Opening the window to the cogenesis with Affleck-Dine mechanism in gravity mediation

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The observed baryon and dark matter densities are equal up to a factor of 5. This observation indicates that the baryon asymmetry and dark matter have the same origin. The Affleck-Dine baryogenesis is one of the most promising mechanisms in this context. Q balls, which are often formed in the early Universe associated with the Affleck-Dine baryogenesis, decay both into supersymmetric particles and into quarks. Recently, it was pointed out that annihilation of squarks into quarks gives a dominant contribution to the Q-ball decay rate and the branching ratio of Q-ball decay into supersymmetric particles changes from the previous estimate. In this paper, the scenario of baryon and dark matter cogenesis from Q ball in gravity mediation is revisited in respect of the improved Q-ball decay rates. It is found that the successful cogenesis takes place when a wino with mass 0.4 to 1 TeV is dark matter.

I. INTRODUCTION

The existence of the baryon asymmetry and the dark matter is a long standing challenge in cosmology and particle physics. In supersymmetric (SUSY) extensions of the Standard Model (SM), the lightest SUSY particle (LSP) is a good candidate for dark matter if the R-parity is conserved. Furthermore, the Affleck-Dine mechanism can provide the baryon asymmetry [1,2]. In the gravity-mediated SUSY breaking model, the Affleck-Dine mechanism often predicts the formation of Q balls in the early universe [3,4]. The Q ball is a spherical condensate of scalar fields. It generally consists of squarks and sleptons, and eventually decays both into quarks and into SUSY particles before the Big Bang Nucleosynthesis (BBN), and the observed baryon asymmetry is released. Through the cascade decays, the SUSY particles produced by the Q-ball decay turn into LSPs, which can account for the dark matter in the Universe. In this case, the baryon asymmetry and dark matter have the same origin and the resultant ratio of baryon to dark matter can be $O(1)$ naturally [3,4,8,12].

When we consider the case that the pair annihilation of the LSPs is ineffective and assume that the Affleck-Dine field $\phi$ takes a circular orbit in the complex $\phi$ plane, the resultant ratio of baryon to dark matter from the Q-ball decay is related only with the mass of the LSP and the branching ratio of the Q-ball decay into baryons and SUSY particles. In the previous works, the branching ratio of the Q-ball decay into SUSY particles is believed to be comparable with that into quarks [3,8,13]. In this case, the mass of dark matter should be $O(1)\text{GeV}$ [3].

However, it was pointed out that the many body processes like the squark annihilation may be dominant and then the branching ratio may change drastically [13]. In this letter, we reexamine the branching ratio into SUSY particles in respect of the many body process.

Since the effective mass of the squark inside the Q ball is smaller than that of the free squark, the Q ball cannot decay into squarks. We assume that the Q ball is kinematically allowed to decay into binos, winos (LSPs), and SM particles. When the Q-ball decay rate is saturated due to the Pauli exclusion principle [16], the branching ratio is determined only by the number of degrees of freedom in the final state. Finally, we show that the branching ratio into SUSY particles can be $O(0.01)$. By using this branching ratio, we provide a successful scenario of the baryon and dark matter cogenesis through the Q-ball decay, and show that the wino LSP with mass of 0.4 to 1 TeV can naturally explains the observed baryon to dark matter ratio in the case that the pair annihilation of the LSPs is ineffective.

This letter is organized as follows. In Sec. II, we briefly review the property of Q balls in gravity mediation. In Sec. III, we compare the saturated decay and annihilations and then derive the branching ratios. In Sec. IV, we discuss the thermal history in our scenario. Sec. V is devoted to the conclusion.

II. Q BALL PROPERTIES IN GRAVITY MEDIATION

In SUSY extensions of the standard model, there are many flat directions in the scalar potential. The flat directions are lifted by the SUSY breaking effect, and we can take the following potential for the flat direction to see the property of the Q ball in gravity mediation:

$$V = m_\phi^2 |\phi|^2 \left( 1 + K \log \frac{|\phi|^2}{M_p^2} \right),$$

where $m_\phi$ is the mass of the field $\phi$, $K$ is a constant that determines the slope of the potential, and $M_p$ is the reduced Planck mass.
where \( m_\phi \) is the mass of the flat direction and \( M_P \) is the reduced Planck mass (\( \simeq 2.4 \times 10^{18} \) GeV). In gravity mediation, \( m_\phi \) is the same order of the gravitino mass \( m_{3/2} \).

The second term in the parenthesis comes from the one-loop radiative corrections, and typically \(|K| \sim 0.01-0.1\). In many cases, the gluino loops have dominant contributions to the radiative corrections and lead to \( K < 0 \), and then there exists a Q-ball solution [14]. The energy of the Q ball \( M_Q \), the radius \( R \), the rotation speed of the field \( \omega_0 \), and the field amplitude at the center of the Q ball \( \phi_0 \) are given by

\[
M_Q \approx m_\phi(\phi_0)Q, \\
R \approx \frac{1}{|K|^{1/2}m_\phi(\phi_0)}, \\
\omega_0 \approx m_\phi(\phi_0), \\
\phi_0 \approx (2\pi^{3/2})^{-1/2}|K|^{3/4}m_\phi(\phi_0)Q^{1/2},
\]

where \( m_\phi(\phi_0) \) is the mass defined at the energy scale \( \phi_0 \).

The rotation speed \( \omega_0 \) has a further important meaning as \( \omega_0 = dM_Q/dQ \); i.e., the Q-ball energy per unit charge.

As discussed in detail in Sec. [14], the decay temperature of Q balls should be sufficiently suppressed for the pair annihilation of LSPs to be ineffective. This indicates that the charge of Q balls should be \( Q \gtrsim 10^{26} \) and thus the magnitude of the scalar field is \( \phi_0 \gtrsim 10^{13}m_\phi(\phi_0) \). At this energy scale, the mass of the flat direction \( m_\phi(\phi_0) \) is lower than the mass of squarks at the electro-weak scale due to \( K < 0 \), and the Q ball cannot decay into squarks.

### III. Q-BALL DECAY RATES INTO BINO-WINO, AND QUARKS

The fermion production rates from the Q ball have upper bounds due to the Pauli exclusion principle [16]. The upper bound of the each massless fermion flux \( j \) from the Q-ball surface is calculated as

\[
\mathbf{n} \cdot j \lesssim 2 \int \frac{d^3k}{(2\pi)^3} \theta(\omega_0/2 - |k|) \frac{\theta(k \cdot \mathbf{n})}{|k|} \mathbf{k} \cdot \mathbf{n},
\]

where \( \mathbf{n} \) is the outward-pointing normal. We double the flux and take the upper limit of integration as \( \omega_0/2 \), because one of the decay products has energy less than \( \omega_0/2 \). We obtain the upper bound for the production rate from the Q ball by multiplying Eq. (7) by the area of the Q-ball surface \( 4\pi R^2 \). The decay rate is saturated when \( g\phi_0 > \omega_0 \) for the interaction \( g\phi_0 \xi \eta \) (\( \xi, \eta \): massless fermions). The condition \( g\phi_0 > \omega_0 \) is almost always satisfied due to the large Q value (see Eq. [5]).

In the case of the massive fermion \( \chi \), the upper bound of the flux is lower than Eq. (7). We consider the process of squark \( \rightarrow \) quark \( +\chi \), and treat the quark as a massless particle. The fermion \( \chi \) can obtain the energy in the range of \([m_\chi, \omega_0]\), and the quark obtain the energy in the range of \([0, \omega_0 - m_\chi]\). Taking this into account, we just change the integral of Eq. (7) as

\[
\frac{1}{8\pi^2} \int_0^{\omega_0 - m_\chi} k^2dk,
\]

for \( \omega_0 > m_\chi > \omega_0/2 \), and as

\[
\frac{1}{8\pi^2} \left[ \int_0^{\omega_0/2} k^2dk + \int_{m_\chi}^{\omega_0/2} k^2dk \right],
\]

for \( m_\chi < \omega_0/2 \). Thus, the \( \chi \) flux is given by

\[
\mathbf{n} \cdot j_\chi \simeq \frac{\omega_0^3}{96\pi^2} \times f(m_\chi/\omega_0),
\]

\[
(f(x) \equiv \begin{cases} 4(1-x)^3 & \text{for } 1/2 < x < 1, \\ 4[(1/2)^3 + ((1/2)^3 - x^3)] & \text{for } x < 1/2. \end{cases}
\]

Q balls can also decay into quarks via heavy gluino/higgsino exchange \( \phi \phi \rightarrow qq \). This reaction rate is also saturated by the Pauli exclusion principle. The detailed discussion is given in Ref. [18]. The saturated flux is Eq. (7) with \( \omega_0 \) replaced by \( 2\omega_0 \), which is the total energy available in this process. Thus, we obtain the each quark flux as

\[
\mathbf{n} \cdot j_\text{quark} \simeq \frac{(2\omega_0)^3}{96\pi^2}.
\]

This is larger than Eq. (7) by a factor of 8. Notice that this flux is valid only for \( M > \omega_0 \), where \( M \) is the gluino/higgsino mass, and we assume it in this letter. In Appendix, we show \( N(\geq 3) \) body processes are not saturated and negligible.

Now, let us compare the branching ratios of the Q-ball decay into binos, winos, and quarks. The bino or wino production rate is given by Eq. (10), while the quark production rate is given by Eq. (12). Here we should note that since the saturated production rate is determined by the Pauli exclusion principle, the total quark production rate is Eq. (12) times the number of degrees of freedom for quarks produced in the decay. We can count it once we specify the flat direction. Hereafter, we consider the flat direction \( \delta \sum_i \delta_i \phi \xi_i \) (\( \xi, \eta \): massless fermions). The condition \( \phi_0 > \omega_0 \) is almost always satisfied due to the large Q value (see Eq. [5]).

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We conclude that the total decay rate of the Q ball and the branching ratios of the decay into quarks and bino are calculated as

\[ \sum_{\text{all}} \frac{dN}{dt} \approx \left[ 8 \times 36 \times \frac{3}{4} + f \left( \frac{m_b}{\omega_0} \right) \right] R^2 \omega_0^3 / 24 \pi, \]

(13)

\[ B_{\text{quarks}} \approx \frac{8 \times 36 \times 3/4}{8 \times 36 \times 3/4 + f (m_b/\omega_0)}, \]

(14)

\[ B_{\text{bino}} \approx \frac{f (m_b/\omega_0)}{8 \times 36 \times 3/4 + f (m_b/\omega_0)}. \]

(15)

In Eq. (13) the factor 36 comes from the degrees of freedom for colors (3), flavors (6) and chiralities (2), and the factor 3/4 comes from the $U(1)_Y$ charge conservation. We do not include the quarks from the process of squark → quark + bino, because the quarks production rates are determined by the Pauli exclusion principle and the phase space of the quarks produced by the squark decay is a subset of that of the quarks produced by the squark annihilation. The binos eventually decay into winos (LSPs). Note that the thermal relics of winos do not overclose the Universe\(^2\).

### IV. COGENESIS IN GRAVITY MEDIATION

In this section, we show that our scenario of the baryon and dark matter cogeneration works well. We consider the Affleck-Dine mechanism using the flat direction without non-renormalizable superpotential, because our scenario requires Q balls with $Q \gtrsim 10^{26}$ for the pair annihilation to be ineffective. The scalar potential for the flat direction is typically written as

\[ V(\phi) = (m_\phi^2 - c_H H^2)|\phi|^2 + \frac{m_\phi^2}{M_*^2}(a_m \phi^n + h.c.) + \frac{H^2}{M_*^2}(a_H \phi^n + h.c.) + \ldots, \]

(16)

where $M_*$ is a cut-off scale and ... denotes higher order Planck-suppressed terms. The terms proportional to $H^2$ are induced via the interaction with the inflaton, and $c_H$, $a_m$, and $a_H$ are $O(1)$ constants. Here we assume $c_H > 0$. Owing to the Hubble induced terms and higher order Planck-suppressed terms, the flat direction has a large expectation value during inflation $\phi \simeq M_*$, and then begins to oscillate and rotate around $\phi = 0$ when $H \simeq m_\phi$. Soon after the oscillation, Q balls are formed. Here we assume that the second term in Eq. (16) which kicks $\phi$ in the phase direction is large enough for $\phi$ to take a circular orbit. In this case anti-Q balls are not produced, which leads to the simple relation between baryon and dark matter densities\(^3\). The charge of the Q ball is determined by

\[ Q \sim \beta \left( \frac{M_*}{m_\phi} \right)^2 \sim 3 \times 10^{28} \left( \frac{M_*}{M_Y} \right)^2 \left( \frac{2 \text{ TeV}}{m_\phi(\phi_0)} \right)^2, \]

(17)

where $\beta = 2 \times 10^{-2}$ \(^2\).

The Q-ball decay temperature is estimated as

\[ T_d = \left( \frac{90}{4 \pi g_*} \right)^{1/4} \sqrt{\Gamma_Q M_Y}, \]

\[ \simeq 10 \text{ MeV} \left( \frac{m_\phi(\phi_0)}{2 \text{ TeV}} \right)^{1/2} \left( \frac{10^{28}}{Q} \right)^{1/2}, \]

(18)

where $g_*$ is the effective relativistic degrees of freedom at the decay time, and $\Gamma_Q = (1/Q) \sum_{\text{all}} dN/dt$ is the decay rate of the Q ball. In the second line of Eq. (18), we set $g_* = 10.75$ and $\sum_{\text{all}} dN/dt / \sim 200 \times R^2 \omega_0^3 / 24 \pi Q$. We find that the Q ball decays before the BBN but after the sphaleron process freezes out \(^2\). Winos produced from the Q-ball decay do not annihilate when the following condition is satisfied:

\[ Y_{\tilde{w}}^{(NT)} \lesssim \sqrt{\frac{45}{8 \pi^2 g_*}} \frac{1}{(\sigma v)M_Y T_d}, \]

\[ \simeq 1.1 \times 10^{-10} \times \left( \frac{m_{\tilde{b}}}{10 \text{ MeV}} \right) \left( \frac{10^{-24} \text{ cm}^3/\text{s}}{\sigma v} \right) \left( \frac{10 \text{ MeV}}{T_d} \right), \]

(19)

As mentioned above, we consider the flat direction without non-renormalizable superpotential. In this case, the Q balls dominate the Universe soon after inflation, and the baryon-to-entropy ratio is given by

\[ Y_b \simeq 10^{-10} \left( \frac{T_d}{10 \text{ MeV}} \right) \left( \frac{2 \text{ TeV}}{m_\phi(\phi_0)} \right) \left( \frac{10^4}{\Delta} \right), \]

(20)

where we include the dilution factor $\Delta$. There is some mechanism to produce entropy after the reheating of the inflation, such as thermal inflation \(^{22}\) and domain wall decay \(^{23}\). We do not specify the dilution mechanism and assume that the baryon asymmetry produced from the Q-ball decay is consistent with the observation. A dilution mechanism may also dilute the undesirable relics such as thermal relic of the stable bino. Thus, the successful bino LSP scenario may be realized as a simple extension of the present wino LSP scenario. However, most of the dilution mechanisms produce SUSY particles at the same time. This is a reason why we focus on the wino LSP scenario.

In the case of the wino LSP, the thermal relic abundance can be ignored for $m_{\tilde{w}} \ll 1$ TeV \(^{24}\). The baryon-to-dark matter ratio is determined only by the Q-ball decay rate of the Q ball and anti-Q ball. When both Q balls and anti-Q balls are formed, cogenesis re-
and SM particles kinematically, and considered the wino have assumed that the Q ball can decay into binos, winos, and quarks produced in the decay. We have assumed that the Q ball can decay into quarks is enhanced by the number of degrees of freedom for quarks produced in the decay. We have assumed that the Q ball can decay into quarks is enhanced by the number of degrees of freedom for quarks produced in the decay. We have assumed that the Q ball can decay into binos, winos, and SM particles kinematically, and considered the wino as LSP. In this case, we show that the branching into binos can be $O(0.01)$ for the $\bar{u}d\bar{d}$ flat direction and predict that the dark matter is the wino with mass of $0.4-1$ TeV.  

**Appendix: Q-ball decay rates through the $N \geq 3$ body scattering processes**

Not only the decay process but also the $N$ body scattering processes can occur in the Q ball. The rate of the charge emission from the Q ball through the $N$ body scattering process can be roughly estimated as

$$\left(\frac{dN}{dt}\right)_N \sim Q \times n_0^{-N-1} \times \Gamma_N,$$  \hspace{1cm} (A.1)

$$\Gamma_N = \int d\text{Lips} |\mathcal{M}|^2 \prod_{i=\text{initial}} \frac{1}{2E_i},$$  \hspace{1cm} (A.2)

$$d\text{Lips} \equiv (2\pi)^4 \delta \left( \sum_{j=\text{all}} p_j \right) \prod_{i=\text{final}} \frac{d^3k_i}{(2\pi)^3 2E_i},$$  \hspace{1cm} (A.3)

where $n_0 \sim \omega_0 \phi_0^2$ is the squark number density in the Q ball. Let us show that the rates of the $N$ body scattering processes are not saturated for $N \geq 3$.

The mass of the field interacting with the Q ball is $O(\phi_0)$, but the typical interaction energy is $O(\omega_0)$. Thus, we can estimate the rates of the $N$ body scattering processes in the leading order of $\omega_0/\phi_0 \sim Q^{-1/2}$. The number of particles in the final state should be minimized in the leading order as

$$N_{\text{ext}} = \begin{cases} N, & N: \text{even} \\ N + 1, & N: \text{odd}. \end{cases}$$  \hspace{1cm} (A.4)

Then, the number of fermion propagators can be counted as

$$N_{\text{prop}} = \begin{cases} 3N/2 - 2, & N: \text{even} \\ 3N/2 - 3/2, & N: \text{odd}. \end{cases}$$  \hspace{1cm} (A.5)

However, as shown in Fig. 2 there should be a factor of $M$ from the chirality flip, where $M$ is the Majorana gluino mass or the higgsino mass, and we assume $\omega_0 < M \ll \phi_0$. The number of mass insertions is

$$N_{\text{mass}} = \begin{cases} N/2, & N: \text{even} \\ (N - 1)/2, & N: \text{odd}. \end{cases}$$  \hspace{1cm} (A.6)

![FIG. 1. Solutions of Eq. (22) for the case of $m_\tilde{\omega} = m_\tilde{\chi}_1$ as a function of $m_\phi(\phi_0)$. There is no solution for $m_\phi(\phi_0) \lesssim 1.2$ TeV and are two independent solutions for $m_\phi(\phi_0) \lesssim 1.2$ TeV (green and blue lines). The red and magenta dotted lines show the two asymptotic solutions $m_\tilde{\omega} = m_\phi(\phi_0)$ and $m_\tilde{\omega} = 360$ GeV, respectively.](image1)

![FIG. 2. Examples of the diagrams for the $N$ body scattering processes.](image2)
The gauge boson is massless if it has no tree level interaction with the Q ball. Hereafter, we conservatively take the gauge boson as a massless field. Thus, from Eqs. (A.5) and (A.6), we can estimate $|\mathcal{M}|^2$ as

$$|\mathcal{M}|^2 \sim \begin{cases} \phi_0^{8-6N} M^N \omega_0^N, & N: \text{even} \\ \phi_0^{6-6N} M^{N-1} \omega_0^{N+1}, & N: \text{odd} \end{cases} \quad (A.7)$$

Here we determine the $\omega_0$ dependence from dimensional analysis. On the other hand, the kinematics is determined only by $\omega_0$. We conclude that the charge emission rates from Q ball through the $N$ body scattering process can be estimated as

$$\left(\frac{dN}{dt}\right)_N \sim Q(\omega_0 \phi_0^2)^{N-1} \Gamma_N,$$

$$\sim \begin{cases} \omega_0 Q^{1-2N} (M/\omega_0)^N, & N: \text{even} \\ \omega_0 Q^{3-2N} (M/\omega_0)^{N-1}, & N: \text{odd} \end{cases} \quad (A.8)$$

where we have used $Q \sim \phi_0^2/\omega_0^2$. We should compare this with the saturated emission rate from the Q ball $(dN/dt)_{\text{sat}} \sim \omega_0$ (see Eqs. (3) and (13)) and find that the rate is not saturated for $N \geq 3$.

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