015) [arXiv:1412.3764 [hep-ph]]; W. Buchmuller, E. Dudas, L. Heurtier, A. Westphal, C. Wieck and M. W. Winkler, JHEP **1504**, 058 (2015) [arXiv:1501.05812 [hep-th]].
[9] K. Harigaya, M. Ibe, M. Kawasaki and T. T. Yanagida, arXiv:1506.05250 [hep-ph].
[10] A. Anisimov and M. Dine, Nucl. Phys. B **619**, 729 (2001) [hep-ph/0008058].
[11] M. Fujii, K. Hamaguchi and T. Yanagida, Phys. Rev. D **63**, 123513 (2001) [hep-ph/0102187].
[12] K. Mukaida and K. Nakayama, JCAP **1301**, 017 (2013) [JCAP **1301**, 017 (2013)] [arXiv:1208.3399 [hep-ph]].
[13] J. A. Harvey and M. S. Turner, Phys. Rev. D **42**, 3344 (1990).
[14] K. A. Olive *et al.* [Particle Data Group Collaboration], Chin. Phys. C **38**, 090001 (2014).
[15] E. J. Copeland, A. R. Liddle, D. H. Lyth, E. D. Stewart and D. Wands, Phys. Rev. D **49**, 6410 (1994) [astro-ph/9401011].
[16] G. R. Dvali, Q. Shafi and R. K. Schaefer, Phys. Rev. Lett. **73**, 1886 (1994) [hep-ph/9406319].
[17] P. A. R. Ade *et al.* [Planck Collaboration], Astron. Astrophys. **571**, A25 (2014) [arXiv:1303.5085 [astro-ph.CO]].
[18] S. Nakamura and M. Yamaguchi, Phys. Lett. B **638**, 389 (2006) [hep-ph/0602081]; M. Kawasaki, F. Takahashi and T. T. Yanagida, Phys. Lett. B **638**, 8 (2006) [hep-ph/0603265]; T. Asaka, S. Nakamura and M. Yamaguchi, Phys. Rev. D **74**, 023520 (2006) [hep-ph/0604132]; M. Dine, R. Kitano, A. Morisse and Y. Shirman, Phys. Rev. D **73**, 123518 (2006) [hep-ph/0604140]; M. Endo, K. Hamaguchi and F. Takahashi, Phys. Rev. D **74**, 023531 (2006) [hep-ph/0605091]; M. Kawasaki, F. Takahashi and T. T. Yanagida, Phys. Rev. D **74**, 043519 (2006) [hep-ph/0605297]; M. Endo, M. Kawasaki, F. Takahashi and T. T. Yanagida, Phys. Lett. B **642**, 518 (2006) [hep-ph/0607170].
[19] M. Bolz, A. Brandenburg and W. Buchmuller, Nucl. Phys. B **606**, 518 (2001) [Erratum-ibid. B **790**, 336 (2008)] [hep-ph/0012052].
[20] J. Pradler and F. D. Steffen, Phys. Rev. D **75**, 023509 (2007) [hep-ph/0608344].
[21] F. Takahashi and M. Yamada, arXiv:1510.07822 [hep-ph].
[22] M. Yamada, in preparation.
[23] M. Kawasaki, M. Yamaguchi and T. Yanagida, Phys. Rev. Lett. **85**, 3572 (2000) [hep-ph/0004243].
[24] P. A. R. Ade *et al.* [Planck Collaboration], arXiv:1502.01589 [astro-ph.CO].
[25] K. Enqvist and J. McDonald, Phys. Rev. Lett. **83**, 2510 (1999) [hep-ph/9811412]; K. Enqvist and J. McDonald, Phys. Rev. D **62**, 043502 (2000) [hep-ph/9912478]; M. Kawasaki and F. Takahashi, Phys. Lett. B **516**, 388 (2001) [hep-ph/0105134].
[26] S. Kasuya, M. Kawasaki and F. Takahashi, JCAP **0810**, 017 (2008) [arXiv:0805.4245 [hep-ph]].
[27] K. Harigaya, A. Kamada, M. Kawasaki, K. Mukaida and M. Yamada, Phys. Rev. D **90**, no. 4, 043510 (2014) [arXiv:1404.3138 [hep-ph]].
[28] M. Kawasaki and M. Yamada, Phys. Rev. D **91**, no. 8, 083512 (2015) [arXiv:1502.03550 [hep-ph]].
[29] M. Fujii, K. Hamaguchi and T. Yanagida, Phys. Lett. B **538**, 107 (2002) [hep-ph/0203189].
[30] F. Capozzi, G. L. Fogli, E. Lisi, A. Marrone, D. Montanino and A. Palazzo, Phys. Rev. D **89**, 093018 (2014) [arXiv:1312.2878 [hep-ph]].
[31] S. Coleman, Nucl. Phys. **B262** (1985) 263.
[32] A. Kusenko, Phys. Lett. **B405** (1997) 108.
[33] K. Enqvist and J. McDonald, Phys. Lett. B **425**, 309 (1998); Nucl. Phys. B **538**, 321 (1999).
[34] A. Kusenko and M. E. Shaposhnikov, Phys. Lett. B **418**, 46 (1998).
[35] S. Kasuya and M. Kawasaki, Phys. Rev. D **61**, 041301(R) (2000); S. Kasuya and M. Kawasaki, Phys. Rev. D **62**, 023512 (2000); S. Kasuya and M. Kawasaki, Phys. Rev. D **64**, 123515 (2001).
[36] M. Fujii and K. Hamaguchi, Phys. Lett. B **525**, 143 (2002) [hep-ph/0110072]; M. Fujii and K. Hamaguchi, Phys. Rev. D **66**, 083501 (2002) [hep-ph/0205044].
[37] L. Roszkowski and O. Seto, Phys. Rev. Lett. **98**, 161304 (2007) [hep-ph/0608013].
[38] I. M. Shoemaker and A. Kusenko, Phys. Rev. D **80**, 075021 (2009); S. Kasuya and M. Kawasaki, Phys. Rev. D **84**, 123528 (2011) [arXiv:1107.0403 [hep-ph]]; F. Doddato and J. McDonald, JCAP **1307**, 004 (2013) [arXiv:1211.1892 [hep-ph]]; S. Kasuya, M. Kawasaki and M. Yamada, Phys. Lett. B **726**, 1 (2013) [arXiv:1211.4743 [hep-ph]].
[39] A. Kamada, M. Kawasaki and M. Yamada, Phys. Lett. B **719**, 9 (2013) [arXiv:1211.6813 [hep-ph]]; A. Kamada, M. Kawasaki and M. Yamada, Phys. Rev. D **91**, no. 8, 081301 (2015) [arXiv:1405.6577 [hep-ph]].