Non-thermal Dark Matter from Affleck-Dine Baryogenesis[1]

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Abstract

In this talk we discuss the origin and nature of the dark matter in the Affleck-Dine (AD) baryogenesis. The AD baryogenesis via most of the flat directions predict formations of large Q-balls, and a great number of the lightest supersymmetric particles (LSPs) are produced nonthermally via the late-time decays of these Q-balls. In order to avoid the overclosure of the universe by these nonthermally produced LSPs, an LSP with a large pair-annihilation cross section, like Higgsino- or Wino-like neutralino, is required, instead of the standard Bino-like neutralino. This reveals new cosmologically interesting parameter regions in various SUSY breaking models, which have not attracted much attention so far.

Introduction The origins of the dark matter and the baryon asymmetry in the present universe are big puzzles in particle cosmology. In the framework of supersymmetry (SUSY), there is an ideal dark matter candidate, the lightest SUSY particle (LSP). On the other hand, the minimal SUSY standard model (MSSM) also provides an interesting mechanism to generate baryon asymmetry quite effectively, by using the flat directions in the scalar potential carrying baryon and/or lepton number: that is, the Affleck-Dine (AD) baryogenesis [3]. In this talk, we point out that the Higgsino- or Wino-like neutralino naturally becomes the dominant component of the dark matter, if either of them is the LSP and if the AD mechanism is responsible for the generation of the baryon asymmetry in the present universe. This reveals new cosmologically interesting parameter regions in various SUSY breaking models, where the most extensively studied LSP as a dark matter candidate has been the Bino-like neutralino. We also comment on the detection possibility of these nonthermal neutralino dark matter.

Affleck-Dine baryogenesis Let us start by briefly reviewing the AD baryogenesis [3, 4], adopting the flat direction $\bar{U}\bar{D}\bar{D}$, for example.[2] We assume that there exists the following nonrenormalizable operator in the superpotential.[3]

$$W = \frac{\lambda}{M_{pl}} \bar{U}\bar{D}\bar{D}\ \bar{U}\bar{D}\bar{D}\,, \tag{1}$$

[1] This talk is based on the works in Ref. [1, 2].

[2] A classification of general MSSM flat directions and discussions about other flat directions are available in Ref. [2]. Our main conclusion in this talk is applicable to other flat directions as well, except for the leptonic flat directions.

[3] The case without superpotential was also discussed in Ref. [2]. Actually, it has various attractive points: there is no cosmological gravitino problem; the baryon asymmetry and dark matter density in the present universe are determined only by the potential of the AD field, independently of the reheating temperature of the inflation. Recently, it has also been pointed out [5] that this model naturally explains the ratio of the mass density of dark matter to that of baryons.

where $M_{pl} = 2.4 \times 10^{18}$ GeV is the reduced Planck scale and λ is a coupling constant. We denote the flat direction field by ϕ hereafter, and rewrite the above superpotential as follows:

$$W = \frac{1}{6M^3}\phi^6 \,, \tag{2}$$

where we have defined the effective scale M including the coupling λ.

The relevant scalar potential for the flat direction field ϕ is given by [4]

$$V(\phi) = (m_\phi^2 - c_H H^2)|\phi|^2 + \frac{m_{3/2}}{6M^3}(a_m \phi^6 + h.c.) + \frac{1}{M^6}|\phi|^{10} \,. \tag{3}$$

Here, the potential terms proportional to the soft mass squared m_ϕ^2 and the gravitino mass $m_{3/2}$ come from the SUSY breaking at the true vacuum. We assume that the SUSY breaking is mediated by gravity, and take $m_\phi \simeq m_{3/2}|a_m| \simeq 1$ TeV.[5] The Hubble mass term $-c_H H^2|\phi|^2$ is induced by the SUSY breaking due to the finite energy density of the inflaton [4]. ($H \equiv \dot{R}/R$ is the Hubble parameter, R is the scale factor of the expanding universe, and the dot denotes the derivative with cosmic time t.) c_H is a real constant of order unity, which depends on the couplings between the inflaton and the ϕ field in the Kähler potential. Hereafter, we take $c_H \simeq 1$ (> 0), which is crucial to let ϕ have a large expectation value during inflation.

Now let us estimate the baryon asymmetry. The baryon number density is given by, in terms of the AD field,

$$n_B = \frac{1}{3} i(\dot{\phi}^* \phi - \phi^* \dot{\phi}) \,. \tag{4}$$

The equation of motion for the AD field is given by

$$\ddot{\phi} + 3H\dot{\phi} + \frac{\partial V}{\partial \phi^*} = 0 \,. \tag{5}$$

Then, the equation of motion for the baryon number density is written as follows:

$$\dot{n}_B + 3Hn_B = \frac{2}{3}\mathrm{Im}\left(\frac{\partial V}{\partial \phi}\phi\right) = \frac{2}{3}\frac{m_{3/2}}{M^3}\ \mathrm{Im}\left(a_m \phi^6\right) \,, \tag{6}$$

By integrating this equation, we obtain the baryon number at the cosmic time t as

$$\left[R^3 n_B\right](t) = \frac{2}{3}\ \frac{|a_m| m_{3/2}}{M^3} \int^t\ R^3 |\phi|^6 \sin\theta\, dt \,, \tag{7}$$

where $\theta \equiv \arg(a_m) + 6\ \arg(\phi)$.

There are three stages in the evolution of the ϕ field. (i) First, during the inflation, the ϕ filed has a large expectation value due to the negative Hubble mass term at the origin: $|\phi| \simeq (HM^3)^{1/4}$. At this stage, the curvature along the phase direction is much

[4]The thermal effects are negligible in this flat direction as long as the reheating temperature T_R of the inflation is low enough as $T_R \lesssim 10^8$ GeV [2].

[5]The case of anomaly-mediated SUSY breaking [6], where $m_{3/2} \gg m_\phi$, was also discussed in Ref. [2].

smaller than the Hubble parameter, and hence the ϕ field has in general an arbitrary phase. Therefore, we naturally expect that $\sin\theta = \mathcal{O}(1)$.

(ii) After the end of inflation, the AD field slowly rolls down toward the origin following the gradual decrease of the Hubble parameter H as $|\phi(t)| \simeq (H(t)M^3)^{1/4} \propto t^{-1/4}$. At this slow rolling regime, the right-hand side of Eq. (7) increases as $\propto t^{3/2}$. Here, we have assumed the matter-dominated universe $R \propto t^{2/3}$, which is true as long as $T_R \lesssim 2 \times 10^{10}\ \mathrm{GeV}\,(m_\phi/10^3\ \mathrm{GeV})^{1/2}$.

(iii) Finally, the soft mass term of the AD field eventually dominates the negative Hubble mass term at the time when $H(t_{\rm osc}) \simeq m_\phi$, and causes the coherent oscillation of the AD field around the origin. After this time, the amplitude of the AD field rapidly decreases as $|\phi| \propto t^{-1}$, and then the production of the baryon number terminates at the time $H_{\rm osc} \simeq m_\phi$.

Using the above arguments and Eq. (7), we obtain the baryon number density at the time $t = t_{\rm osc}$:

$$n_B(t_{\rm osc}) = \frac{8}{27}\delta_{\rm eff}|a_m|m_{3/2}\left(H_{\rm osc}M^3\right)^{1/2}\,, \tag{8}$$

where $\delta_{\rm eff} \equiv \sin\theta (= \mathcal{O}(1))$. After the completion of the reheating process of the inflation, this leads to the following baryon asymmetry:

$$\frac{n_B}{s} \;=\; \frac{1}{4}\frac{T_R}{M_{pl}^2 H_{\rm osc}^2} n_B(t_{\rm osc})\,, \tag{9}$$

where s is the entropy density, and T_R is the reheating temperature of the inflation. In terms of the density parameter, it is

$$\Omega_B h^2 \;\simeq\; 0.04 \times \delta_{\rm eff}|a_m| \left(\frac{m_{3/2}}{m_\phi}\right)\left(\frac{1\ \mathrm{TeV}}{m_\phi}\right)^{1/2}\left(\frac{M}{M_{\rm pl}}\right)^{3/2}\left(\frac{T_R}{100\ \mathrm{GeV}}\right)\,, \tag{10}$$

where h is the present Hubble parameter in units of 100 km sec^{-1}Mpc^{-1} and $\Omega_B \equiv \rho_B/\rho_c$ (ρ_B and ρ_c are the energy density of the baryon and the critical energy density in the present universe, respectively.) Here, we have used $H_{\rm osc} \simeq m_\phi$. Therefore, the empirical baryon asymmetry $\Omega_B h^2 \simeq 0.02$ can be explained by taking a relatively low reheating temperature $T_R \gtrsim 100$ GeV and a reasonable set of other parameters. However, this is not the whole story.

Formation and Decay of Q-ball Let us consider the epoch just after the generation of the baryon asymmetry finishes, i.e., just after the ϕ field starts its coherent oscillation. If we take into account the one-loop correction, the mass term of the flat direction field becomes

$$V_{\rm mass}(\phi) = m_\phi^2\left[1 + K\log\left(\frac{|\phi|^2}{M_G^2}\right)\right]|\phi|^2\,, \tag{11}$$

where M_G is the renormalization scale at which the soft mass m_ϕ is defined, and the $K\log(|\phi|^2)$ term represents the one-loop correction, which mainly comes from the gaugino

loops. K is estimated in the range from -0.01 to -0.1 [7, 8, 9]. This potential, flatter than ϕ^2, causes spatial instabilities of the homogeneous ϕ's oscillation [10], and the oscillation of the ϕ field fragments into inhomogeneous lumps. These lumps eventually form non-topological solitons, "Q-balls" [11]. We should emphasize here the fact that almost all the baryon asymmetry generated by the AD baryogenesis is absorbed in the Q-balls [12]. Hence, the baryon asymmetry in the present universe must be provided by the decays of these Q-balls.

The decay temperature of the Q-ball is given by [13],

$$T_d \lesssim 2 \text{ GeV} \times \left(\frac{0.03}{-K}\right)^{1/2} \left(\frac{m_\phi}{1 \text{ TeV}}\right)^{1/2} \left(\frac{10^{20}}{Q_i}\right)^{1/2} , \tag{12}$$

where Q_i is the initial charge of the individual Q-ball. In the case of flat directions with the scalar potential (3), it turns out to be [2]

$$Q_i \simeq 10^{20} \times \delta_{\rm eff} |a_m| \left(\frac{m_{3/2}}{m_\phi}\right) \left(\frac{1 \text{ TeV}}{m_\phi}\right)^{3/2} \left(\frac{M}{M_{\rm pl}}\right)^{3/2} . \tag{13}$$

Therefore, the Q-ball decay occurs below $\sim$(a few) GeV. Actually, in most of the flat directions, the charge of the formed Q-ball is as large as $Q_i \sim 10^{19}$–10^{26} [2], and hence the decay temperature T_d of the Q-ball naturally lies in the range of $T_d \sim 1$ MeV–(a few) GeV.

Nonthermal production of LSP Because the Q-ball consists of squarks, its decay produces quarks and supersymmetric particles. The quarks become baryon asymmetry, which is requisite to the big-bang nucleosynthesis, while the supersymmetric particles eventually decay into LSPs. Thus, the number density of the LSPs, $n_{\rm LSP}$, is related to the baryon asymmetry:

$$n_{\rm LSP} \geq n_\phi \geq 3 n_B , \tag{14}$$

where n_ϕ is the number density of the ϕ field (= squarks) in the Q-balls. An important point here is that the Q-ball decay occurs *below* the freeze-out temperature of the LSPs, which is typically given by $T_f \sim m_{\rm LSP}/20$. ($m_{\rm LSP}$ is the mass of the LSP.) Namely, the LSPs are never thermalized after they are produced by the Q-ball decay. If there is no pair annihilation of LSPs afterwards, therefore, the relation in Eq. (14) maintains until present, which results in [8]:

$$\Omega_{\rm LSP} \geq 3 \frac{m_{\rm LSP}}{m_p} \Omega_B = 12 \left(\frac{m_{\rm LSP}}{100 \text{ GeV}}\right) \left(\frac{\Omega_B}{0.04}\right) , \tag{15}$$

where m_p is the nucleon mass. This conflicts with the observed dark matter density $\Omega_{DM} \simeq 0.3$, unless the LSP mass is extremely small:

$$m_{\rm LSP} \leq 3 \text{ GeV} \left(\frac{\Omega_{\rm LSP}}{0.3}\right) \left(\frac{0.04}{\Omega_B}\right) . \tag{16}$$

Consequently, the formation and late time decays of the Q-balls is a serious obstacle for the AD baryogenesis. This is actually the case in the standard parameter regions where the thermal relics of Bino-like LSPs lead to a cosmologically interesting mass density of dark matter.

Here, we consider a simple solution to this problem [1]: an LSP with a large pair-annihilation cross section, like Higgsino- or Wino-like neutralino. If significant pair-annihilations of LSPs occur after they are produced, the relation in Eq. (14) no longer holds. The final abundance of the LSP is obtained by solving the Boltzmann equation analytically, which leads to [1, 2]

$$Y_{\rm LSP}(T) \simeq \left[\frac{1}{Y_{\rm LSP}(T_d)} + \sqrt{\frac{8\pi^2 g_*(T_d)}{45}} \langle \sigma v \rangle M_{\rm pl} (T_d - T) \right]^{-1} , \tag{17}$$

where $Y_{\rm LSP} \equiv n_{\rm LSP}/s$, $g_*(T)$ is the effective degrees of freedom at temperature T, and $\langle \sigma v \rangle$ is the the thermally averaged annihilation cross section of the LSP. If initial abundance $Y_{\rm LSP}(T_d)$ is large enough, the final abundance $Y^0_{\rm LSP}$ for $T \ll T_d$ is given by

$$Y^0_{\rm LSP} \simeq Y^{\rm approx}_{\rm LSP} \equiv \left[\sqrt{\frac{8\pi^2 g_*(T_d)}{45}} \langle \sigma v \rangle M_{\rm pl} T_d \right]^{-1} . \tag{18}$$

Therefore, in this case, the final abundance $Y^0_{\rm LSP}$ is determined only by the Q-ball decay temperature T_d and the annihilation cross section of the LSP $\langle \sigma v \rangle$, independently of the initial value $Y_{\rm LSP}(T_d)$ as long as $Y_{\rm LSP}(T_d) \gg Y^{\rm approx}_{\rm LSP}$. In terms of the density parameter $\Omega_{\rm LSP}$, it is rewritten as

$$\Omega_{\rm LSP} \simeq 0.5 \left(\frac{0.7}{h} \right)^2 \times \left(\frac{m_{\rm LSP}}{100\ {\rm GeV}} \right) \left(\frac{10^{-7}\ {\rm GeV}^2}{\langle \sigma v \rangle} \right) \times \left(\frac{100\ {\rm MeV}}{T_d} \right) \left(\frac{10}{g_*(T_d)} \right)^{1/2} , \tag{19}$$

Therefore, the LSP needs a pair annihilation cross section as large as $\langle \sigma v \rangle \sim 10^{-8}$–$10^{-6}\ {\rm GeV}^{-2}$ so as to obtain the correct mass density as a dominant component of dark matter for the typical decay temperature of Q-balls. Interestingly, Higgsino- and Wino-like LSPs have pair annihilation cross sections just in the desired range, and naturally leads to the correct mass density of dark matter.

Parameter space and possibility of detections We have studied in detail the parameter space where the nonthermal LSP from the Q-ball decay becomes the dominant component of the dark matter, adopting several SUSY breaking models. The Wino-like neutralino dark matter is realized in a wide parameter regions in the anomaly-mediated SUSY breaking models [6] and the no-scale models with nonuniversal gaugino masses [14]. Higgsino dark matter is realized in the minimal supergravity scenario, in the so-called "focus point" region [15]. Here, we show an example of the latter case in Fig. 1. We have also investigated direct and indirect detection of the neutralino dark matter in these regions, and found that there is an intriguing possibility to detect them in various next generation dark matter search experiments [2]. Actually, the possibility of direct and indirect detection of them is much larger than that of the standard Bino-like neutralino.

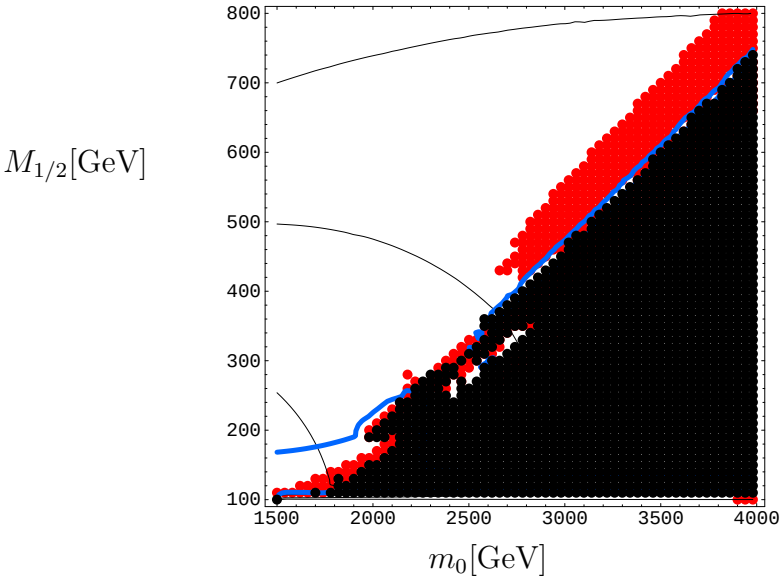

Figure 1: The allowed region in the mSUGRA scenario with $\tan\beta = 15$ and $A_0 = 0$ in the $(m_0$–$M_{1/2})$ plane. In the red shaded region, non-thermally produced LSPs via decays of Q-balls result in $0.05 \leq \Omega_{\rm LSP}h^2 \leq 0.5$ for 1 MeV $\leq T_d \leq$ 10 GeV. The black shaded region is where the electroweak symmetry breaking cannot be implemented. The region below the blue (thick) line is excluded by the chargino mass bound $m_{\chi^\pm} \gtrsim 105$ GeV. The contours of the light Higgs boson mass are given by the black (thin) lines, which correspond to $m_h = 117,\ 120,\ 122$ GeV, respectively. From Ref. [2].

Summary The formation and late time decays of Q-balls are inevitable consequences of the Affleck-Dine baryogenesis via most of the flat directions. In order to avoid the overclosure of the universe by the nonthermal LSPs, which are generated from Q-ball decay, the LSP should have a large pair annihilation cross section. The possible candidates are Higgsino- and Wino-like neutralinos.

We emphasize that, if the Higgsino- or Wino-like neutralino dark matter is indeed detected at future experiments, in the parameter regions we have shown, then it suggests the existence of nonthermal source of these LSPs, since thermal abundance of them would be too small to be the dominant component of the dark matter. The Q-balls produced via the Affleck-Dine baryogenesis is the most promising candidate of such a nonthermal source of the LSPs.

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