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Cosmology with a heavy Polonyi field

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Abstract. We consider a cosmologically consistent scenario with a heavy Polonyi field. The Polonyi field with a mass of $\mathcal{O}(100)$ TeV decays before the Big-Bang Nucleosynthesis (BBN) and avoids the severe constraint from the BBN. However, the abundance of the Lightest Supersymmetric Particle (LSP) produced from the decay often exceeds the observed dark matter density. In our scenario, the dark matter density is obtained by the LSP abundance with an aid of entropy production, and baryon asymmetry is generated by the Affleck-Dine mechanism. We show that the observed baryon-to-dark matter ratio of $\mathcal{O}(0.1-1)$ is naturally explained in sequestering models with a QCD axion.

Keywords: cosmology of theories beyond the SM, supersymmetry and cosmology, baryon asymmetry

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1 Introduction

Cosmological observations have revealed the existence of the dark matter (DM) and the baryon asymmetry, while their origins remain unknown for a long time. The cosmic densities of the DM and baryonic components have been precisely measured by the observation of the Cosmic Microwave Background (CMB) radiation, and the observed baryon-to-DM ratio is $\Omega_B/\Omega_{\text{DM}} = 0.18$ [1]. We need a cosmologically consistent scenario explaining both the baryon asymmetry and the DM abundance.

Since those origins cannot be explained in the framework of the Standard Model (SM), there should be underlying physics beyond the SM. Supersymmetry (SUSY)¹ is one of the most interesting models since it achieves the unification of the SM gauge couplings and can relax the hierarchy problem. In addition, supersymmetric models with a conserved R parity predict the stability of the Lightest Supersymmetric Particle (LSP), which becomes a good candidate for DM. Moreover, SUSY extensions of the SM contain a lot of flat directions with $B-L$ charges [4], which can produce $B-L$ asymmetry [5]. In the early universe, one of the flat directions, which we call the Affleck-Dine field, may receive an angular kick from SUSY breaking and R symmetry breaking effects and rotates in its complex plane, which corresponds to the generation of the $B-L$ asymmetry. The $B-L$ asymmetry is converted into the baryon asymmetry through the sphaleron process [6–8]. This mechanism, known as “the Affleck-Dine (AD) mechanism” [5, 9], can produce baryon number more effectively than most baryogenesis scenarios. For these advantages, we focus on supersymmetric theories in this paper.

The SUSY must be spontaneously broken since superparticles have not been discovered yet. The simplest SUSY breaking model is the Polonyi model [10] in which an F -term of an elementary singlet field Z breaks the SUSY in the hidden sector.² This model is attractive because of its simplicity. However, such a singlet field Z , called the Polonyi field, may cause

¹For reviews see refs. [2, 3].

²In the Polonyi model, the SUSY breaking scale can be easily obtained by dynamical transmutation, by assigning a vanishing R charge to the Polonyi field and breaking the R symmetry by gaugino condensation.

a cosmological difficulty [11]. Since the Polonyi potential has no symmetry enhanced points, the minimum of the potential during the inflation is deviated from the true vacuum. After inflation, the Polonyi field begins to oscillate around the true minimum with an amplitude of the order of the Planck scale. The energy density of its coherent oscillation immediately dominates the universe after the inflaton decays. Since the Polonyi field very weakly interacts with particles in the observable sector, its late-time decay upsets the standard scenario of the Big-Bang Nucleosynthesis (BBN), which is referred to as “the Polonyi problem”.

There are mainly two ways of solving the problem. One possible way is to dilute the Polonyi density by some mechanisms, for example, by thermal inflation [12, 13]. In this case, however, pre-existing baryon asymmetry is also diluted, and it is known to be difficult to produce a sufficient amount of baryon number beforehand. Therefore, baryon asymmetry should be produced after the dilution, and we need an intricate model proposed in the context of the thermal inflation [14–16]. The other simple solution is to make the Polonyi field heavy enough to decay before the BBN. Even in this case, we need some dilution since the abundance of LSPs produced from the Polonyi decay exceeds the observed DM abundance unless pair annihilation is very effective [17, 18]. When such dilution is needed, the AD mechanism is the leading candidate for baryogenesis since most baryogenesis scenarios including the thermal leptogenesis [19] cannot produce a sufficient amount of baryon asymmetry beforehand.

In this paper, we consider the case where the Polonyi field is heavy enough to decay before the BBN, i.e., $m_Z \sim \mathcal{O}(100)\text{TeV}$, where m_Z is a mass of the Polonyi field. We construct a cosmologically consistent scenario in the presence of the heavy Polonyi field. In our scenario, the baryon asymmetry is created by dynamics of the AD field which takes the vacuum expectation value (VEV) of the order of the Planck scale in the early universe. Both the Polonyi density and the baryon asymmetry are diluted by late-time entropy production, for example, by the thermal inflation. After the entropy production, the Polonyi field decays into LSPs which explain the observed DM abundance. We show that the baryon-to-DM ratio is simply determined by the LSP mass and the branching fraction of the decay of the Polonyi field into superparticles. The correct baryon-to-DM ratio is obtained when the LSP mass is of $\mathcal{O}(1)\text{TeV}$ and the branching fraction of the decay of the Polonyi into superparticles is of $\mathcal{O}(10^{-3})$.

In order to realize the branching fraction of $\mathcal{O}(10^{-3})$, we consider a specific model satisfying the following conditions. Firstly, we assume that the SUSY breaking sector is sequestered from the visible sector in superspace Lagrangian, which is called “the sequestering model” [20, 21]. The squarks/sleptons and gauginos have vanishing masses at the tree level and acquire loop-suppressed masses through quantum corrections, such as anomaly-mediated SUSY breaking effects [21, 22] (see also refs. [23, 24]) or one loop corrections from Planck-suppressed interactions [21, 25]. Sfermions and gauginos are lighter than the gravitino with a mass of $\mathcal{O}(100)\text{TeV}$. In addition, we introduce a (pseudo-)Nambu-Goldstone boson (NGB). The NGB can be identified with the QCD axion [26–29], which solves the strong CP problem [30–32]. In this model, the Polonyi field mainly decays into NGBs, but their contribution to the DM abundance is negligible since the NGB is much lighter than the LSP (as is the case with the QCD axion) or decays into SM particles. On the other hand, the decay of the Polonyi into superparticles is suppressed for the sequestered potential, and the branching ratio is of $\mathcal{O}(10^{-3})$. The DM abundance is determined by the LSP abundance produced through the Polonyi decay.

This paper is organized as follows. In section 2, we briefly review the Polonyi problem and explain why the entropy production is needed. We explain the AD mechanism in section 3. In section 4, we show that the baryon-to-DM ratio is simply determined by the LSP

mass and the branching fraction. We also introduce the sequestering model as a candidate to realize our scenario. The final section is devoted to summary and discussions.

2 The Polonyi problem

In this section, we explain the cosmological problem of the Polonyi model.

Let us briefly introduce the Polonyi model [10]. In this model, the only ingredient in the SUSY breaking sector is an elementary field Z which is singlet under any symmetry. We call it as the Polonyi field. The superpotential in the hidden sector is given by³

$$W_{\text{hid}} = \mu^2 M_{\text{pl}} \left(1 + c \frac{Z}{M_{\text{pl}}} + \dots \right), \quad (2.1)$$

where μ is a parameter with mass dimension 1, and c is a dimensionless parameter of $\mathcal{O}(1)$. Hereafter, we use M_{pl} as the reduced Planck mass ($M_{\text{pl}} \simeq 2.4 \times 10^{18}$ GeV). Higher order terms are expressed by the ellipsis. Hereafter, we will neglect these terms for simplicity. Note that the parameter μ breaks the R symmetry since Z has an R charge of 0. At the true minimum, the VEV of Z is assumed to be of the order of the Planck scale. The parameter μ is related to the gravitino mass as $|\mu|^2 \simeq \langle |W_{\text{hid}}| \rangle / M_{\text{pl}} \simeq m_{3/2} M_{\text{pl}}$. The F -term of Z is given by $|F_Z| \simeq m_{3/2} M_{\text{pl}}$, which implies spontaneous SUSY breaking. The mass of the Polonyi field is of the order of $m_{3/2}$ for generic Kähler potentials.

The Polonyi model is attractive because of its simplicity and has been studied extensively so far. However, such a singlet field causes cosmological difficulties [11]. In the early universe, non-zero vacuum energy in the inflaton sector largely breaks SUSY, which gives the Hubble induced mass term for the Polonyi field Z [9, 33]. Since Z is singlet under any symmetry, the minimum determined by the Hubble induced term deviates from the true minimum. The deviation is generically expected to be of the order of the fundamental energy scale of the supergravity, which we assume to be the Planck scale. Considering these SUSY breaking effects, the Polonyi field evolves as follows. During the inflation, the Polonyi field Z sits at the minimum determined by the Hubble induced term. After the end of the inflation, the Hubble parameter decreases to m_Z , and then Z begins to oscillate around the true minimum with an amplitude of the order of the Planck scale. The energy density of the oscillating field Z scales as a^{-3} , where a is the scale factor, and that of radiation scales as a^{-4} . Thus, the oscillating Polonyi field immediately dominates the universe after an inflaton decays into radiation.

When the Polonyi field starts its oscillation before the inflaton decays, the Polonyi energy density-to-entropy ratio after the inflaton decay is given by

$$\frac{\rho_Z}{s} \simeq \frac{T_{\text{inf}}}{8} \left(\frac{z_0}{M_{\text{pl}}} \right)^2, \quad (2.2)$$

where T_{inf} denotes the temperature when the inflaton decays, s denotes the entropy density, and z_0 represents the oscillation amplitude of Z , which is expected to be of the order of the Planck scale. With the decay rate of the Polonyi field,

$$\Gamma_Z = \frac{d_Z}{8\pi} \frac{m_Z^3}{M_{\text{pl}}^2}, \quad (2.3)$$

³Hereafter, we use the same letter Z to denote the scalar component of the supermultiplet.

the decay temperature is given by

$$T_Z = \left(\frac{90}{\pi^2 g_*(T_Z)} \right)^{1/4} \sqrt{\Gamma_Z M_{\text{pl}}} \simeq 4 \text{ MeV} \times d_Z^{1/2} \left(\frac{m_Z}{100 \text{ TeV}} \right)^{3/2}, \quad (2.4)$$

where d_Z is a numerical constant, and $g_*(T_Z)$ is the effective number of degrees of freedom at a temperature of T_Z . Here, we used $g_*(T_Z) = 10.75$. One can find that if m_Z is smaller than $\mathcal{O}(100)$ TeV, the Polonyi field decays during and after the BBN. In this case, $\rho_Z/s \lesssim \mathcal{O}(10^{-14}-10^{-11})$ GeV is required in order not to upset the success of the BBN [34, 35]. From eq. (2.2), one can find that the constraint from the BBN cannot be avoided without dilution even if $T_{\text{inf}} \simeq \mathcal{O}(10)$ MeV. The required dilution factor Δ for $\rho_Z/s \lesssim 10^{-14}$ GeV is as follows:

$$\Delta \equiv \frac{s_f a^3}{s_i a^3} \gtrsim 1.3 \times 10^{22} \left(\frac{T_{\text{inf}}}{10^9 \text{ GeV}} \right), \quad (2.5)$$

where s_f and s_i denote the entropy density after and before the entropy production, respectively. Such huge entropy production, however, also dilutes pre-existing baryon asymmetry. We then need an intricate model in which the baryon asymmetry is produced after the dilution [14–16].

There is another simple way to avoid the problem, which we focus on in this paper.⁴ When the Polonyi is as heavy as $\mathcal{O}(100)$ TeV, it decays before the onset of the BBN, and the constraint becomes much milder [34, 35]. Even in this case, however, there is an incidental problem: LSPs are abundantly produced from the decay of the Polonyi, and the LSP density tends to exceed the observed DM density.⁵ The abundance of the LSPs is given by

$$Y_{\text{LSP}} \equiv \frac{n_{\text{LSP}}}{s} \simeq \min \left[\frac{\Gamma_Z}{\langle \sigma v \rangle s(T_Z)}, \frac{N_{\text{LSP}} n_Z(T_Z)}{s(T_Z)} \right], \quad (2.6)$$

where n_{LSP} and n_Z denote the number density of the LSP and the Polonyi field, respectively. $\langle \sigma v \rangle$ represents a thermally averaged cross section of the pair annihilation. N_{LSP} is the averaged number of superparticles produced by the decay of one Polonyi field. When the first term is relevant, the pair annihilation between LSPs proceeds after the Polonyi decay, and the relic LSP density is approximately proportional to T_Z^{-1} . For example, when the neutral wino is the LSP with a mass of $\mathcal{O}(0.1-1)$ TeV, the decay temperature T_Z is needed to be larger than $\mathcal{O}(1-10)$ GeV [18] in order for the wino abundance not to exceed the observed DM density.⁶ This requires that m_Z is larger than $\mathcal{O}(5000)$ TeV. Assuming that the gravitino mass is generically comparable to m_Z and that the wino mass is generated by the anomaly-mediation, such a heavy Polonyi mass is incompatible with the wino mass of $\mathcal{O}(0.1-1)$ TeV.

In this paper, we consider the case where $m_Z \simeq \mathcal{O}(100)$ TeV and the LSP density is diluted by entropy production. As we will show, it is possible to generate the baryon asymmetry before the dilution since the required dilution factor is much smaller than eq. (2.5) for $m_Z = \mathcal{O}(100)$ TeV. The most probable candidate for baryogenesis is the Affleck-Dine mechanism because it can create huge baryon number to survive the dilution.⁷

⁴For other solutions, see refs. [36–41].

⁵The R-parity violation makes the LSP unstable, and it serves to avoid the LSP overproduction. In this paper, however, we try to construct the scenario under the assumption of the R-parity conservation.

⁶Reference [18] has taken into account Sommerfeld effect and coannihilation among charged and neutral winos.

⁷Another possibility is to generate the baryon asymmetry after the dilution, e.g. by the electroweak baryogenesis [6] and primordial magnetic field [42, 43].

3 Affleck-Dine mechanism without superpotential

The Affleck-Dine mechanism [5] is a promising candidate for the baryogenesis in cosmological scenarios with dilution. In this section, we briefly explain the AD mechanism in the case where the AD field does not appear in the superpotential. We show that the resultant baryon number density is comparable to the number density of the Polonyi field.

The minimal SUSY Standard Model (MSSM) contains a lot of flat directions which have no scalar potentials at the renormalizable level and in SUSY limit [4]. In the AD mechanism, a flat direction with a $B - L$ charge creates the baryon asymmetry. We call it “the Affleck-Dine field”. In the early universe, SUSY breaking effects and non-renormalizable terms affect its evolution. In particular, A -term scalar potentials violating $B - L$ global symmetry rotate the AD field in the complex plane, and can effectively generate $B - L$ number [9]. The $B - L$ asymmetry is converted into the baryon asymmetry through the sphaleron process [6–8].

In order to generate huge baryon number comparable to the number density of the Polonyi field, the AD field value at the onset of its oscillation should be of the order of the Planck scale. To obtain the large field value, we assume that the AD field does not appear in the superpotential.⁸ In this case, SUSY breaking effects including the A -terms are provided by the Kähler potential. We consider the following terms in the Kähler potential:⁹

$$\begin{aligned} \mathcal{L}_{\text{AD}} &= \int d^2\theta d^2\bar{\theta} \left[-3M_{\text{pl}}^2 \exp\left(-\frac{K}{3M_{\text{pl}}^2}\right) \right] \\ &\supset \int d^2\theta d^2\bar{\theta} \left[f_1 |\Phi|^2 + \left(f_2 \frac{\Phi^n}{nM_{\text{pl}}^{n-2}} + \text{h.c.} \right) + f_3 \frac{|\Phi|^4}{M_{\text{pl}}^2} + \dots \right], \end{aligned} \quad (3.1)$$

where Φ denotes the AD field, and f_i ($i = 1, 2, 3$) is an arbitrary real function of Z and Z^\dagger which satisfies $f_i = f_i^\dagger$ for $i = 1, 3$. From these terms, the potential for the AD field Φ is given by

$$V(\Phi) = m_\phi^2 |\Phi|^2 - \frac{m_{3/2}^2}{nM_{\text{pl}}^{n-2}} (a_n \Phi^n + \text{h.c.}) + c_4 m_{3/2}^2 \frac{|\Phi|^4}{M_{\text{pl}}^2} + \dots, \quad (3.2)$$

where m_ϕ and $m_{3/2}$ denote the soft scalar mass of the AD field and the gravitino mass, respectively. a_n and c_4 are $\mathcal{O}(1)$ dimensionless parameters. Higher order terms are expressed by the ellipsis, and we will neglect them for simplicity.

As we will discuss later, m_ϕ is assumed to be smaller than $m_{3/2}$. When $m_\phi \ll m_{3/2}$ and $c_4 = 0$, there exists charge/color breaking minima smaller than the Planck scale because of the relatively large A -terms [44]. The AD field needs to avoid dropping the minima during its evolution, which makes the AD mechanism less effective [45–47]. The quartic term, however, lifts the potential near the Planck scale, and the global minima disappear when its coefficient is positive.

In addition, the AD field acquires the so-called Hubble induced terms since the non-zero vacuum energy in the early universe largely violates SUSY [9, 33]. We assume that the AD field has a negative Hubble induced mass term so that it takes a large field value at the onset of its oscillation:

$$V_H = -c_H H^2 |\Phi|^2, \quad (3.3)$$

⁸For example, a $U(1)_R$ symmetry can prohibit appearance of the AD field in the superpotential.

⁹The second term in the second line is equivalent to a superpotential term suppressed by the gravitino mass.

where H is the Hubble parameter, and c_H is a positive dimensionless parameter of $\mathcal{O}(1)$. Due to this negative mass term, the AD field takes its field value of the order of the Planck scale until $H(t) \simeq m_{3/2}$. When $H(t) \lesssim m_{3/2}$, the position of the local minimum of the AD field, which is determined by a balance between the negative Hubble induced mass term and the positive quartic term, becomes smaller than the Planck scale. Since the position is quickly driven towards the origin, the AD field cannot track the local minimum and starts to roll down to the origin when $H(t) \simeq m_{3/2}$ (for details, see appendix A and [48]).

Let us estimate the produced baryon number density. The evolution equation for the $B-L$ density is expressed as

$$\dot{n}_{B-L} + 3Hn_{B-L} = 2\beta \text{Im} \left[\frac{\partial V}{\partial \Phi} \Phi \right], \quad (3.4)$$

where n_{B-L} expresses the $B-L$ density and β denotes the $B-L$ charge of the AD field. The dot denotes the time derivative. The right-hand side is the source of the $B-L$ asymmetry. By solving the evolution equation, one can find that the asymmetry is produced most effectively at the onset of the oscillation ($H \simeq m_{3/2}$). After that, the AD field value decreases due to the expansion of the universe, and $B-L$ violating effects become negligible. The produced $B-L$ density is then estimated as

$$\begin{aligned} n_{B-L}(t_{\text{osc}}) &\sim 2\beta |a_n| \sin [n\theta_i + \arg(a_n)] \frac{m_{3/2}^2 |\Phi_{\text{osc}}|^n}{H_{\text{osc}} M_{\text{pl}}^{n-2}} \\ &\equiv \epsilon H_{\text{osc}} |\Phi_{\text{osc}}|^2, \end{aligned} \quad (3.5)$$

where the subscripts of $_{\text{osc}}$ show the values when the AD field starts to oscillate. θ_i is the initial phase of the AD field. ϵ is estimated as

$$\epsilon \simeq 2\beta |a_n| \sin [n\theta_i + \arg(a_n)] \frac{m_{3/2}^2 |\Phi_{\text{osc}}|^{n-2}}{H_{\text{osc}}^2 M_{\text{pl}}^{n-2}}. \quad (3.6)$$

As mentioned above, the AD field starts to roll down to the origin from the field value of the order of the Planck scale when $H(t) = H_{\text{osc}} \simeq m_{3/2}$. Therefore, the $B-L$ asymmetry is estimated as $n_{B-L}(t_{\text{osc}}) \simeq \epsilon m_{3/2} |\Phi_{\text{osc}}|^2$, where $\epsilon \simeq \mathcal{O}(1)$ and $|\Phi_{\text{osc}}| \simeq M_{\text{pl}}$.

After the oscillation, the AD field decays into radiation, and the conserved $B-L$ asymmetry is converted into the baryon asymmetry through the sphaleron effect [6–8]. The baryon asymmetry is related to the $B-L$ asymmetry as

$$n_B = \frac{8}{23} n_{B-L}, \quad (3.7)$$

where n_B expresses the baryon number density. As mentioned in section 2, entropy production is often necessary in order to avoid the overproduction of LSPs by the decay of the Polonyi. We consider the case where the density of the Polonyi field is diluted by the entropy production. Then, the baryon asymmetry is also diluted. Assuming that the inflaton decays after the onset of the oscillation of the AD field, the yield of the baryon number is estimated as

$$Y_B \equiv \frac{n_B}{s} = \frac{8}{23} \frac{1}{\Delta} \left. \frac{3T_{\text{inf}} n_{B-L}}{4\rho_{\text{inf}}} \right|_{\text{osc}} \simeq \frac{2}{23} \frac{\epsilon}{\Delta} \frac{T_{\text{inf}}}{m_{3/2}} \left(\frac{|\Phi_{\text{osc}}|}{M_{\text{pl}}} \right)^2, \quad (3.8)$$

where ρ_{inf} denotes the energy density of the oscillating inflaton. Note that the baryon number density is comparable to the density of the Polonyi field because both the AD field and

the Polonyi field simultaneously begin their oscillation with the same amplitude of the order of the Planck scale.

Let us make a comment on Q-ball formation. When the potential for the AD field is shallower than a quadratic potential, the AD field fragments into non-topological solitons, called Q-balls [49], just after the onset of the oscillation [50–52]. Since Q-balls absorb the produced $B - L$ charge, the formation of Q-balls could significantly affect the estimation of the baryon asymmetry. In our scenario, the AD field value at the onset of the oscillation is as large as the Planck scale. Then, the formed Q-balls may be too large to decay before the BBN if Q-ball formation occurs, which renders the AD mechanism ineffective. Hence, the beta function for the soft mass of the AD field may need to be positive in order to prohibit the Q-ball formation. This requires the AD field to involve scalar fields which have large Yukawa couplings.

4 Baryon-to-DM ratio

In this section, we show that the baryon-to-DM ratio is simply given by the LSP mass and the branching fraction of the decay of the the Polonyi into superparticles in our scenario. We also explain that our scenario is realized in the so-called sequestering model [20, 21] with a (pseudo-)NGB, which can be identified with the QCD axion.

4.1 Scenario

Before we calculate the baryon-to-DM ratio, let us summarize our scenario. When $H(t) \simeq m_Z \simeq m_{3/2} \simeq \mathcal{O}(100)$ TeV, both the Polonyi and the AD fields roll down to their origins with the amplitudes of the order of the Planck scale. At that time, the AD field generates the $B - L$ asymmetry which is later converted to the baryon asymmetry by the sphaleron process. Then, the entropy production occurs and dilutes both the Polonyi density and the baryon asymmetry. After the dilution, the Polonyi field decays into superparticles which consequently decay into LSPs before the epoch of the BBN. Thus, the DM density is determined by the abundance of the nonthermally produced LSPs, assuming that thermal relic density of LSPs is negligible.

First, we estimate the DM abundance. From eq. (2.2), the LSP-to-entropy ratio is estimated as

$$\frac{\rho_{\text{LSP}}}{s} = m_{\text{LSP}} \frac{2\text{Br}_{\text{SUSY}} n_Z}{\Delta s_i} \simeq \frac{\text{Br}_{\text{SUSY}} T_{\text{inf}} m_{\text{LSP}}}{\Delta 4m_Z} \left(\frac{z_0}{M_{\text{pl}}} \right)^2, \quad (4.1)$$

where Br_{SUSY} denotes a branching fraction for the Polonyi decay into two superparticles,¹⁰ and m_{LSP} denotes the LSP mass. The number of the produced superparticles is almost equal to that of the LSPs due to the R -parity conservation. Note that the pair annihilation between LSPs is not efficient since the Polonyi field decays after the dilution. We assume that decay products other than the LSPs do not contribute to the DM abundance. Hence, the DM abundance is obtained by the nonthermally produced LSPs.

Next, let us compare the DM abundance with the baryon asymmetry created by the AD mechanism. The ratio of the density of the Polonyi field to the $B - L$ number remains the same after they begin their oscillations since the densities of both components decrease as a^{-3} . From eqs. (3.8) and (4.1), we obtain the following relation:

$$\frac{\Omega_B}{\Omega_{\text{LSP}}} = \frac{8}{23} \frac{\epsilon}{\text{Br}_{\text{SUSY}}} \frac{m_p m_Z}{m_{\text{LSP}} m_{3/2}} \left(\frac{|\Phi_{\text{osc}}|}{z_0} \right)^2$$

¹⁰Gravitinos are not produced from the Polonyi decay assuming that the decay is kinematically forbidden ($m_Z < 2m_{3/2}$). The abundance of gravitinos produced during the reheating becomes negligible after the dilution.

$$\simeq 0.33\epsilon \left(\frac{10^{-3}}{\text{Br}_{\text{SUSY}}} \right) \left(\frac{1 \text{ TeV}}{m_{\text{LSP}}} \right) \left(\frac{|\Phi_{\text{osc}}|}{z_0} \right)^2, \quad (4.2)$$

where m_p represents the proton mass ($m_p \simeq 0.938 \text{ GeV}$). Here, we assume $m_Z \simeq m_{3/2}$. Note that z_0 and $|\Phi_{\text{osc}}|$ are of the order of the Planck scale and that ϵ is of $\mathcal{O}(1)$. One can find that the baryon-to-DM ratio is determined by the LSP mass and the branching fraction of the decay of the Polonyi into superparticles. Assuming that the LSP mass is of $\mathcal{O}(1) \text{ TeV}$, Br_{SUSY} is required to be of $\mathcal{O}(10^{-3})$ in order to realize the observed value, $\Omega_B/\Omega_{\text{DM}} \simeq 0.18$.

Our scenario needs the entropy production (see appendix B for the realization by thermal inflation). Let us estimate the required amount of the entropy production. The ratio of the observed DM density to the entropy density is given by [1]

$$\frac{\rho_{\text{DM}}^{(\text{obs})}}{s_0} \simeq 4.4 \times 10^{-10} \text{ GeV}, \quad (4.3)$$

where $\rho_{\text{DM}}^{(\text{obs})}$ denotes the observed DM energy density, and s_0 denotes the present entropy density. Comparing eq. (4.1) with eq. (4.3), the dilution factor of the required entropy production is estimated as

$$\Delta \simeq 1.9 \times 10^{12} \left(\frac{T_{\text{inf}}}{10^9 \text{ GeV}} \right) \left(\frac{m_{\text{LSP}}}{1 \text{ TeV}} \right) \left(\frac{\text{Br}_{\text{SUSY}}}{10^{-3}} \right) \left(\frac{300 \text{ TeV}}{m_Z} \right) \left(\frac{z_0}{M_{\text{pl}}} \right)^2, \quad (4.4)$$

where we assume that the inflaton decays after the Polonyi and the AD fields start to oscillate. This is the case for $T_{\text{inf}} \lesssim 10^{12} \text{ GeV} (m_Z/300 \text{ TeV})^{1/2}$. m_Z should be of $\mathcal{O}(100) \text{ TeV}$ in order to relax the constraint from the BBN. Comparing eq. (4.4) with eq. (2.5), one can find that the required dilution factor is much smaller than the case of $m_Z \lesssim \mathcal{O}(100) \text{ TeV}$. Note that the dilution factor is estimated assuming that the entropy production occurs before the Polonyi decays.

4.2 Sequestering model and decay process of the Polonyi field

When the Polonyi field decays through dimension 5 operators suppressed by the Planck scale (see eq. (2.3)), the branching fraction of the decay into superparticles is generally comparable to that into the SM particles [56]. In order to suppress the branching fraction into superparticles ($\text{Br}_{\text{SUSY}} \simeq \mathcal{O}(10^{-3})$), we consider the so-called sequestering model [20, 21],¹¹ in which the SUSY breaking sector is sequestered from the visible sector.

The Kähler potential and the superpotential are given by

$$K = -3M_{\text{pl}}^2 \log \left[1 - \frac{f_{\text{vis}} + f_{\text{hid}}}{3M_{\text{pl}}^2} \right], \quad W = W_{\text{vis}} + W_{\text{hid}}, \quad (4.5)$$

where the subscripts of $_{\text{vis}}$ and $_{\text{hid}}$ denote the visible and the hidden sectors (SUSY breaking sector), respectively. We also assume that the SM gauge sector does not directly couple to the hidden sector:

$$\mathcal{L}_{\text{gauge}} = \int d^2\theta [\tau_{\text{vis}} \mathcal{W}^a \mathcal{W}_a + \text{h.c.}], \quad (4.6)$$

where \mathcal{W}^a denotes field strength supermultiplets of the visible SM gauge sector, and τ_{vis} is a holomorphic function which depends only on visible sector fields.

¹¹The sequestering model has been introduced in the context of extra dimension [21]. It is also realized in a four dimensional strongly coupled CFT [57–60].

In this setup, gaugino masses vanish at the tree level because the Polonyi field does not appear in the gauge sector. The quantum corrections to the gaugino masses arise only at loop-suppressed level, which mainly come from the anomaly mediation [21, 22]. Then, the lightest gaugino is the neutral wino with a mass of $\mathcal{O}(1)$ TeV when $m_{3/2} \simeq \mathcal{O}(100)$ TeV. This is compatible with our scenario with the neutral wino LSP.

Soft scalar masses also vanish at the tree level when the Kähler potential is given by eq. (4.5). They acquire loop-suppressed contribution from the anomaly mediation [21, 22], Planck-suppressed interactions [25] and so on. If the MSSM scalars acquire their masses only from the anomaly mediation, slepton masses would become negative. This is problematic in terms of the phenomenology. Thus, there should be other sources to give them positive masses. One of such candidates is one loop correction from the Planck-suppressed interactions [21, 25]. When a cut-off scale is taken around the gravitational scale, one loop correction can exceed the anomaly-mediated masses which appear at two loop level.¹² In this case, sfermion masses are of $\mathcal{O}(10)$ TeV when $m_{3/2} \simeq \mathcal{O}(100)$ TeV.¹³

When the soft masses and the supersymmetric masses (μ term) of the Higgs fields are of $\mathcal{O}(10)$ TeV and of $\mathcal{O}(1)$ TeV respectively, the $B\mu$ term ($\sim \mu m_{3/2}$) is comparable to the scalar masses, which leads to the successful electroweak gauge symmetry breaking. The higgsino with mass of $\mu \sim \mathcal{O}(1)$ TeV could be the LSP instead of the neutral wino. When the soft masses are of $\mathcal{O}(1)$ TeV, the $B\mu$ term is generally too large to realize the electroweak symmetry breaking. In the Next-to-MSSM,¹⁴ however, the supersymmetric Higgs mass term is generated as the breaking term of the scale invariance, and the (effective) $B\mu$ term appears at the loop-suppressed level.

Since the SUSY breaking sector is now sequestered from the AD field, the functions f_i ($i = 1, 2, 3$) in eq. (3.1) do not contain the Polonyi field.¹⁵ Even in this case, the potential for the AD field involves the holomorphic A -terms and the quartic term of $\mathcal{O}(m_{3/2}^2)$ due to the explicit breaking of the conformal symmetry. By requiring that the vacuum energy vanishes, the coefficients in eq. (3.2), a_n and c_4 , are estimated as $a_n \simeq -f_2(n-1)$ and $c_4 \simeq f_3$ when $f_1 = 1$. Note that the estimated values contain uncertainties of $\mathcal{O}(1)$.

Let us consider the decay process of the Polonyi field (for details, see [56, 72]). Firstly, the Polonyi field generally decays into 2 gravitinos at the tree level when $m_Z > 2m_{3/2}$. This decay process is incompatible with our scenario since the branching fraction of the decay of the Polonyi into superparticles is required to be of $\mathcal{O}(10^{-3})$. Hence, we assume that the decay into 2 gravitinos is kinetically forbidden ($m_Z < 2m_{3/2}$).

The decay into matter scalars comes from the kinetic terms for the sequestered potential: $\mathcal{L}_K = g_{ij^*} \partial_\mu \phi^i \partial^\mu \phi^{*j}$, where ϕ^i denotes the matter scalar fields, and $g_{ij^*} = \frac{\partial^2 K}{\partial \phi^i \partial \phi^{*j}}$. The kinetic terms are converted into the following form up to a total derivative: $\mathcal{L} \sim \frac{Z}{M_{\text{pl}}} \phi^i \partial^2 \phi^{*j}$. Using the equation of motion, interaction terms from the kinetic terms are proportional to the scalar mass squared. Thus, the branching fraction of the decay mode $Z \rightarrow \phi^i \phi^{*j}$ is

¹²When the one loop correction determines scalar masses, the mass spectra of MSSM scalar particles become UV sensitive, which is contrary to the anomaly-mediated masses. Thus, we lose a solution to the SUSY FCNC problem unless the universality condition is imposed at the UV scale. There also exists other UV insensitive models which solve the negative slepton mass problem [61–65].

¹³The lightest Higgs boson mass acquires radiative corrections from stop one loop diagrams [66–70]. Stop mass of $\mathcal{O}(10)$ TeV is compatible with the relatively heavy observed Higgs boson mass of 126 GeV.

¹⁴For a review see ref. [71].

¹⁵In the early universe, the inflaton sector breaks the SUSY, which generates the Hubble induced mass term. In order to generate the negative Hubble induced mass term for the AD field, the inflaton sector should not be sequestered from the visible sector.

suppressed by a factor of $\mathcal{O}(m_\phi^4/m_Z^4) \sim \mathcal{O}(10^{-4}-10^{-5})$ when the scalar mass is smaller than the Polonyi mass. The Polonyi field also decays into matter scalar fields through one-loop diagrams by Planck-suppressed interactions, but the rates of these decays are the same order with that of the tree-level decay. Similarly, the branching fraction into matter fermions is proportional to fermion mass squared and is negligible. The decay rate into higgsinos with masses of $\mu \sim \mathcal{O}(1)$ TeV is the same order with that into matter scalar fields since it is suppressed by a factor of $\mathcal{O}(\mu^2/m_Z^2) \sim \mathcal{O}(10^{-4})$.

Decay into three-body final states is suppressed for the sequestered potential. In general, the decay of $Z \rightarrow \phi^i \chi^j \chi^k$, where χ^i denotes the matter fermions, occurs through the following interaction:

$$\mathcal{L}_{\text{three}} = -\frac{1}{2} e^{\frac{K}{2M_{\text{pl}}^2}} \left(\frac{K_Z}{M_{\text{pl}}^2} W_{ijk} - 3\Gamma_{Zi}^l W_{jkl} \right) Z \phi^i \chi^j \chi^k + \text{h.c.}, \quad (4.7)$$

where the subscripts represent the derivative by the scalar fields, and $\Gamma_{jk}^i = g^{il*} g_{jl*k}$. One can find that this term vanishes if the Kähler potential is given by the form of eq. (4.5). For the same reason, the decay of $Z \rightarrow \phi^i \phi^j \phi^k$ does not occur, either.

Since the Polonyi field is not directly coupled with the gauge sector, it does not decay into gauge bosons and gauginos at the tree level. However, it can decay into them through the anomaly-mediated effects. When the mass of the Polonyi field is dominated by a supersymmetric mass term, the interaction terms between Z and the gaugino λ are given by [73]

$$\mathcal{L}_{\text{anomaly}} = \frac{\alpha_g b_0 m_Z}{24\pi M_{\text{pl}}} \frac{K_Z}{M_{\text{pl}}} Z^* \lambda \lambda + \text{h.c.}, \quad (4.8)$$

where $\alpha_g = g^2/4\pi$ represents a gauge coupling constant, and $b_0 = 3T_G - T_R$ is the coefficient of the beta function. Since the SUSY breaking mass term is comparable to the supersymmetric mass term, the interaction terms are deviated from eq. (4.8) by $\mathcal{O}(1)$. From eq. (4.8), the decay rate is estimated as [73]

$$\Gamma(Z \rightarrow 2\lambda) \simeq \frac{N_g \alpha_g^2 b_0^2}{4608\pi^3} \frac{|K_Z|^2}{M_{\text{pl}}^2} \frac{m_Z^3}{M_{\text{pl}}^2}, \quad (4.9)$$

where N_g is the number of gauginos. The decay rate of Z into 2 gauge bosons is also the same as eq. (4.9). The most important process is the decay into gluons and gluinos. We can estimate its rate by using $N_g = 8$ and $b_0 = 3$.

In summary, the Polonyi field mainly decays into gluinos and gluons through the anomaly-mediated effects for the sequestered Kähler potential. If it is the leading process, however, the Polonyi field becomes long-lived, and the constraint from the BBN is again severe even with $m_Z \simeq \mathcal{O}(100)$ TeV. We need some other efficient decay processes. Note that those decay processes should not yield large DM abundance. As a suitable decay process, we consider the decay of the Polonyi field into a (pseudo-)NGB. To be specific, we introduce the QCD axion [26–29]. The axion is a pseudo-NGB associated with the spontaneous breaking of the Peccei-Quinn (PQ) symmetry and appears as a phase direction of the PQ field:

$$P = v_{\text{PQ}} \exp\left(\frac{s + ia}{\sqrt{2}v_{\text{PQ}}}\right), \quad (4.10)$$

where P , s and a denote the PQ field, the saxion field and the axion field, respectively. v_{PQ} is the VEV of the PQ field.

Let us assume that the PQ field also belongs to the visible sector, which is natural as the PQ field must directly couple to SM charged particles. The PQ field interacts with the Polonyi field through the kinetic terms as follows:

$$\mathcal{L} \sim \frac{Z}{M_{\text{pl}}} P \partial^2 P^* + \text{h.c.} \quad (4.11)$$

After the PQ symmetry breaking, the axion field appears as a massless direction as eq. (4.10). Expanding of eq. (4.11) in terms of the saxion and axion leads to mixing of the kinetic terms between the Polonyi field and (s)axion. In order to estimate the rate of the Polonyi decay into axions, we need to diagonalize the kinetic terms and transform the bases into mass eigenstates. We then obtain the following interactions:

$$\mathcal{L} = c_R \frac{z_R}{\sqrt{2}M_{\text{pl}}} \partial_\mu \hat{a} \partial^\mu \hat{a} + c_I \frac{z_I}{\sqrt{2}M_{\text{pl}}} \partial_\mu \hat{a} \partial^\mu \hat{a}, \quad (4.12)$$

where c_R and c_I are coefficients of $\mathcal{O}(1)$, and z_R and z_I denote a real component and an imaginary component of the Polonyi field ($Z = \frac{1}{\sqrt{2}}(z_R + iz_I)$), respectively. Here, we used \hat{a} to show the mass eigenstate of the massless direction.

The rate of the Polonyi decay into axions¹⁶ is estimated as

$$\Gamma(Z \rightarrow 2 \text{ axions}) = \frac{c_a^2}{64\pi} \frac{m_Z^3}{M_{\text{pl}}^2}, \quad (4.13)$$

where we define c_a as $c_a^2 \equiv c_R^2 + c_I^2$. The Polonyi decay temperature T_Z is estimated as

$$T_Z \simeq 7.1 \text{ MeV} c_a \left(\frac{m_Z}{300 \text{ TeV}} \right)^{3/2}. \quad (4.14)$$

Note that the Polonyi density does not dominate the universe at its decay since we assume that the dilution occurs before the decay. Even when the Polonyi field is a subdominant component of the universe, it must decay before the BBN in order not to destroy synthesized light elements. Hence, the Polonyi should be as heavy as $\mathcal{O}(100)$ TeV.

The rate of the loop-suppressed decay (eq. (4.9)) is much smaller than that of the tree-level decay (eq. (4.13)), and we obtain the branching fraction of the Polonyi decay into superparticles as

$$\text{Br}_{\text{SUSY}} = \frac{\Gamma(Z \rightarrow 2 \text{ superparticles})}{\Gamma(Z \rightarrow 2 \text{ axions})} \sim 1 \times 10^{-3}. \quad (4.15)$$

Since the axion mass is typically much smaller than the LSP mass, the abundance of axions produced from the Polonyi decay is negligible compared with the LSP abundance.¹⁷ Therefore, the DM abundance is determined by the abundance of the decay products of the suppressed decay into superparticles. The axion also gives a negligible contribution to the dark radiation. From eqs. (4.2) and (4.15), it is found that the observed baryon-to-DM ratio of $\Omega_B/\Omega_{\text{DM}} \simeq 0.18$ is explained in the sequestering model with the (pseudo-)NGB.

¹⁶The decay products of the Polonyi field could contain the axino which is a superpartner of the axion. Since the decay into axinos could lead to the overproduction of LSPs, we assume that such a decay process is kinematically forbidden, $2m_{\tilde{a}} > m_Z$, where $m_{\tilde{a}}$ represents the axino mass.

¹⁷We also assume that the density of the coherent oscillation of the axion field does not exceed the observed DM density, which implies $v_{\text{PQ}}/N_{\text{DW}} \sim 10^{9-13}$ GeV.

5 Summary and discussions

In this paper, we have considered the cosmological consistent scenario in the presence of the heavy Polonyi field. When the Polonyi field is heavier than $\mathcal{O}(100)$ TeV, it decays before the BBN, and the constraint from the BBN becomes much milder. LSPs produced by the decay of the Polonyi are diluted by entropy production. With the entropy production, the most promising candidate for the baryogenesis is the AD mechanism.

In our scenario, the DM is explained by the LSPs produced from the decay of the Polonyi field, and the baryon asymmetry is created by the AD mechanism without superpotential. We have shown that the baryon-to-DM ratio is simply determined by the LSP mass (m_{LSP}) and the branching fraction of the decay of the Polonyi into DM (Br_{SUSY}). The observed ratio $\Omega_B/\Omega_{\text{DM}} \simeq 0.18$ is realized for the LSP mass of $\mathcal{O}(1)$ TeV and the branching fraction of $\mathcal{O}(10^{-3})$. The Polonyi density is diluted before its decay by the entropy production, for example, by thermal inflation (see appendix B).

In general, the branching fraction, Br_{SUSY} , is of $\mathcal{O}(1)$ when the Polonyi field is connected with the observable sector through non-renormalizable interactions. In order to suppress Br_{SUSY} ($\simeq \mathcal{O}(10^{-3})$), we have considered the sequestering model with a (pseudo-) NGB, which can be identified with the QCD axion. In this model, the Polonyi field decays into NGBs at the tree level, which do not contribute to the DM abundance. On the other hand, it decays into superparticles mainly through anomaly-induced interactions and hence is suppressed compared with the decay into NGBs.

The phenomenological consequences of the model is as follows. The model predicts the existence of the wino DM and the gluino at a TeV scale, and they would be probed by DM search experiments or colliders. When a scalar mass is mainly generated from anomaly-mediated SUSY breaking effects and is of the $\mathcal{O}(1)$ TeV, a scalar particle may be discovered by the LHC search. When the origin of scalar masses depends on the UV structure, there may be CP violation or flavor violation. The model also predicts the existence of a (pseudo-) NGB.

Let us comment on the implication of our study to the cosmological moduli problem [74, 75]. There may exist other singlet scalar fields called “moduli fields”, which are motivated from UV physics such as the superstring theory. The late-time decay of those fields could also cause cosmological difficulties: as is the case with the Polonyi field, the abundance of nonthermally produced LSPs could exceed the observed DM density even when the moduli field is heavy enough to decay before the BBN. In that case, the entropy production is needed. However, enough baryon asymmetry cannot be produced even with the AD mechanism because the branching fraction Br_{SUSY} is generally of the order unity for generic moduli fields (see eq. (4.2)). Note that the AD mechanism discussed in this paper works the most effectively. The cosmological moduli problem is very severe from the viewpoint of baryogenesis.

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A Evolution of the AD field

Following ref. [9, 48], we make a remark about the evolution of the AD field in the potential given by eq. (3.2). In particular, we pay attention to the onset of the oscillation of the AD field.

Including the negative Hubble induced mass term, the potential for the radial component of the AD field, $\phi \equiv |\Phi|/\sqrt{2}$, is given by

$$V(\phi) \simeq -\frac{c_H}{2} H^2 \phi^2 + c_4 m_{3/2}^2 \frac{\phi^4}{4M_{\text{pl}}^2} + \dots, \quad (\text{A.1})$$

where we omit the soft SUSY breaking mass term assuming that $H \gg m_\phi$. When $H \gtrsim m_{3/2}$, the AD field sits at some value of the order of the Planck scale. This is because higher dimensional operators lift the scalar potential. When $m_\phi \lesssim H \lesssim m_{3/2}$, local minimum determined by the potential is given by

$$\phi_{\text{min}} \simeq \sqrt{\frac{c_H}{c_4}} \frac{M_{\text{pl}} H}{m_{3/2}}. \quad (\text{A.2})$$

One can find that ϕ_{min} decreases as $a^{-3/2}$ from the Planck scale during the epoch of the inflaton oscillation. When ϕ_{min} starts to decrease when $H \simeq m_{3/2}$, the AD field cannot track the local minimum and then starts to roll down to the origin.

In order to look more closely at the onset of the oscillation of the AD field, we write down the equation of motion for ϕ as

$$\ddot{\phi} + 3H\dot{\phi} - c_H H^2 \phi + c_4 m_{3/2}^2 \frac{\phi^3}{M_{\text{pl}}^2} = 0. \quad (\text{A.3})$$

Using the number of e-folding $N \equiv \ln(a/a_i)$ as a time variable, this is rewritten as

$$\frac{d^2}{dN^2} \phi + \frac{3}{2} \frac{d}{dN} \phi - c_H \phi + \frac{c_4 m_{3/2}^2}{M_{\text{pl}}^2 H^2} \phi^3 = 0, \quad (\text{A.4})$$

where we take a_i as the scale factor when $H_i \simeq m_{3/2}$. Rescaling the AD field value ϕ as

$$\psi \equiv \frac{\phi}{\phi_{\text{min},i}} e^{\frac{3N}{2}}, \quad \phi_{\text{min},i} = \sqrt{\frac{c_H}{c_4}} \frac{M_{\text{pl}} H_i}{m_{3/2}}, \quad (\text{A.5})$$

we can eliminate the dependence on the time variable in the coefficients. In terms of ψ , the equation of motion is rewritten as

$$\frac{d^2}{dN^2} \psi - \frac{3}{2} \frac{d}{dN} \psi - c_H \psi + c_H \psi^3 = 0. \quad (\text{A.6})$$

Note that the coefficient of the friction term is negative. This implies that the AD field cannot track the local minimum and starts to oscillate around the origin.¹⁸ $B-L$ asymmetry is effectively produced at the onset of the oscillation ($H \simeq m_{3/2}$). After that, the asymmetry is conserved since the amplitude of the AD field decreases, and the $U(1)_{B-L}$ symmetry breaking terms become ineffective.

¹⁸Reference [48] numerically confirms this behavior in the context of the evolution of the PQ field.

B Thermal inflation

Our scenario needs the entropy production by some mechanism in order to dilute the Polonyi density and the baryon asymmetry, which requires another component to dominate the energy density of the universe. Its energy density must be constant or decrease more slowly than that of the oscillating Polonyi field. In this appendix, we introduce thermal inflation [12, 13] as one example of dilution mechanisms.

Let us introduce a specific model of the thermal inflation. The thermal inflation is a short epoch of accelerated expansion of the universe at a low-energy scale. This mechanism requires a scalar field corresponding to a flat direction at the renormalizable level, which is called “the flaton”. We assume an approximate Z_4 symmetry and the superpotential given by

$$W = \frac{\lambda_X}{4M_{\text{pl}}} X^4 + g_\xi X \xi \bar{\xi}, \quad (\text{B.1})$$

where λ_X and g_ξ are dimensionless coupling constants, and X is the supermultiplet of the flaton field with a Z_4 charge of 1. ξ and $\bar{\xi}$ are massless $SU(3)_C$ gauge charged fields. Note that X is singlet under the SM gauge symmetry. The massless gauge charged fields interact with thermal bath, which generates the thermal mass term for X . Here, we ignore higher dimensional terms since we focus on the field value much smaller than the Planck scale.

Including SUSY breaking effects and the thermal mass term, the potential is given by

$$V(X) = V_0 + (c_T T^2 - m_0^2) |X|^2 + \frac{m_{3/2}}{4M_{\text{pl}}} (\lambda_X X^4 + \text{h.c.}) + \frac{|\lambda_X|^2}{M_{\text{pl}}^2} |X|^6 + \dots, \quad (\text{B.2})$$

where V_0 represents the vacuum energy which causes the thermal inflation, and c_T is a coefficient of the order of the square of g_ξ . Hereafter, we assume that X has a tachyonic mass term around the origin ($m_0 > 0$). We also assume that the flaton sector is sequestered from the SUSY breaking sector and that m_0 is smaller than $m_{3/2} \simeq \mathcal{O}(100)$ TeV.

The evolution of the flaton X is as follows: we assume that X obtains a positive Hubble induced mass term during the primordial inflation. Then, X is expected to sit around the origin just after the inflation. Even when $H \lesssim m_0$, X can be trapped around the origin due to the thermal mass term, and then the vacuum energy V_0 causes the accelerated expansion of the universe at a low-energy scale. The thermal inflation lasts until the temperature decreases to the critical value of $T_c \simeq c_T^{-1/2} m_0$. After the end of the thermal inflation, X starts to roll down to the true minimum due to the negative mass term. When H decreases to the decay rate of X , X decays into radiation with huge entropy production, and the radiation dominated universe is realized.

The true minimum of the potential is determined by the A -term and non-renormalizable terms when $m_0 \ll m_{3/2}$. The flaton VEV at present is given by

$$\langle |X| \rangle \equiv M \simeq \left(\frac{m_{3/2} M_{\text{pl}}}{3|\lambda_X|} \right)^{1/2} \simeq 4.9 \times 10^{11} \text{ GeV} |\lambda_X|^{-1/2} \left(\frac{m_{3/2}}{300 \text{ TeV}} \right)^{1/2}. \quad (\text{B.3})$$

Hereafter, we assume that λ_X is of $\mathcal{O}(1)$. One can find that the VEV is much larger than the electroweak scale. Therefore, the $SU(3)_C$ charged matter ξ ($\bar{\xi}$) with mass of the order of $g_\xi \langle |X| \rangle$ is expected to be much heavier than the electroweak scale at present.¹⁹ V_0 is

¹⁹Note that the thermal mass term exists only when $g_\xi \langle |X| \rangle \ll T$. When X is trapped at the origin during the thermal inflation, the charged matters ξ ($\bar{\xi}$) behave as relativistic particles in thermal bath.

determined as follows by requiring that the vacuum energy vanishes at the true minimum:

$$V_0 \simeq \frac{1}{18} M^2 m_{3/2}^2 \simeq 1.2 \times 10^{33} \text{ GeV}^4 |\lambda_X|^{-1} \left(\frac{m_{3/2}}{300 \text{ TeV}} \right)^3. \quad (\text{B.4})$$

When X has its large VEV, it is decomposed as

$$X = \left[M + \frac{\chi}{\sqrt{2}} \right] \exp \left(i \frac{a_\chi}{\sqrt{2} M} \right), \quad (\text{B.5})$$

where χ and a_χ are canonically normalized real scalar fields. We obtain a mass of the radial component χ around the true minimum:

$$m_\chi \simeq \sqrt{\frac{2}{3}} m_{3/2} \simeq 240 \text{ TeV} \left(\frac{m_{3/2}}{300 \text{ TeV}} \right). \quad (\text{B.6})$$

If the R symmetry was not broken, the superpotential for X would have R symmetry and the phase component a_χ would be a massless R axion. However, the R symmetry breaking, namely the non-zero VEV of the superpotential, generates the R symmetry breaking A -terms and the phase component a_χ also obtains its mass as $m_{a_\chi} \simeq \sqrt{4/3} m_{3/2}$.²⁰

After the end of the thermal inflation, the energy density of the oscillating flaton field dominates that of the universe. Thus, the reheating occurs when H decreases to the decay rate of the flaton χ . It mainly decays into gluons through one loop diagrams of ξ and $\bar{\xi}$. The decay rate is given by [53, 54]

$$\Gamma(\chi \rightarrow 2g) = \frac{1}{4\pi} \left(\frac{\alpha_s}{4\pi} \right)^2 \frac{m_\chi^3}{M^2}. \quad (\text{B.7})$$

The reheating temperature T_{RH}^χ is estimated as

$$\begin{aligned} T_{\text{RH}}^\chi &\simeq \left(\frac{90}{\pi^2 g_*(T_{\text{RH}}^\chi)} \right)^{1/4} \sqrt{\Gamma(\chi \rightarrow 2g) M_{\text{pl}}} \\ &\simeq 470 \text{ GeV} |\lambda_X|^{1/2} \left(\frac{m_{3/2}}{300 \text{ TeV}} \right) \left(\frac{\alpha_s}{0.1} \right), \end{aligned} \quad (\text{B.8})$$

where we use $g_*(T_{\text{RH}}^\chi) = 106.75$. Although χ also decays into gluinos, which might lead to overproduction of LSPs, they can annihilate before decoupling from the thermal bath. LSPs produced from the flaton field are negligible in the case of the wino/higgsino LSP, which is compatible with the sequestering models (see section 4.2).

We can estimate the dilution factor Δ as follows:

$$\begin{aligned} \Delta &= \frac{4}{3} \frac{V_0}{2\pi^2/45 g_{*s}(T_c) T_c^3 T_{\text{RH}}^\chi} \\ &\simeq 1.3 \times 10^{12} |\lambda_X|^{-1} \left(\frac{m_{3/2}}{300 \text{ TeV}} \right)^3 \left(\frac{300 \text{ TeV}}{T_c} \right)^3 \left(\frac{470 \text{ GeV}}{T_{\text{RH}}^\chi} \right). \end{aligned} \quad (\text{B.9})$$

One can find that the estimated dilution factor is the same order with the required one in eq. (4.4) when $T_c \simeq 300 \text{ TeV}$. Assuming that $m_0 \ll T_c \simeq \mathcal{O}(300) \text{ TeV}$, the coefficient c_T should be taken as $c_T \simeq g_\xi^2 \simeq (m_0/T_c)^2$.

²⁰Since $m_\chi < 2m_{a_\chi}$, the decay of χ into the phase components is forbidden.

In the context of the thermal inflation, we need to take into account a secondary oscillation of the Polonyi field. Its potential during the thermal inflation is given by

$$\begin{aligned}
 V &= \frac{1}{2}m_Z^2 z^2 + \frac{c'_H}{2}H^2(z - z'_0)^2 \\
 &= \frac{1}{2}(m_Z^2 + c'_H H^2) \left(z - \frac{c'_H H^2}{m_Z^2 + c'_H H^2} z'_0 \right)^2 + \dots, \tag{B.10}
 \end{aligned}$$

where z denotes the amplitude of Z ($z \equiv |Z|/\sqrt{2}$), and z'_0 represents the local minimum determined by the Hubble induced terms. c'_H is a coefficient of the Hubble mass term and is of $\mathcal{O}(1)$. Note that z'_0 is expected to be of the order of the Planck scale. One can find that the Polonyi field does not sit at the true minimum but at the local minimum determined by $c'_H(H_{\text{th}}/m_Z)^2 z'_0$ just after the thermal inflation, where H_{th} represents the Hubble parameter during the thermal inflation and $H_{\text{th}} \ll m_Z$. Therefore, the secondary oscillation occurs with the amplitude of $c'_H(H_{\text{th}}/m_Z)^2 z'_0$. The abundance of LSP produced from the above contribution is given by

$$\begin{aligned}
 \frac{\rho_{\text{LSP,sec}}}{s_0} &= 2\text{Br}_{\text{SUSY}} \frac{m_{\text{LSP}}}{m_Z} \frac{\rho_{Z,\text{sec}}}{s_f} \Big|_{H=\Gamma_\chi} = 2\text{Br}_{\text{SUSY}} \frac{m_{\text{LSP}}}{m_Z} \frac{3T_{\text{RH}}^\chi}{4V_0} \rho_{Z,\text{sec}} \Big|_{H=H_{\text{th}}} \\
 &\simeq 2.9 \times 10^{-19} \text{ GeV} c_H'^2 \left(\frac{\text{Br}_{\text{SUSY}}}{10^{-3}} \right) \left(\frac{T_{\text{RH}}^\chi}{470 \text{ GeV}} \right) \left(\frac{m_{\text{LSP}}}{1 \text{ TeV}} \right) \left(\frac{z'_0}{M_{\text{pl}}} \right)^2, \tag{B.11}
 \end{aligned}$$

where we use $m_Z \simeq m_{3/2}$. Since this is much smaller than the observed DM density, one can find that the density of LSPs produced from the secondary oscillation is negligible.

Let us comment on domain walls. Since the superpotential of X has the Z_4 symmetry, there exists four degenerate minima in the potential for X (see eq. (B.2)). The flaton randomly falls into one of them after the thermal inflation, and then domain walls (DWs) are formed. The DWs dominate the energy density of the universe, which leads to a cosmological disaster [55]. A bias for the degenerate minima is needed so that DWs collapse before they dominate the universe. In order to avoid the DW problem, the following condition should be satisfied:²¹

$$\delta V_{\text{bias}} \gg \frac{\sigma^2}{M_{\text{pl}}^2}, \tag{B.12}$$

where δV_{bias} represents the bias of the energy density, and σ represents the tension of DWs. The necessary bias is so small that it does not change the scenario of the thermal inflation.

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²¹Assuming that the energy density of DWs obeys the scaling relation $\rho_{\text{DW}} \sim \sigma H$, DWs dominate the universe when $H_{\text{dom}} \sim \sigma/M_{\text{pl}}^2$. On the other hand, DWs decay due to the bias when $H_{\text{dec}} \sim \delta V_{\text{bias}}/\sigma$. The condition is derived by requiring $H_{\text{dec}} \gg H_{\text{dom}}$.

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