Axino Dark Matter from Q-Balls in Affleck-Dine Baryogenesis and the $\Omega_b - \Omega_{\rm DM}$ Coincidence Problem

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We show that the $\Omega_b - \Omega_{\rm DM}$ coincidence can naturally be explained in a framework where axino is cold dark matter which is predominantly produced in nonthermal processes involving decays of Q-balls formed in Affleck-Dine baryogenesis. In this approach, the similarity of Ω_b and $\Omega_{\rm DM}$ is a direct consequence of the (sub-)GeV scale of the mass of the axino, while the reheating temperature T_R must be low, some 10^2 GeV, or less.

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1. Introduction.—The origin of nonbaryonic cold dark matter (DM) and of baryon asymmetry in the Universe are among the longest lasting puzzles in cosmology as well as in particle physics today. In particular, the question of why the observed values of baryon density Ω_b and of dark matter $\Omega_{\rm DM}$ are so close to each other, $\Omega_{\rm DM}/\Omega_b = 5.65 \pm 0.58$ [1], remains a mystery.

A standard paradigm is that the nonbaryonic cold dark matter is made up of some weakly interacting massive particle (WIMP) which freezes out of thermal equilibrium in the early Universe. Perhaps the most popular WIMP candidate is the lightest neutralino of the Minimal Supersymmetric Standard Model (MSSM) as the lightest supersymmetric particle (LSP). It remains stable due to the conservation of R-parity. This economical scenario does not, however, explain the proximity of Ω_b and $\Omega_{\rm DM}$.

The same is generally true for conventional mechanisms of baryogenesis or leptogenesis. This may indicate that the observed baryon-to-DM density ratio is just a pure accident, or else a result of some underlying, and as yet unknown, more fundamental theory. An alternative approach is to try to identify a physical mechanism which would simultaneously produce both baryon asymmetry and DM in the proportions consistent with observations. It is clear that this basically necessitates abandoning standard paradigms for producing both types of species in the Universe. This may be one important lesson to learn from these considerations.

A number of attempts at explaining baryon-to-DM ratio have been suggested in the literature. For instance, recently a right handed sneutrino [2] and a sneutrino condensate as an AD field [3] have been proposed.

A few years ago, Enqvist and McDonald (EMD) proposed [4,5] an attractive solution based on a variant of Affleck-Dine (AD) baryogenesis [6]. In that scenario, an AD condensate forms during inflation and develops a large vacuum expectation value (VEV) along a D-flat direction in the MSSM. D-flat directions are configurations of scalar fields for which the D-part of the potential vanishes. Prevalent in theories with many scalar fields like the MSSM, they are of much interest to cosmology [7].

In the standard AD scenario, after the end of inflation, the scalar field condensate slowly rolls towards the origin and, after a few dozens of coherent oscillations, produces a nonzero baryon number in presence of baryon number-violating couplings of the fields making up the flat direction. Originally, Kusenko and Shaposhnikov argued that the AD condensate can instead fragment into nontopological solitons called Q-balls [8]. If their baryonic charge is large enough, as in models with gauge mediated SUSY breaking, Q-balls remain effectively stable until today, and contribute to the DM density, despite severe astrophysical constraints [9]. On the other hand, EMD demonstrated that, under nontrivial but natural conditions (that we summarize below), in a large class of supergravity (SUGRA) models with gravity mediated SUSY breaking (GRMSB) Q-balls subsequently decay into baryonic matter and neutralino WIMPs assumed to be the LSP [4].

In the EMD scenario, the baryon-to-DM ratio can easily be estimated to be in the right ballpark, as we shall see below. This otherwise attractive framework suffers, however, from a serious problem: neutralino production in Q-ball decays is in fact too efficient, and density Ω_χ can only agree with observations for low neutralino mass $m_\chi \sim 1$ GeV, well below LEP limits [10]. Moreover, this puts into a potential jeopardy the AD mechanism in a large class of GRMSB supergravity models.

In this Letter, we suggest a way out from the above problems of the EMD scenario which at the same time preserves its successful features, in particular, an explanation of the $\Omega_b/\Omega_{\rm DM}$ ratio. We propose that the DM is not made up of the neutralino but instead of an axino, a superpartner of the axion. The axino is a neutral Majorana, chiral fermion. It arises in SUSY models incorporating a Peccei-Quinn solution to the strong CP problem in QCD. Unlike for the neutralino or gravitino, its mass is strongly model dependent and can be much smaller than the (gravity mediated) SUSY breaking scale [11–13]. Similarly to the axion, its interactions are suppressed by the PQ scale

$f_a \simeq 10^{11}$ GeV, well below the sensitivity of LEP. The axino has a number of properties which make it a promising candidate for cold dark matter [14,15]. Earlier papers considered warm axino relics [13,16]. As we will show below, axinos are naturally produced at low temperatures of a few GeV, consistent with the Q-ball scenario of EMD but still before the period of Big Bang Nucleosynthesis (BBN).

2. The Enqvist-McDonald scenario.—We now briefly present the main features of the EMD variant of the AD baryogenesis. It is assumed that the AD field ϕ is given by a D-flat direction in the MSSM. D-flat directions are configurations of scalar fields for which the D-part of the potential vanishes. Its potential is, in general, lifted by soft supersymmetric (SUSY) breaking terms and nonrenormalizable terms [17,18].

The potential of the AD field, including inflaton-induced terms, reads

$$V(\phi) \simeq \left\{(m_\phi^2 - c_1 H^2)\left[1 + K \ln\left(\frac{|\phi|^2}{\Lambda^2}\right)\right]\right\}|\phi|^2 + \left[(c_2 H + A m_{3/2})\frac{\lambda \phi^n}{n M^{n-3}} + \text{H.c.}\right] + \lambda^2 \frac{|\phi|^{2n-2}}{M^{2n-6}}, \tag{1}$$

where m_ϕ is the soft SUSY breaking mass for the AD field and a radiative correction is given by $K \ln|\phi|^2$. A flat direction dependent constant, K, takes values from -0.01 to -0.1 [19,20]. Λ denotes a renormalization scale, and $-c_1 H^2$, with $c_1 \sim 1$, is the negative mass-squared term induced by the energy density of the inflaton [18]. Terms proportional to A and c_2 are the trilinear terms from low energy SUSY breaking and those induced by the inflaton, respectively, while $m_{3/2}$ denotes the gravitino mass. The nonrenormalizable terms in Eq. (1) come from the superpotential $W = \lambda/nM^{n-3}\phi^n$, where λ is the Yukawa coupling and M is some large scale acting as a cut-off. In SUGRA, it is natural to assume $M = M_P \simeq 2.4 \times 10^{18}$ GeV which is the reduced Planck mass.

Since during inflation the Hubble parameter $H \gg m_\phi \sim m_{3/2}$, the AD field settles down at the minimum of the potential (1) which is given by

$$|\phi| \simeq \left(\sqrt{\frac{c_1}{n-1}}\frac{H M_P^{n-3}}{\lambda}\right)^{1/(n-2)} \simeq \left(\frac{H M_P^{n-3}}{\lambda}\right)^{1/(n-2)}. \tag{2}$$

It is clear that the AD field can naturally develop a very large VEV, which is possible in nonminimal Kähler potentials [18], or if large enough trilinear term A is induced by the inflaton [21].

We have neglected in Eq. (1) thermal mass terms $h^2 T^2 |\phi|^2$, where h denotes couplings of the AD field to other particles [22]. They would play a role if the AD field VEV were relatively small. We have also neglected two loop thermal effects due to the running of gauge coupling which generate a term $\alpha T^4 \ln(|\phi|^2/T^2)$, where $|\alpha| = \mathcal{O}(10^{-2})$ [23]. They will not be important below.

As H decreases, the AD field traces the instantaneous minimum after inflation, begins to oscillate when $H_{\rm osc}^2 \simeq m_\phi^2$ and, after a few dozen turns, produces a nonzero baryon number and then fragments into Q-balls.

The baryon number density for the AD field ϕ is given by $n_b = iq(\dot\phi^* \phi - \phi^* \dot\phi)$ where q is the baryonic charge for the AD field. By using the equation of motion of the AD field, the charge density can be rewritten as

$$n_b(t) \simeq \frac{1}{a(t)^3}\int^t dt' a(t')^3 \frac{2 q m_{3/2}}{M_P^{n-3}} \text{Im}(A\phi^n), \tag{3}$$

with $a(t)$ being the scale factor. When the AD field starts to oscillate around the origin, the baryon number density is induced by the relative phase between A and c_2. With the entropy density after reheat $s = 4\pi^2 g_* T^3/90$, we can express the baryon asymmetry as

$$\frac{n_b}{s} = \frac{T_R n_b}{4 M_P^2 H^2}\bigg|_{t_{\rm osc}} \simeq \frac{q|A| m_{3/2}}{2}\frac{T_R |\phi_{\rm osc}|^n}{H_{\rm osc}^3 M_P^{n-1}} \sin\delta. \tag{4}$$

Here, $t_{\rm osc}$ denotes the time of the start of the oscillation, and $\sin\delta$ is the effective CP phase.

From now on, we consider the case of $n = 6$ because a promising AD field for our scenario, a $\bar u \bar d \bar d$ direction belongs to this class. Let us first evaluate n_b/s. The baryon asymmetry, Eq. (4), for the relevant case is estimated as

$$\frac{n_b}{s} \simeq \frac{q|A|\sin\delta}{2\lambda^{3/2}}\frac{m_{3/2} T_R}{m_\phi^{3/2} M_P^{1/2}} \simeq 1 \times 10^{-10}\frac{q|A|\sin\delta}{\lambda^{3/2}}\left(\frac{m_{3/2}}{100\ \text{GeV}}\right)\left(\frac{10^3\ \text{GeV}}{m_\phi}\right)^{3/2} \times \left(\frac{T_R}{100\ \text{GeV}}\right) \tag{5}$$

which is of the right order. (The previously made assumption $M \sim M_P$ is crucial, for otherwise Q-balls would decay too early or would evaporate.) The low reheat temperature T_R after inflation is required to explain the appropriate baryon asymmetry. In this case, the AD condensate fragments into Q-balls [19,20].

The growth of perturbations of the AD field and its subsequent fragmentation into Q-balls crucially depends on the logarithmic correction to the ϕ^2 mass term in $V(\phi)$, Eq. (1). An essential requirement is that $V(\phi)$ is flatter than quadratic, or that $K < 0$ [20]. This can be achieved in SUGRA models with a nonminimal Kähler potential [19].

In order to discuss the evolution of Q-balls, first we briefly summarize their relevant properties in GRMSB models. The radius of a Q-ball, R, is estimated as $R^2 \simeq 2/(|K| m_\phi^2)$ [19]. The charge is roughly given by $Q \simeq \frac{4}{3}\pi R^3 n_b(t_i) \simeq \frac{4}{3}\pi R^3 (H_i/H_{\rm osc})^2 n_b|_{t_{\rm osc}}$, where the suffix i represents the time when the spatial imhomogeneity becomes nonlinear, which can be evaluated as [24]

$$Q \sim 6 \times 10^{-3} \frac{2q|A|\sin\delta}{\lambda^{3/2}} \frac{m_{3/2} M_P^{3/2}}{m_\phi^{5/2}}$$
$$\simeq 1 \times 10^{20} \frac{q|A|\sin\delta}{\lambda^{3/2}} \left(\frac{m_{3/2}}{100\ \mathrm{GeV}}\right)\left(\frac{1\ \mathrm{TeV}}{m_\phi}\right)^{5/2}. \quad (6)$$

Unless $Q > \mathcal{O}(10^{18})$, Q-balls will evaporate before decaying [25]. For Q as in Eq. (6), Q-ball decay temperature is $T_d \simeq 1$ GeV to 1 MeV [4,26]. For example [26],

$$T_d \lesssim 2\ \mathrm{GeV} \times \left(\frac{0.03}{|K|}\right)^{1/2}\left(\frac{m_\phi}{1\ \mathrm{TeV}}\right)^{1/2}\left(\frac{10^{20}}{Q}\right)^{1/2}$$

which is lower than the typical freeze-out temperature of WIMPs, $T_f \simeq m_\chi/24$. Thus, the LSPs generated in Q-ball decays do not subsequently thermalize. Nor will the baryon asymmetry be washed out by sphaleron effects since $T_d < T_{ew}$ [4]. Note that Q-balls decay prior to BBN and thus do not spoil its successful predictions.

In the EMD scheme, T_R also must be rather low. (This justifies neglecting thermal effects in Eq. (1).) In fact, unless $T_R \lesssim 10^{3-5}$ GeV, Q-balls could thermalize [4]. In order to preserve the Ω_b–Ω_{DM} relation, one needs to suppress the neutralino population from freeze-out. For this to happen, it would be sufficient to assume $T_R \lesssim T_f$.

It is easy to see why in the EMD scenario, the ratio Ω_b/Ω_χ should be less than 1. The Q-ball is basically a huge "bag" of squarks. It decays predominantly via $\tilde{q} \rightarrow q + \chi$. Thus, for one unit of a baryon number, at least $N_\chi \geq 3$ units of nonbaryonic number density are created. (This number can be larger than 3 if one takes into account additional decays of squarks into heavier charginos and neutralinos which then cascade decay into the lightest neutralino, which are model dependent.) In other words, the LSP number density n_χ after Q-ball decay is given by $n_\chi = N_\chi f_B n_b$, where f_B is the fraction of baryon asymmetry carried by the AD field ϕ that is transferred into Q-balls. From lattice calculations, $f_B \simeq 1$ [27]. Assuming that the LSPs subsequently do not undergo any significant self-annihilation, and since in general $\Omega h^2 = mY = mn/s$, this can be recast into

$$\frac{\Omega_b}{\Omega_\chi} = \frac{m_n Y_b}{m_\chi Y_\chi} = \frac{m_n}{m_\chi}\frac{1}{f_B N_\chi}, \quad (7)$$

where m_n denotes the mass of a nucleon and m_χ the mass of the neutralino. It is clear that Eq. (7) implies Ω_b/Ω_χ to be less than one but not $\ll 1$. In the EMD scenario, not only are both types of matter simultaneously produced but also a right ratio of their abundances is predicted.

Unfortunately, this attractive picture runs into a serious problem of over-producing neutralinos, as noticed already by EMD themselves [5,28]. Since $Ym \simeq 3.9 \times 10^{-10}(\Omega h^2/0.11)$ GeV, one can rewrite Eq. (7) as

$$m_\chi \simeq 1.5\ \mathrm{GeV}\left(\frac{3}{N_\chi}\right)\left(\frac{1}{f_B}\right)\left(\frac{0.86\times 10^{-10}}{n_b/s}\right)\left(\frac{\Omega_\chi h^2}{0.11}\right). \quad (8)$$

In order to remain consistent with the values of n_b and $\Omega_\chi h^2$ derived from observations, the neutralino mass has to be $\mathcal{O}(1\ \mathrm{GeV})$ which, in the MSSM, is excluded by LEP [10]. Here, we have neglected a possible contribution to the LSP density from freeze-out. If it were significant, the problem would become only worse. Moreover, the condition (8) puts into question an attractive AD mechanism in a large class of SUGRA models.

To circumvent these problems, one has to review assumptions in the above discussions, namely: (i) LSPs produced in Q-ball decay do not annihilate; (ii) The LSP is the lightest neutralino of the MSSM.

If we relax assumption (i), the neutralino LSP with the mass of $\mathcal{O}(10^2\ \mathrm{GeV})$, consistent with LEP, becomes acceptable. Indeed, allowing for subsequent LSP self-annihilation, the LSP abundance will be reduced by $\sim\langle\sigma_\chi v\rangle(T_d - T)$ [29]. If the cross section σ_χ for the LSP (self-)annihilation is large enough, e.g., when the LSP is Higgs-ino or Wino-like [29], the relation between Ω_b and Ω_{DM} is lost. One interesting exception is when the energy density of universe is dominated by Q-ball itself [30].

Alternatively, if we lift assumption (ii), the $\Omega_b - \Omega_{\mathrm{DM}}$ relation may be preserved. One way is to consider, e.g., models with the Higgs sector supplemented by a singlet. If its fermionic partner, the singlino, is the LSP then, for some specific choices of parameters [31], it could be possible to circumvent the LEP bound and perhaps also to suppress the LSP abundance from freeze-out. In the rest of the Letter, we will investigate axino LSP as DM.

3. Axino dark matter from Q-balls.—In general, axinos, like gravitinos, can be produced in both thermal processes (TP) and in nonthermal processes (NTP), e.g., in late decays. TP consists of the scatterings and the decays of particles in the thermal bath. NTP is given by the decay of the Next-to-LSP (NLSP) relic (which, for simplicity, we assume to be the neutralino) from freeze-out or from the decay of Q-balls in our scenario.

The relevant Boltzmann equations can be written as

$$\dot{n}_\chi + 3Hn_\chi = -\langle\sigma_\chi v\rangle(n_\chi^2 - n_{\chi,\mathrm{eq}}^2) + \gamma_Q - \Gamma_\chi n_\chi, \quad (9)$$

$$\dot{n}_{\tilde{a}} + 3Hn_{\tilde{a}} = \langle\sigma v\rangle_{ij} n_i n_j + \langle\sigma v\rangle_i n_i + \Gamma_\chi n_\chi, \quad (10)$$

where $\langle\sigma_\chi v\rangle$ is the usual neutralino freeze-out term, γ_Q denotes the contribution to χ from Q-balls decay, Γ_χ is the decay rate of the neutralino, $\langle\sigma v(i + j \rightarrow \tilde{a} + \ldots)\rangle_{ij}$ and $\langle\sigma v(i \rightarrow \tilde{a} + \ldots)\rangle_i$ are the scattering cross section and the decay rate for the thermal production of axinos.

The total NLSP abundance is given by

$$Y_\chi = N_\chi f_B \frac{n_b}{s} + Y_\chi^{\mathrm{TP}} \quad (11)$$

where $Y_\chi^{\mathrm{TP}} \simeq H/s_{|_{T=m_\chi}} m_\chi/T_f/\langle\sigma v\rangle_{\mathrm{ann}}$, as usual. Since $n_{\tilde{a}} = n_\chi$, owing to R-parity conservation, the resulting number density of axino is given by

$$Y_{\tilde{a}} = Y_{\tilde{a}}^{\mathrm{NTP}} + Y_{\tilde{a}}^{\mathrm{TP}}, \quad (12)$$

with

$$Y_{\tilde{a}}^{\mathrm{NTP}} = 1 \times 10^{-10} N_\chi \left(\frac{f_B}{1}\right)\left(\frac{n_b/s}{1 \times 10^{-10}}\right) + Y_\chi^{\mathrm{TP}}, \qquad (13)$$

where $Y_{\tilde{a}}^{\mathrm{TP}}$ denotes the axinos produced by thermal processes. Since typically $Y_\chi^{\mathrm{TP}} \sim 10^{-11}$, one can see that nonthermal production of axinos due to the thermal relic NLSPs decay can easily be negligible compared to that from Q-ball production, and its contribution to $\Omega_{\tilde{a}} h^2$ is further suppressed by the ratio $m_{\tilde{a}}/m_\chi$. The thermally produced axino $Y_{\tilde{a}}^{\mathrm{TP}}$ also can be subdominant, say $Y_{\tilde{a}}^{\mathrm{TP}} \lesssim 10^{-11}$, for $T_R \lesssim 100$ GeV [15,32]. Hence, the axino dark matter density is estimated as

$$\left(\frac{\Omega_{\tilde{a}} h^2}{0.11}\right) \simeq \left(\frac{m_{\tilde{a}}}{1.5\ \mathrm{GeV}}\right)\left(\frac{N_\chi f_B}{3}\right)\left(\frac{\Omega_b h^2}{0.02}\right). \qquad (14)$$

One can see that the baryon asymmetry and the dark matter abundance are readily linked.

As mentioned above, axino mass is strongly model dependent; in particular, it critically depends on how the visible and hidden sectors are coupled [11–13]. At tree level, either $m_{\tilde{a}} = \mathcal{O}(m_{3/2})$ or $\mathcal{O}(m_{3/2}^2/f_a) = \mathcal{O}(\mathrm{keV})$. However, in the latter case, trilinear terms can generate a substantial 1-loop correction of order $f_Q^2/8\pi^2 A$, where f_Q is the Yukawa coupling of the heavy quark to a singlet field containing the axion, which gives $m_{\tilde{a}}$ in the range of a few tens of GeV or less [12,13].

The final check point is the compatibility with successful predictions of BBN. However, this is not really a problem for the axino LSP because its interactions are less suppressed than those of the gravitino, roughly by $(M_P/f_a)^2$ and, so long as the NLSP is heavier than about 150 GeV, axinos are produced before the time of BBN [15]. In contrast, the gravitino LSP would be produced in late NLSP neutralino decays, which faces strong constraints from BBN [33,34].

4. Conclusions.—We have shown that the framework with cold dark matter axino LSP produced in Q-ball decays can explain the abundance of dark matter and the baryon asymmetry simultaneously and may be an answer to the $\Omega_b \sim \Omega_{\mathrm{DM}}$ coincidence. In this approach, the similarity between Ω_b and Ω_{DM} is explained by basically only the axino mass of order of (sub-)GeV. The essential property of Q-ball decays is that one can predict the number of SUSY particles per one baryonic charge from Q-ball N_χ. A characteristic feature is low reheat temperature T_R of 10^2 GeV.

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