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Physics Letters B

www.elsevier.com/locate/physletb

Unified origin of baryons and dark matter

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ARTICLE INFO

Article history:

Received 10 September 2008

Received in revised form 19 September 2008

Accepted 22 September 2008

Available online 27 September 2008

Editor: G.F. Giudice

ABSTRACT

We investigate the possibility that both the baryon asymmetry of the universe and the observed cold dark matter density are generated by decays of a heavy scalar field which dominates the universe before nucleosynthesis. Since baryons and cold dark matter have common origin, this mechanism yields a natural explanation of the similarity of the corresponding energy densities. The cosmological moduli and gravitino problems are avoided.

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1. Introduction

Supersymmetry is a well-motivated framework for TeV scale physics and it has several conceptually nice features. Apart from being the unique extension of the space–time symmetry it provides us with the most compelling scenarios of gauge coupling unification. Not only for short-distance physics, supersymmetry provides interesting possibilities for the explanation of the structure of the universe. The supersymmetric Standard Model contains a natural candidate for cold dark matter of the universe, and also the fact that quadratic divergences are absent allows us to naturally postpone the cut-off of the theory as high as the Planck scale. It makes us possible to calculate high-scale or high-temperature phenomena such as inflation and reheating in a reliable framework.

However, a closer inspection of cosmological scenarios for dark matter production and baryogenesis in supersymmetric models reveals that there is often an inconsistency in the underlying assumptions. For example, the most popular scenarios for dark matter and for baryogenesis are known to be incompatible. It is widely accepted that the relic density of the neutralino from thermal decoupling naturally explains the amount of cold dark matter, and there is a good explanation of the baryon component of the universe by the thermal leptogenesis scenario [1]. The thermal neutralino dark matter is realized in the gravity mediation scenario, i.e., the gravitino mass $m_{3/2} \sim 100$ GeV, whereas the thermal leptogenesis needs a high reheating temperature after inflation, $T_R \gtrsim 10^9$ GeV [2–4]. Considering the thermal production of the gravitinos and their decays, the bound from big bang nucleosynthesis

(BBN) excludes such a possibility (see e.g. [5–9]). It is possible to avoid the constraints from the BBN by assuming a very heavy ($m_{3/2} \gtrsim 100$ TeV) or a stable gravitino, but the most popular explanation of dark matter via the thermal relic neutralinos will be lost in such cases since the lightest supersymmetric particle will be the wino [10,11] (or higgsino [12]), which typically leads to a too small abundance, or the gravitino.

Not only that, for $m_{3/2} \sim 100$ GeV, there is a serious cosmological moduli (or Polonyi) problem [13]. It is possible to give a large mass to the Polonyi field by assuming a non-trivial interactions in order to avoid the late time decay [14], but in such a case the decay into gravitinos causes a disaster [15,16]. There is also a rather disturbing bound on the reheating temperature coming from modular cosmology [17].

It is clear from the above discussion that we need to abandon some of the “standard” assumptions. We propose in this Letter an extreme but cosmologically safe scenario in which radiation, dark matter and baryon asymmetry are all generated right before the BBN era so that there is no constraint from the high-temperature gravitino production, moduli dynamics, etc. (See [18–22] for earlier works.) We find that there is a consistent framework to realize this scenario by a late-time decay of a scalar condensate ϕ in supersymmetric models. We assume that coherent oscillations of a scalar field ϕ first dominate the energy density of the universe. The ϕ field then decays later producing a large amount of radiation and all the dangerous relics (e.g. gravitinos) get diluted to a negligible density. The standard BBN takes place after ϕ has decayed at sufficiently high temperature. The dark matter is produced by the ϕ decay and the correct abundance is obtained in the wino/higgsino LSP scenario due to pair annihilation [23] which is less effective than in the so-called thermal scenarios. The baryon asymmetry is also generated in ϕ decays. The decay channel of a ϕ

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field into two gravitinos is forbidden by R -parity, and we assume that the decay into a gravitino and the fermionic superpartner of ϕ ($\tilde{\phi}$) is kinematically forbidden. Interestingly, by unifying the origin of dark matter and baryon asymmetry, this scenario explains the one of the puzzling issues of our universe that the energy densities of dark matter and baryon asymmetry are close to each other, $\Omega_b \sim \Omega_{\text{CDM}}$.¹

The Letter is organized as follows: in Section 2 we present the baryogenesis mechanism, and show that it is indeed possible to generate the observed amount of baryon asymmetry while satisfying the BBN constraints on the reheating temperature after the ϕ decay. In Section 3, we discuss the abundance of the dark matter from the ϕ decay. In particular, in our preferred scenario the ratio $\Omega_{\text{CDM}}/\Omega_b \sim 5$ implies a large gravitino mass, $m_{3/2} \sim 100$ TeV which fits nicely to the wino/higgsino LSP scenario.

2. Baryogenesis

2.1. Basic idea

We consider a chiral superfield $\Phi = (\phi, \tilde{\phi}, F_\phi)$ which couples to the matter fields via a higher-dimensional term in the superpotential [20],

$$\mathcal{W} \supset \frac{1}{M} \Phi U D D, \quad (1)$$

with $U = (\tilde{u}^c, u^c, F_u)$ and $D = (\tilde{d}^c, d^c, F_d)$ denoting the up- and down-type quark superfields. Here, we suppressed color and generation indices, and absorbed dimensionless couplings into M which will be taken to be of the order of the Planck scale, $M \sim M_P = 2.44 \times 10^{18}$ GeV, unless stated otherwise. Due to this operator ϕ effectively carries baryon number (+1).

Let us now define the ϕ number asymmetry

$$q_\phi := i(\dot{\phi}^* \phi - \phi^* \dot{\phi}). \quad (2)$$

q_ϕ is given by the difference between the number densities n_ϕ and n_{ϕ^*} of particles ϕ and antiparticles ϕ^* . q_ϕ can be interpreted as angular momentum of the ϕ field rotating in the complex plane [29].

The scenario we shall describe in the following consist of the following sequence of steps: first, a positive q_ϕ is generated. Then there is an era of coherent ϕ oscillations where a significant fraction of the energy density of the universe is carried by these oscillations, and q_ϕ is conserved. Finally, ϕ number is converted into the baryon number by its decay, and it reheats the universe up to $O(100$ MeV) consistently with nucleosynthesis.

Before we describe the mechanism in detail, let us briefly explain the main differences to Ref. [20]. For the mechanism to work, one has to make sure that dangerous ϕ -number violating interaction terms are absent or sufficiently suppressed. One can forbid these terms by imposing a symmetry. In Ref. [20], a model with a \mathbb{Z}_{4R} symmetry is presented, which ensures the ϕ -number conservation at a sufficient level and also prevents ϕ from dominating the universe. Below, we will consider a different model with an anomaly-free \mathbb{Z}_9 discrete baryon symmetry in addition to the usual R -parity (see Table 1), and will, as already stated, assume that ϕ dominates the universe at an early epoch.

In order to obtain the baryon asymmetry before BBN, in Ref. [20] enhanced couplings of ϕ to the baryons are assumed such that the ϕ lifetime is short enough. Moreover, in the case of the ϕ mass of the order 100 GeV, the universe would be always matter (ϕ or dark matter) dominated once ϕ dominates the universe at

Table 1

Charge assignments under \mathbb{Z}_9 and R -parity. We denote the MSSM superfields according to $Q = (q, \tilde{q}, F_q)$ with q representing the quark doublets, etc.

Field	Q	U	D	L	E	N	H	\bar{H}	ϕ	$\tilde{\phi}$
\mathbb{Z}_9	+1	-1	-1	0	0	0	0	0	+3	-3
R -parity	-	-	-	-	-	-	+	+	-	-

an early time. This is not compatible with the requirement of successful BBN and, therefore, a model without (early) ϕ domination is constructed there.

The situation is, however, different if ϕ is heavy. As the temperature after the ϕ decay is higher than the BBN temperature, ϕ is allowed to dominate the energy density of the universe. This high temperature after ϕ decay also plays a crucial role in the generation of cold dark matter (cf. Section 3).

2.2. ϕ evolution

Let us start by considering the dynamics of the ϕ field. The evolution of ϕ is described by its equation of motion,

$$\ddot{\phi} + (3H + \Gamma_\phi)\dot{\phi} + \frac{\partial V}{\partial \phi^*} = 0, \quad (3)$$

where $V = V(\phi, \phi^*, \dots)$ denotes the scalar potential, H the Hubble rate and Γ_ϕ the ϕ decay rate. Eq. (3) translates into an equation of motion for q_ϕ ,

$$\dot{q}_\phi + 3Hq_\phi = -i\left(\phi \frac{\partial V}{\partial \phi} - \phi^* \frac{\partial V}{\partial \phi^*}\right). \quad (4)$$

Hence, a non-vanishing right-hand side of (4) can be used for the first step, i.e. to create non-zero q_ϕ dynamically. Before explaining this in detail, recall that we need also to satisfy the condition of ϕ number conservation in the stage of ϕ oscillation. This means that in the ϕ oscillation era the ϕ number violating terms have to be absent (or sufficiently suppressed). The most dangerous term of this type is $\mu^2 \phi^2 + \text{h.c.}$

In order to enforce the absence of those dangerous terms, we impose a discrete \mathbb{Z}_9 symmetry which is an anomaly free subgroup of baryon number symmetry [30–34]. The charge assignment is listed in Table 1, where $\tilde{\phi}$ is introduced in order to give a mass term for the fermionic superpartner of ϕ without introducing $\mu^2 \phi^2$ term in the Lagrangian.² With this choice $\phi = 0$ can always be a minimum of the potential which is necessary to preserve the R -parity.

The symmetry does, however, allow for a Φ^6 term in the Kähler potential and in the superpotential. In the following we discuss the two cases in which the Φ^6 term in superpotential is absent and present.

Case A: No Φ^6 term in the superpotential

In the case where there is a Φ^6 term only in the Kähler potential, the potential for the ϕ field is given by

$$V = m_\phi^2 |\phi|^2 + m_{3/2}^2 M^2 F(|\phi|^2/M^2) + \left[\kappa \frac{m_{3/2}^2}{M^4} \phi^6 + \text{h.c.} \right] + \text{higher-order terms}, \quad (5)$$

where κ is expected to be order one. The gravitino mass parameter $m_{3/2}$ represents the supersymmetry breaking scale. The function $F(x)$ is a general (polynomial) function.

² The analysis will remain unchanged when we include the dynamics of the $\tilde{\phi}$ field. Although a possible mass mixing term, $m^2 \phi \tilde{\phi} + \text{h.c.}$, will distribute the baryon number to the $\tilde{\phi}$ field, Eq. (4) will be the same once we include the $\tilde{\phi}$ field in q_ϕ in Eq. (2). The decay of $\tilde{\phi}$ (or more precisely the other mass eigenstate) happens about the same time as the ϕ decay provided the mixing is order one.

¹ See, for example, [21,24–28] for earlier attempts to explain the similarity: $\Omega_b \sim \Omega_{\text{CDM}}$.

The presence of the ϕ^6 term in (5) can lead to a dynamical generation of $q_\phi \neq 0$ as follows: for $H \gg m_\phi$, ϕ oscillation is negligible and we can treat ϕ to be constant. We can integrate the equation

$$\dot{q}_\phi + 3Hq_\phi = \frac{m_{3/2}^2}{M^4} \text{Im}[\kappa\phi^6], \quad (6)$$

so that

$$q_\phi(t = m_\phi^{-1}) \sim |\kappa| \frac{m_{3/2}^2}{2m_\phi M^4} \phi_{\text{ini}}^6, \quad (7)$$

where ϕ_{ini} is the initial amplitude of ϕ after inflation which is generically $O(M)$ with the potential in Eq. (5) (with $m_{3/2}^2$ replaced by $O(H^2)$). The number density of ϕ and ϕ^* particles is given by ρ_ϕ/m_ϕ where $\rho_\phi \simeq m_\phi^2 |\dot{\phi}|^2 + |\dot{\phi}|^2$. Since the oscillation of ϕ starts with amplitude ϕ_{ini} , we obtain for the dimensionless ϕ asymmetry

$$\varepsilon := \frac{q_\phi}{n_\phi + n_{\phi^*}} \sim |\kappa| \left(\frac{m_{3/2}}{m_\phi} \right)^2. \quad (8)$$

Here we have taken $\phi_{\text{ini}} \sim M$. We have checked that this expression yields roughly the correct order of magnitude for ε (as long as ϕ_{ini} is comparable to M) by solving the equation of motion numerically.

The resulting ϕ asymmetry stays constant after ϕ starts to oscillate because of the r.h.s. of Eq. (6) becomes numerically irrelevant when the amplitude drops (far) below M . As a consequence, $R^3 q_\phi$ (with R being the scale factor) is approximately conserved during the ϕ oscillation era, until ϕ decays.

Case B: Φ^6 term in the superpotential

The case with Φ^6 term in the superpotential is qualitatively the same as case A. The potential of the ϕ field in this case is given by

$$V = m_\phi^2 |\phi|^2 + m_{3/2}^2 M^2 F(|\phi|^2/M^2) + \left[\kappa' \frac{m_{3/2}}{M^3} \phi^6 + \text{h.c.} \right] + \kappa'' \frac{|\phi|^{10}}{M^6} + \text{higher-order terms}, \quad (9)$$

where κ' and κ'' are $O(1)$ coefficients. In the early universe, $m_{3/2}$ which represents the SUSY breaking effect is replaced by the Hubble rate H . Now the minimum of the potential is generically $\phi \sim (HM^3)^{1/4}$ which sets the initial amplitude of the ϕ oscillation to be $\phi_{\text{ini}} \sim (m_\phi M^3)^{1/4}$ since the ϕ oscillations start when $H \sim m_\phi$.

The equation for the evolution of the ϕ number asymmetry q_ϕ in Eq. (4) is

$$\dot{q}_\phi + 3Hq_\phi = \frac{m_{3/2}^2}{M^3} \text{Im}[\kappa'\phi^6]. \quad (10)$$

With the value of ϕ at the minimum of the potential, $\phi \sim (HM^3)^{1/4}$, before ϕ oscillation, we find

$$q_\phi \sim |\kappa'| \frac{m_{3/2}}{m_\phi} \frac{\phi_{\text{ini}}^6}{M^3}. \quad (11)$$

Therefore, the asymmetry factor ε is estimated to be

$$\varepsilon \sim |\kappa'| \frac{m_{3/2}}{m_\phi}, \quad (12)$$

which is larger than the case without the Φ^6 term in the superpotential if $m_{3/2} \ll m_\phi$ and $\kappa' \sim 1$. In conclusion, the presence of the Φ^6 term leads only to a quantitatively different result.

2.3. Baryogenesis via ϕ decay

So far we have seen that, due to the presence of higher-order terms, a ϕ number asymmetry is induced which is conserved in

the regime of ϕ oscillations until ϕ decays. Let us now consider the conversion of ϕ number to baryon number of the universe through the decay arising from the coupling (1), $\phi \rightarrow qq\bar{q}$. The corresponding decay rate is given by

$$\Gamma_\phi = \xi \frac{m_\phi^3}{M^2}, \quad (13)$$

where ξ is obtained by a standard calculation. In the simplest case where all couplings of ϕ to the quark superfields equal one, we obtain $\xi = 27/(256\pi^3) \simeq 3 \times 10^{-3}$.

The temperature T_d of the thermal bath after ϕ decay is calculated by equating Hubble rate H and Γ_ϕ ,

$$T_d \simeq 120 \text{ MeV} \left(\frac{\xi}{10^{-2}} \right)^{1/2} \left(\frac{m_\phi}{1500 \text{ TeV}} \right)^{3/2} \left(\frac{M_{\text{P}}}{M} \right). \quad (14)$$

Since ϕ has a large hadronic branching fraction, T_d has to fulfill [35] $T_d \gtrsim 4 \text{ MeV}$. The corresponding lower bound on the ϕ mass is

$$m_\phi \gtrsim 150 \text{ TeV} \left(\frac{10^{-2}}{\xi} \right)^{1/3} \left(\frac{M}{M_{\text{P}}} \right)^{2/3}, \quad (15)$$

where we have exploited that $\sqrt{\frac{\pi^2 g_*}{90}} \simeq 1$ for temperatures of the order MeV.

Assuming that ϕ dominates the energy density of the universe before its decay, the number density of ϕ just before decay is obtained by

$$m_\phi(n_\phi + n_{\phi^*}) \simeq \frac{\pi^2}{30} g_* T_d^4. \quad (16)$$

Using the relation between baryon asymmetry and ϕ number, $n_b = q_\phi = \varepsilon(n_\phi + n_{\phi^*})$, we can estimate the baryon asymmetry as

$$\begin{aligned} \frac{n_b}{s} &\simeq \frac{3}{4} \varepsilon \frac{T_d}{m_\phi} \\ &\sim \begin{cases} 10^{-10} \cdot |\kappa| \left(\frac{\xi}{10^{-2}} \right)^{1/2} \left(\frac{m_{3/2}}{50 \text{ TeV}} \right)^2 \left(\frac{m_\phi}{1500 \text{ TeV}} \right)^{-3/2} \left(\frac{M_{\text{P}}}{M} \right), & \text{(A)} \\ 10^{-10} \cdot |\kappa'| \left(\frac{\xi}{10^{-2}} \right)^{1/2} \left(\frac{m_{3/2}}{2 \text{ TeV}} \right) \left(\frac{m_\phi}{1500 \text{ TeV}} \right)^{-1/2} \left(\frac{M_{\text{P}}}{M} \right). & \text{(B)} \end{cases} \end{aligned} \quad (17)$$

Cases A and B correspond to the model without and with the Φ^6 term in the superpotential, respectively. By comparing the observed value $(n_b/s)^{\text{obs}} = (8.7 \pm 0.3) \times 10^{-11}$ [36], the lower bound on m_ϕ in Eq. (15) requires either enhanced SUSY breaking terms $|\kappa| \gg 1$ or the large gravitino mass $m_{3/2} \gtrsim 10 \text{ TeV}$ for case A whereas $m_{3/2} \sim 1 \text{ TeV}$ is possible in case B (assuming $M \simeq M_{\text{P}}$).

As we shall see in the next section, the reference value of $m_\phi \sim 1500 \text{ TeV}$ is motivated by considerations on dark matter. Notice that the required size of the SUSY breaking scale $m_{3/2}$ from baryogenesis is indeed at a favorable value for low energy phenomenology. In particular, in case A, $m_{3/2} \sim 50 \text{ TeV}$ indicates that loop corrections to soft masses (anomaly mediation) are important. In many models with such situation, the wino [10,11] or higgsino [37–39] becomes the LSP and that is indeed consistent with the discussion below. For case B, the wino or higgsino LSP is realized by either assuming small value or phase of κ' such that gravitino mass is enhanced or simply assuming the wino or higgsino LSP in the scenarios of the gravity mediation type by relaxing the universality of the gaugino and/or scalar masses.

3. Dark matter from ϕ decay

Since every ϕ decay produces (at least) one superpartner, Ref. [20] concludes that the number density of LSPs exceeds the one of baryons, $n_{\text{LSP}} \gtrsim n_b$. However, n_{LSP} is modified by LSP pair annihilation processes in a heavy ϕ scenario. These processes are effective as long as the corresponding rate exceeds the Hubble rate.

In the MSSM, the dark matter candidates which pair annihilate strongly are the wino and the higgsino. Both particles have large annihilation cross sections through weak interaction, and thus the thermal abundance cannot explain the energy density of the dark matter. On the other hand, non-thermal production from ϕ decays renders the wino/higgsino a viable dark matter candidate.

In order to find out to what extent the LSPs annihilate, one describes the evolution of number densities of ϕ quanta and LSPs, n_ϕ and n_χ , and energy density of the thermal bath, ρ_{rad} , by Boltzmann equations [23],

$$\frac{dn_\phi}{dt} + 3Hn_\phi = -\Gamma_\phi n_\phi, \quad (18a)$$

$$\frac{dn_\chi}{dt} + 3Hn_\chi = v_{\text{LSP}}\Gamma_\phi n_\phi - \langle\sigma v\rangle n_\chi^2, \quad (18b)$$

$$\frac{d\rho_{\text{rad}}}{dt} + 4H\rho_{\text{rad}} = (m_\phi - v_{\text{LSP}}m_\chi)\Gamma_\phi n_\phi + m_\chi\langle\sigma v\rangle n_\chi^2. \quad (18c)$$

v_{LSP} denotes the number of LSPs produced by a ϕ decay. The Boltzmann equations can be integrated, and the relic density of χ can be approximated by [40]

$$\frac{n_\chi}{s} \sim (4\langle\sigma v\rangle M_P T_d)^{-1}, \quad (19)$$

as long as the LSPs are not equilibrated, i.e. $T_d \lesssim m_\chi/30$. The relic χ abundance is then

$$\begin{aligned} \Omega_\chi h^2 \simeq 0.1 & \left(\frac{2.5 \times 10^{-3}}{m_\chi^2 \langle\sigma v\rangle} \right) \left(\frac{10^{-2}}{\xi} \right)^{1/2} \left(\frac{m_\chi}{100 \text{ GeV}} \right)^3 \\ & \times \left(\frac{m_\phi}{1500 \text{ TeV}} \right)^{-3/2} \left(\frac{M}{M_P} \right). \end{aligned} \quad (20)$$

The thermal average of the annihilation cross section is typically $\langle\sigma v\rangle \sim 10^{-3}/m_\chi^2$ for the particles which have $SU(2)_L$ quantum numbers such as wino and higgsino. Therefore the non-thermal component can explain the dark matter of the universe for $m_\phi \sim 10^3\text{--}5$ TeV depending on m_χ .

For concreteness, let us focus on the case of the wino LSP. The annihilation cross section is [23] (cf. the extensive list [41])

$$\langle\sigma_{\tilde{W}^0\tilde{W}^0 \rightarrow W^+W^-} v\rangle = \frac{g_2^4}{2\pi} \frac{1}{m_\chi^2} \frac{[1 - \frac{m_W^2}{m_\chi^2}]^{3/2}}{[2 - \frac{m_W^2}{m_\chi^2}]^2}. \quad (21)$$

In Fig. 1 we show the relic wino density $\Omega_\chi h^2$ (where $h \simeq 0.7$ is the present normalized Hubble expansion rate [36]) as a function of m_ϕ . To produce Fig. 1, we solve the set of Boltzmann equations (18) (extended to include the charged wino NLSP) and take into account coannihilation. The coannihilation effect between LSP and NLSP becomes important for large m_ϕ , and explains the deviation of the contours from straight lines (see in particular the $m_\chi = 100$ GeV contour in Fig. 1). If $T_d \gtrsim m_\chi/30$, the winos are in thermal equilibrium, and the relic abundance does not depend on m_ϕ any more. This explains why the contours become horizontal for large m_ϕ in Fig. 1.³ For $T_d \lesssim m_\chi/30$, the estimate (20) turns out to be a reasonable approximation. In particular, the temperature T_d can be as low as 100 MeV without overclosing the universe. If we fix M to be the Planck scale, we find that m_ϕ should exceed 10^3 TeV.

It is instructive to study the dependence of the baryon density in Eq. (17) and the relic dark matter density in Eq. (19) on the

³ For such large m_ϕ , the standard WIMP scenario works and thus the bino is a good dark matter candidate. However, this would require even larger ϕ masses, and according to Eq. (17), large m_ϕ would make baryogenesis difficult.

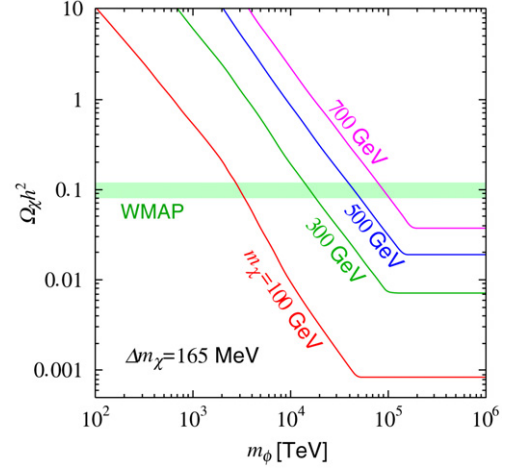


Fig. 1. $\Omega_\chi h^2$ as a function of the mass m_ϕ for various wino masses m_χ . We use for the mass difference between charged and neutral winos $\Delta m_\chi = 165$ MeV, and fix $M = M_P$. The shaded bar corresponds to the 2σ region of Ω_{CDM} as reported by WMAP [36]. As usual, we fix the present energy density of the universe so that Ω can formally become larger than one ('overclosure').

physical parameters. For case A in Eq. (17), amazingly, the ratio is independent of the ϕ mass,

$$\begin{aligned} \frac{n_\chi}{n_b} & \sim |\kappa|^{-1} (4\langle\sigma v\rangle \xi m_{3/2}^2)^{-1} \left(\frac{M}{M_P} \right)^2 \\ & \sim 10^4 |\kappa|^{-1} \left(\frac{m_\chi}{M_{3/2}} \right)^2 \left(\frac{M}{M_P} \right)^2. \end{aligned} \quad (22)$$

Our scenario relies on a large gravitino mass which is realized in anomaly mediation [10,11] where the wino mass is suppressed by a loop factor, e.g. $m_\chi/m_{3/2} \sim g_2^2/(16\pi^2)$ for the wino χ . This implies $n_\chi/n_b \sim (\text{few}) \times 10^{-2}$ for $M \sim M_P$. Hence, our scenario predicts for the ratio of dark matter to baryon densities

$$\frac{\Omega_\chi}{\Omega_b} \sim |\kappa|^{-1} \times \text{few} \times 10^{-2} \times \frac{m_\chi}{m_{\text{nucleon}}} \times \left(\frac{M}{M_P} \right)^2. \quad (23)$$

In particular, for m_χ of the order 100 GeV (and $M \simeq M_P$), the observed ratio $\Omega_{\text{CDM}}/\Omega_b \simeq 5$ [36] finds a very natural explanation within the framework described here. The same is true for the higgsino LSP case since it naturally has a mass of the order of the wino mass.

It is interesting to relax the assumption $M \simeq M_P$ and take, for instance, M to be of order GUT or compactification scale, $M \sim M_{\text{GUT}} \simeq 3 \times 10^{16}$ GeV, or the string scale. If so, the ϕ mass can be substantially lower, $m_\phi \sim m_{3/2} \sim 100$ TeV.

Concerning cold dark matter, our analysis coincides with the one of Ref. [23] if we identify ϕ as a modulus. However, our assumption that ϕ is odd under R -parity completely avoids the gravitino problem caused by the ϕ decay [42,43].

4. Conclusions

We have discussed a scenario where both the observed baryon asymmetry and the cold dark matter originate from decays of a heavy scalar field ϕ . In an example we imposed \mathbb{Z}_9 symmetry such that ϕ number is effectively conserved in the ϕ oscillation era. This allows to have an initial asymmetry q_ϕ which is conserved until the ϕ decay, and q_ϕ is converted into baryon asymmetry. The baryon asymmetry is automatically in the right ballpark for ϕ masses which are high enough to evade the BBN constraints. The ϕ decays also produce LSPs. For sufficiently high decay temperature T_d , pair annihilation is still partially effective if the LSP is the wino or the higgsino, and hence the number density of

LSPs gets reduced to an appropriate amount of the dark matter, $n_\chi \ll n_b$. In one of our scenarios (case A) we require a heavy gravitino, $m_{3/2} \sim 100$ TeV, anomaly mediated contributions to gaugino masses become important, so that one naturally obtains a wino or higgsino LSP.

Amazingly, in that context, our mechanism predicts $\Omega_\chi/\Omega_b \sim \text{few} \times 10^{-2} \times (m_\chi/m_{\text{nucleon}})$ independently of the ϕ mass m_ϕ . It can hence very naturally account for $\Omega_\chi \simeq 5\Omega_b$ although each ϕ decay produces (at least) one superpartner and $m_{\text{LSP}} \gtrsim 100 \times m_{\text{nucleon}}$. In our scenario, the role of the ϕ field is thus twofold: it serves both as source of the observed baryon asymmetry and explains why winos or higgsinos can be cold dark matter despite of the large annihilation cross section.

Our mechanism also opens new possibilities in inflation model building. One could, for instance, envisage a sequestered scenario where the inflaton is (geometrically) separated from the MSSM fields. This is usually a problem because then the inflaton reheats the hidden sector. In our scenario, however, such a setup may be viable because (dark) matter and radiation are generated by the ϕ field rather than inflaton decays.

Acknowledgements

We would like to thank K. Choi, K. Hamaguchi, O. Lebedev and H.P. Nilles for useful discussions. We acknowledge support from the Aspen Center for Physics where this work has been started. The research of M.R. is supported by the DFG cluster of excellence Origin and Structure of the Universe, the Graduiertenkolleg “Particle Physics at the Energy Frontier of New Phenomena” and the SFB-Transregio 27 “Neutrinos and Beyond” by the Deutsche Forschungsgemeinschaft (DFG). The research of H.M. is supported in part by World Premier International Research Center Initiative (WPI Initiative), MEXT, Japan, in part by the US Department of Energy under Contract DE-AC03-76SF00098, and in part by the National Science Foundation under grant PHY-04-57315.

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