

## Asymmetric Higgsino Dark Matter

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In the supersymmetric framework, prior to the electroweak phase transition, the existence of a baryon asymmetry implies the existence of a Higgsino asymmetry. We investigate whether the Higgsino could be a viable asymmetric dark matter candidate. We find that this is indeed possible. Thus, supersymmetry can provide the observed dark matter abundance and, furthermore, relate it with the baryon asymmetry, in which case the puzzle of why the baryonic and dark matter mass densities are similar would be explained. To accomplish this task, two conditions are required. First, the gauginos, squarks, and sleptons must all be very heavy, such that the only electroweak-scale superpartners are the Higgsinos. With this spectrum, supersymmetry does not solve the fine-tuning problem. Second, the temperature of the electroweak phase transition must be low, in the (1–10) GeV range. This condition requires an extension of the minimal supersymmetric standard model.

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*Introduction.*—The matter content of our Universe is made of two main components: the dark matter (DM), with  $\Omega_{\text{DM}} \sim 0.20$ , and baryons, with  $\Omega_{\text{b}} \sim 0.04$ . Intriguingly, neither of these numbers can be explained within the standard model (SM) of particle physics. The most intensively studied scenarios involve very different mechanisms to explain these two numbers. The DM relic abundance is explained by a freeze-out of a weakly interacting massive particle number density that occurs when its annihilation rate becomes slower than the expansion rate of the Universe. The baryon relic abundance is explained by an asymmetry between baryons and antibaryons. Under these circumstances, there is no explanation of the fact that the energy densities of the DM and the baryons are surprisingly close to each other, which is then just a coincidence. It would be more satisfying if it could be naturally explained. This can be the case if the DM density were also the result of an asymmetry, rather than of freeze-out [1–12]. This type of scenario comes under the name of asymmetric dark matter.

One of the best motivated extensions of the SM is the minimal supersymmetric standard model (MSSM). Extending this framework, there are several ways to generate the baryon asymmetry. Perhaps the most plausible one is via leptogenesis (for a review, see [13]). Prior to the electroweak phase transition (EWPT), when the Higgsino is of Dirac nature, leptogenesis generates also Higgs and Higgsino asymmetries that are initially of similar size to the lepton asymmetry (see, e.g., [14]). Regardless of the source of the baryon and Higgsino asymmetries, however, the conditions of chemical equilibrium imply that, at temperatures above the EWPT, a nonzero baryon asymmetry

requires that there is also a nonzero Higgsino asymmetry. Could such a Higgsino asymmetry survive and lead to asymmetric Higgsino dark matter after the EWPT, when the Higgsino becomes of Majorana nature? This is the question that we address in this work.

Imagine that the following set of conditions applies: (1) the Higgsino is the lightest supersymmetric particle; (2) before the EWPT, Higgsino number-changing interactions become slow enough that the Higgsino asymmetry remains constant; (3) after the EWPT, Higgsino–anti-Higgsino oscillations are slow, or if they are fast, the Higgsino–anti-Higgsino annihilation rate is slow. Then, a rather large relic Higgsino density can survive. In this work, we study the constraints on the supersymmetric spectrum and on cosmological parameters that follow from imposing this set of conditions.

*The Higgsino asymmetry.*—The asymmetry in a particle  $x$  number density  $\Delta n_x \equiv (n_x - n_{\bar{x}})$  is related to its chemical potential  $\mu_x$  via (for  $\mu_x \ll T$  [15])

$$\Delta n_x^{\text{eq}} = \frac{g_x T^2 \mu_x}{6} K(z_x), \quad (1)$$

where  $g_x$  is the number of internal degrees of freedom of the particle  $x$ ,  $z_x \equiv m_x/T$ ,  $K(z_x \ll 1) = 2(1)$  for bosons (fermions), and  $K(z_x \gg 1)$  is exponentially suppressed. We are interested in relating the Higgsino asymmetry to the baryon asymmetry. Under the conditions that will be of interest to us (see below), all of the sfermions are much heavier than the Higgsinos and the sfermion number densities are negligible because the relevant  $K(z)$  factors are exponentially suppressed. On the other hand, for the quarks and leptons  $K(z) = 1$ .

With these assumptions, the baryon and lepton asymmetries are given by

$$\begin{aligned}\Delta n_B &= \frac{T^2}{6}(2\mu_Q + \mu_u + \mu_d), \\ \Delta n_L &= \frac{T^2}{6}(2\mu_L + \mu_e),\end{aligned}\quad (2)$$

where  $\mu_\psi \equiv \sum_i \mu_{\psi_i}$  ( $i = 1, 2, 3$  is a generation index). We define the comoving asymmetry via  $\Delta Y_x \equiv \Delta n_x/s$ , where  $s$  is the entropy density. We follow the derivation of Refs. [16,17]. Imposing the conditions of fast gauge, Yukawa, and sphaleron interactions, the Majorana nature of the gauginos, the Dirac nature of the Higgsinos (prior to the EWPT), and hypercharge neutrality, we obtain

$$\frac{\Delta Y_{\tilde{h}}}{\Delta Y_B} = -\frac{2K(\tilde{h})}{12 + 3[K(h_u) + K(h_d) + 2K(\tilde{h})]}.\quad (3)$$

Here  $\tilde{h}$  stands for the Dirac Higgsino, and  $h_u, h_d$  stand for the scalar Higgs doublets of hypercharge  $\pm 1/2$ . Equation (3) relies on superequilibrium (SE), that is, chemical equilibrium between Higgs particles and Higgsinos [18,19]. At the end of the epoch of SE, a Higgsino asymmetry is conserved until the EWPT.

As long as the Higgsinos are relativistic, the right hand side of Eq. (3) is (in absolute value) in the range 0.10–0.15. If the Higgsinos become nonrelativistic while SE persists, then  $\Delta Y_{\tilde{h}}$  is quenched by the  $K(\tilde{h})$  factor in Eq. (3). Such quenching of  $\Delta Y_{\tilde{h}}$  is not allowed if Higgsinos are to provide the DM. Thus, we are led to impose that SE must be broken while Higgsinos are still relativistic, freezing  $\Delta Y_{\tilde{h}}$  at a value

$$\Delta Y_{\tilde{h}} \approx -10^{-11}.\quad (4)$$

In the next section we find the conditions on the supersymmetric spectrum that would lead to early breakdown of SE.

*Non-superequilibrium (NSE).*—Our goal is to find the conditions on the particle spectrum such that NSE occurs while the Higgsinos are relativistic  $T_{\text{NSE}} > \mu$ . We use  $\Gamma < H$  as the rough criterion for an interaction to be out of equilibrium, where the Hubble expansion rate is given by  $H \approx 10T^2/m_{\text{pl}}$  and the interaction rate is given by  $\Gamma = n\langle\sigma v\rangle$ ,  $n$  being the number density of target particles.

Higgsino number is violated in Higgs-Higgs scattering ( $hh \rightarrow \tilde{h}\tilde{h}$ ) and Higgs-anti-Higgsino scattering ( $h\tilde{h}^c \rightarrow h^c\tilde{h}$ ). These processes arise from the effective Lagrangian

$$-\mathcal{L}_{\text{eff}} = \frac{1}{\Lambda_u}\tilde{h}_u\tilde{h}_u h_u^* h_u^* + \frac{1}{\Lambda_d}\tilde{h}_d\tilde{h}_d h_d^* h_d^*,\quad (5)$$

generated at tree level by gaugino and at one loop by quark-squark diagrams. Requiring that  $h_u h_u \rightarrow \tilde{h}_u \tilde{h}_u$  is not in equilibrium at  $T \sim \mu$  puts a lower bound on  $\Lambda_u$ ,

$$\Lambda_u \geq 3 \times 10^9 \text{ GeV} \left( \frac{\mu}{1 \text{ TeV}} \right)^{1/2}.\quad (6)$$

If  $h_d$  is not heavy and decoupled, then a similar bound holds for  $\Lambda_d$ .

The gaugino contributions to (5) are given by

$$\frac{1}{\Lambda_u} = \frac{1}{\Lambda_d} = \frac{g'^2}{8M_1} + \frac{g^2}{8M_2}.\quad (7)$$

The stop and sbottom contributions are given by

$$\begin{aligned}\frac{1}{\Lambda_u} &= \frac{3\alpha_W^2 m_t^3}{2m_W^4 s_\beta^4} \sin 2\theta_t \ln \frac{m_{\tilde{t}_2}^2}{m_{\tilde{t}_1}^2}, \\ \frac{1}{\Lambda_d} &= \frac{3\alpha_W^2 m_b^3}{2m_W^4 c_\beta^4} \sin 2\theta_b \ln \frac{m_{\tilde{b}_2}^2}{m_{\tilde{b}_1}^2}.\end{aligned}\quad (8)$$

Equation (6), therefore, implies

$$\begin{aligned}M_i &\geq 10^8 \text{ GeV} \left( \frac{\mu}{1 \text{ TeV}} \right)^{1/2}, \\ \sin 2\theta_i \ln \frac{m_{\tilde{t}_2}^2}{m_{\tilde{t}_1}^2} &\lesssim 10^{-6} \left( \frac{1 \text{ TeV}}{\mu} \right)^{1/2}.\end{aligned}\quad (9)$$

Finally, decoupling the decays and inverse decays of sfermions to Higgsino-fermion [17] at  $T_{\text{NSE}} > \mu$  requires all of the sfermions to be heavy  $m_{\tilde{f}} \gtrsim (10\text{--}40)\mu$ . Here, the stronger bound refers to top squarks and the weaker bound refers to the electron superpartners.

We note that since, as shown in the next section,  $\mu > T_{\text{EWPT}}$  the above spectrum guarantees that SE ends before the EWPT.

*The spectrum.*—The Higgsino spectrum includes two neutral mass eigenstates, with mass splitting  $\Delta m_0$  [20], and a charged Higgsino, split by  $\Delta m_+$  from the lightest neutral Higgsino [21,22]. As concerns  $\Delta m_+$ , this quantity does not violate Higgsino number and in our scenario it is therefore dominated by gauge loops  $\Delta m_+ \approx \frac{1}{2}\alpha_W s_W^2 m_Z \sim 350 \text{ MeV}$ . The results of the previous section have, however, interesting implications for  $\Delta m_0$ . The operators of Eq. (5) lead to

$$\begin{aligned}\Delta m_0 &= \Delta_u + \Delta_d, \quad \Delta_u = \langle h_u \rangle^2 / \Lambda_u, \\ \Delta_d &= \langle h_d \rangle^2 / \Lambda_d.\end{aligned}\quad (10)$$

Equation (6) implies  $\Delta_u \lesssim 10 \text{ KeV}$ . An analogous bound  $\Delta_d \lesssim (10 \text{ keV})/\tan^2\beta$  applies unless the second Higgs doublet  $h_d$  is much heavier than the Higgsinos. We now show that experimental constraints lead us to choose the second possibility, namely  $m_{h_d} \gg \mu$ .

Constraints on the Higgsino spectrum in our scenario come from direct and indirect DM searches. The dominant Higgsino-nucleus interactions are spin-independent inelastic interactions (see, e.g., [22–24]). The cross section with a single nucleon, in the limit of zero mass splitting, is  $\sigma_n \approx 10^{-39} \text{ cm}^2$  for  $\mu > 200 \text{ GeV}$ . This is orders of magnitude above current bounds. The only way to evade these bounds (and maintain Higgsino LSP) is by a large enough mass splitting that will make the inelastic scattering kinematically forbidden [25–27]. The minimum mass splitting required is

a function of the Higgsino mass, but in the entire range of interest for  $\mu$  it is smaller than 400 keV. We learn that the sbottom contribution to  $\Delta m_0$  must be substantial,  $\Delta_d \gtrsim 400$  keV. As a result, the only configuration consistent with early breakdown of SE  $T_{\text{NSE}} > \mu$  is one with a very heavy second doublet  $m_{h_d} \gg \mu$ .

As can be seen from Eqs. (8) and (10), requiring

$$\Delta m_0 \gtrsim 1 \text{ MeV} \quad (11)$$

constrains the sbottom sector to satisfy,

$$\sin 2\theta_b \gtrsim 2 \times 10^{-2} \left( \frac{\tan \beta}{20} \right)^{-2}. \quad (12)$$

Finally, with  $\Delta m_0 \gg 10$  eV, the present Higgsino population is symmetric (see below) and therefore annihilates and may provide signals in indirect searches for DM. The annihilation cross section into  $WW$  and  $ZZ$  pairs is given by [28]

$$\langle \sigma^{\text{ann}} v \rangle \approx 10^{-26} \text{ cm}^3 \text{ sec}^{-1} (1 \text{ TeV}/\mu)^2. \quad (13)$$

The bound on the cross section from the Fermi-LAT data [29], when compared with Eq. (13), requires

$$\mu \gtrsim 190 \text{ GeV}. \quad (14)$$

*Oscillations, damping and expansion.*—At the EWPT, the Higgs acquires a VEV, and the Higgsinos mix with the gauginos. The resulting propagation eigenstates change from Dirac to Majorana fermions, and oscillations begin. On the other hand, incoherent interactions with the plasma continue and damp the oscillations. In this section, we study the time evolution of the Higgsino system ( $\tilde{h}, \tilde{h}$ ) under the simultaneous effects of oscillations, annihilations, damping, and the expansion of the Universe. To do so, we employ the formalism of the density matrix. This formalism was originally developed to study neutrinos [30–33] and we adapt it to our case. Our equations are consistent with those of Ref. [34], which deals with closely related issues (see also [35]). The quantum rate equations for  $Y_\mu \equiv n_\mu/s$  (where  $n_0 \pm n_3$  are the number densities of the Higgsinos and anti-Higgsinos, and  $n_1 \mp in_2$  are the off diagonal elements of the density matrix) read

$$\begin{aligned} \frac{d}{d \log z} \mathbf{Y} &= -\frac{1}{H} \begin{pmatrix} D & V & 0 \\ V & D & \Delta m_0 \\ 0 & -\Delta m_0 & 0 \end{pmatrix} \mathbf{Y}, \\ \frac{d}{d \log z} Y_0 &= \frac{s}{H} \langle \sigma^{\text{ann}} v \rangle \times \left[ 2Y^{\text{eq}} \bar{Y}^{\text{eq}} - \frac{1}{2} Y_0^2 + \frac{1}{2} Y_3^2 \right. \\ &\quad \left. + G(Y_1^2 + Y_2^2) \right], \end{aligned}$$

where  $z \equiv \mu/T$ . The damping (or decoherence) factor  $D$  and the effective matter potential  $V$  are given by

$$D = 2 \sum_f n_f \langle \sigma_{\tilde{h}+f \rightarrow \tilde{h}+f} v \rangle \approx 2G_f^2 T^5,$$

$$V = 8\zeta(3) \alpha_W \eta_B T^3 / (\pi m_W^2).$$

Here  $D$  is proportional to the (flavor-sensitive) elastic scattering cross section and to the total number density of the massless fermions in the plasma. The contribution from (flavor-blind) annihilation is quantitatively small. Here  $V$  is proportional to the elastic scattering amplitude and to the fermion-antifermion asymmetry. In the limit of large damping, the effective rate of oscillation is  $\Gamma_{\text{osc}} \sim (\Delta m_0)^2/D$ , as obtained in [12]. The  $G$  factor measures the ratio between the annihilation cross section with and without including coannihilations.

The initial conditions, at the EWPT, are the following:

$$Y_0 = n_{\tilde{h}}/s, \quad Y_1 = Y_2 = 0, \quad Y_3 = -10^{-11}, \quad (15)$$

where  $n_{\tilde{h}}$  is the solution of the Boltzmann equations (at the EWPT) with constant asymmetry, and the value of  $Y_3 = \Delta Y_{\tilde{h}}$  is taken from Eq. (4).

*Asymmetry-assisted Higgsino DM.*—We aim to solve for  $Y_0(\infty)$ , which gives the final total number density in Higgsinos, and for  $Y_3(\infty)$ , which gives the final Higgsino asymmetry. In particular, to provide  $\Omega_{\text{DM}} h^2 \approx 0.11$ , our scenario needs to fit

$$Y_0^{\text{obs}} \approx 7.6 \times 10^{-13} (1 \text{ TeV}/\mu). \quad (16)$$

We consider three free parameters:  $\mu$ ,  $\Delta m_0$ , and  $T_{\text{EWPT}}$  and ask whether there is a range of these parameters where such a fit is achieved. If the asymmetry is washed out before the symmetric decoupling temperature  $T_{\text{dec}}^{\text{sym}} \approx \mu/25$ , then the present DM relic density is the standard symmetric one, as if an initial asymmetry was never generated:

$$Y_0^{\text{sym}} \approx 5.9 \times 10^{-13} (\mu/1 \text{ TeV}). \quad (17)$$

Comparing Eqs. (16) and (17) we learn that for  $\mu \sim 1.1$  TeV, Higgsinos can account for DM without asymmetry. Since in the presence of an asymmetry the total number density is always larger than in the symmetric case, this puts an upper bound of 1.1 TeV on the Higgsino mass. Thus, the range of interest is  $190 \text{ GeV} \leq \mu \leq 1.1 \text{ TeV}$ . Within this range, the lighter the Higgsino, the larger the asymmetry that is required in order to satisfy (16).

We can obtain  $Y_0(\infty) > Y_0^{\text{sym}}$  if either of the following two conditions applies: (1)  $\Delta m_0$  is small enough that the oscillations are slow, and at least part of the Higgsino asymmetry survives down to  $T < T_{\text{dec}}^{\text{sym}}$ ; (2) The EWPT occurs late enough that annihilations are already slow when oscillations begin  $T_{\text{EWPT}} < T_{\text{dec}}^{\text{sym}}$ . To study the first possibility, we fix  $T_{\text{EWPT}} = 100$  GeV and solve for  $Y_0(\infty)$  as a function of  $\mu$  and  $\Delta m_0$ . We find that the lighter the Higgsinos are, the smaller the mass splitting that is required to provide the DM abundance. The reason is that smaller  $\mu$  leads to smaller  $T_{\text{dec}}^{\text{sym}}$ , and the asymmetry is required to survive to later times. The required  $\Delta m_0$  as

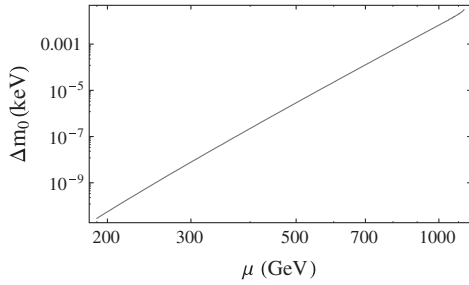


FIG. 1. The required  $\Delta m_0$  as a function of  $\mu$ , with  $T_{\text{EWPT}} = 100$  GeV.

a function of  $\mu$  is shown in Fig. 1. It ranges between ( $10^{-2}$ – $10^{-10}$ ) keV for  $\mu$  in the range (1000–200) GeV. Such a small mass splitting is excluded by direct DM searches. We conclude that it is impossible for the asymmetry to survive once the EWPT takes place.

Second, we fix  $\Delta m_0 = 1$  MeV (the results are not sensitive to changes in  $\Delta m_0$  in the range where it is not excluded by direct searches) and solve for  $Y_0(\infty)$  as a function of  $\mu$  and  $T_{\text{EWPT}}$ . Our numerical result for the required  $T_{\text{EWPT}}$  as a function of  $\mu$  is shown as the smooth line in Fig. 2. Above the line, the final DM abundance is too low. Below the line, the final abundance is too high. The required temperature ranges between (30–0.1) GeV for  $\mu$  in the range (1000–200) GeV. For  $T \gtrsim \text{GeV}(\mu/1 \text{ TeV})^2$ , we can obtain an approximate analytic solution for the required  $T_{\text{EWPT}}$ ,

$$T_{\text{EWPT}} = 33 \text{ GeV} \left( \frac{\mu}{1 \text{ TeV}} \right)^3, \quad (18)$$

shown as the dashed line in Fig. 2.

We conclude that a viable scenario of asymmetric Higgsino dark matter could have occurred as follows: a Higgsino number asymmetry of a size that is about a factor of 10 smaller than the baryon asymmetry exists before the EWPT. At the EWPT, the asymmetry is quickly washed out due to Higgsino–anti-Higgsino oscillations. The resulting symmetric Higgsino population is (for masses below TeV) much larger than the would-be population without an initial asymmetry. It survives if the phase transition occurs at a temperature that is somewhat low, of order (1–10) GeV. We note that such a low temperature requires a rather strong phase transition. Given that the MSSM with the spectrum specified in this work gives  $T_{\text{EWPT}} \approx 150$  GeV, the existence of new degrees of freedom other than those of the MSSM seems necessary. Further analysis of this requirement is beyond the scope of the current paper and we postpone it to future work.

*Conclusions.*—Within the framework of the MSSM, we ask whether the Higgsino could be a viable asymmetric dark matter candidate. We find that the answer is in the affirmative, provided that the following constraints on the supersymmetric spectrum are satisfied: (a) electroweak gauginos are heavier than  $10^8$  GeV; (b) sfermions are

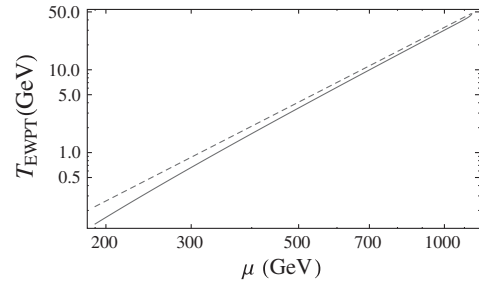


FIG. 2. The required  $T_{\text{EWPT}}$  as a function of  $\mu$ , with  $\Delta m_0 = 1$  MeV. The approximate analytic solution (18) is shown in dashed line.

heavier than  $10^4$  GeV; (c) the stop mixing angle is small, and the sbottom mixing angle is large; (d) Higgsinos are in the range (200–1000) GeV. In addition, the temperature of the electroweak phase transition must be somewhat low, of order (1–10) GeV.

The supersymmetric spectrum is somewhat reminiscent of split supersymmetry models [36,37]. The supersymmetric flavor problem is solved. Grand unification remains a viable possibility. Supersymmetry does not solve the fine tuning problem. It does however explain both the baryon asymmetry and the dark matter abundance and relates the two. The initial source of both asymmetries could be leptogenesis.

This scenario, where the only new particles at the electroweak scale are the Higgsinos, poses a challenge to the LHC. Work on experimental and observational signals is in progress.

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