Protecting the Baryon Asymmetry in Theories with R-parity Violation

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Abstract

We propose a mechanism for hiding the primordial baryon asymmetry from interactions that could wash it out. It requires the introduction of a baryon number carrying singlet which is in equilibrium in the early universe and shares any existing baryon asymmetry. It decouples from the Standard Model particles before all the interactions required to wash out the asymmetry are in equilibrium (\(T \simeq 10\) TeV), and decays after the electroweak phase transition, but before nucleosynthesis. This mechanism can conserve a baryon asymmetry in models (a) with \(B - L = 0\), such as many \(SU(5)\) GUTs, or (b) with \(B - L\) violating interactions in thermal equilibrium, such as SUSY with broken \(R\)-parity. As a result, cosmological constraints on \(R\)-parity violating operators are relaxed considerably.

Making the observed baryon asymmetry of the universe (BAU) \([1, 2, 3]\) in supersymmetric theories \([4]\) with \(R\)-parity violation \([5, 6]\) can be a challenging task. The difficulty is that such theories contain \(B - L\) violating interactions which are naturally in thermal equilibrium above the electroweak phase transition, in conjunction with the anomalous \([7]\) \(B + L\) violating Standard Model processes. Together, these interactions can wash out any BAU present in the early universe \([8, 9]\). The Hubble expansion rate at \(T \simeq 100\) GeV is then so much smaller than particle interaction rates, that it is difficult (but not impossible \([10]\)) to find enough perturbative out-of-equilibrium dynamics to (re)generate an asymmetry. Since Supersymmetric theories with \(R\)-parity violation have recently attracted attention \([3, 4]\), it is of interest to consider how this problem can be avoided. One way around it is to create the asymmetry at \([3]\) or after the electroweak phase transition \([3, 10]\). In this letter, we follow a second approach, which is to assume the asymmetry is generated earlier, but protected by an approximate symmetry. We hide the primordial BAU in a pair of \(SU(3) \times SU(2) \times U(1)\) singlets, \(S\) and \(\bar{S}\), with unit baryon number \((B = 1)\) during the critical period (ie. the time when all interactions required to washed out any BAU are in thermal equilibrium). As \(S\) and \(\bar{S}\) decay the BAU will be transferred back to the Standard Model quarks.
Three ingredients are required to generate a baryon asymmetry \[1\]: baryon number violation, \(C\) and \(CP\) violation, and some out-of-equilibrium process. These are all present in the Standard Model (SM) at the electroweak phase transition; however, there is insufficient perturbative \(CP\) violation \[12\] in the SM, and the non-perturbative \(B + L\) violation would be in thermal equilibrium after the transition for most allowed values of the Higgs mass \[13\] (in which case any asymmetry produced would be destroyed). This suggests that the observed baryon asymmetry cannot be made in the Standard Model, and is evidence for some kind of new physics.

There are numerous extensions of the Standard Model that include viable baryogenesis mechanisms \[1, 2, 3\]. It is particularly interesting that it may be possible to create the BAU at the electroweak phase transition for certain regions of parameter space in the Minimal Supersymmetric Standard model (MSSM) \[14\]. However, we may not live in these regions of parameter space, so an alternative mechanism is certainly desirable. In this letter, we assume that the asymmetry was created before the phase transition. Two generic and popular mechanisms for this are the out-of-equilibrium decay of heavy GUT (Grand Unified Theory) particles produced in the decay of the inflaton \[15\], or the Affleck-Dine mechanism in Supersymmetry and Supergravity \[16\]. In any case, the asymmetry produced is in thermal equilibrium in the early Universe at temperatures above the electroweak phase transition. If interactions capable of washing out the asymmetry are simultaneously present, the asymmetry will be lost. This can happen for an asymmetry with \(B - L = 0\) within the Standard Model, and for any asymmetry in Supersymmetric models with sufficient \(R_p\) violation \[9\].

Recently, there has been considerable interest in supersymmetric models with broken \(R_p\)-parity \[5\], so we will briefly review their baryogenesis difficulties. \(R_p\) is a multiplicative symmetry \[17\] which assigns to each scalar or fermionic field in the model the charge \((-1)^{3B + L + 2S}\), where \(S\) is the particle spin. \(R_p\) conservation in supersymmetric theories was introduced to eliminate renormalizable \(B\) and \(L\) number violating interactions, which, if present simultaneously, would induce proton decay. Requiring \(R\)-parity conservation to ensure proton stability may be too strong a constraint, since it is sufficient to build models that eliminate either \(L\) or \(B\) violating interactions. However, such models may still have problems preserving the cosmological baryon asymmetry; if interactions that take \(B\) or \(L\) to zero are in thermal equilibrium in the presence of the anomalous \(B + L\) violation, then any previously existing asymmetry would be washed out \[8, 9\]. A primordial asymmetry with \(B - L \neq 0\) can be protected if there is no perturbative \(B\) violation, and at least one lepton flavor is effectively conserved. This constrains the \(L\) violating couplings in one generation to be small enough that they are not in chemical equilibrium at the relevant temperatures. This scenario is incompatible with our theoretical prejudice of relating lepton flavor violation to quark flavor violation. Furthermore, it may be in contradiction with the experimental evidence for neutrino oscillations. The main purpose of this letter is to demonstrate that these cosmological constraints on \(R_p\) violating couplings can be circumvented by introducing new particles that can temporarily store the baryon asymmetry.

Let us consider an extension of the MSSM with the additional fields and their quantum numbers given in table 1. The most general superpotential involving these new particles can be written as\[^1\]

\[
W_{\text{new}} = \lambda' \bar{U} T U^c D^c + \lambda T D^c S + \bar{\lambda} T Q Q + m_T \bar{T} T + m_S \bar{S} S,
\]

\(1\)

where \(U^c\) and \(D^c\) (\(Q\)) are the right (left) handed quark fields and we assume the hierarchy \[^1\]The property of gauge coupling unification can be maintained by including a additional pair of Higgs doublets with mass \(m_T\). For simplicity we neglect all interactions involving these Higgs doublets.
Table 1: The particle content and their quantum numbers.

| fields | SU(3)_c | SU(2)_L | U(1)_Y | B |
|--------|---------|---------|---------|---|
| S      | 1       | 1       | 0       | 1 |
| S̄     | 1       | 1       | 0       | −1 |
| T      | 3       | 1       | −2/3    | 2/3 |
| T̄     | 3       | 1       | 2/3     | −2/3 |

$m_T \gg m_S \gtrsim m_z$. For simplicity we have suppressed the flavor indices in eq. (1). After integrating out the heavy fields $T$ and $T̄$ any communication between the MSSM particles and the fields $S$ and $S̄$ proceeds via the non-renormalizable interaction term

$$W_{NR} = \frac{[\lambda, \lambda']}{m_T} U^c D^c D^c S + H.c.,$$

(2)

where we have abbreviated $[\lambda, \lambda'] U^c D^c D^c \equiv (\lambda_k \lambda'_{ij} - \lambda_j \lambda'_{ik}) U^c_i D^c_j D^c_k$ (here, $i, j, k = 1, 2, 3$ are the generation indices).

This model must satisfy three requirements to protect a primordial baryon asymmetry. The singlet $S$ must be initially in chemical equilibrium with the Standard Model fields, so that $S$ and $S̄$ share the primordial BAU irrespective of its origin (This constraint could be relaxed by assuming that sufficient asymmetry was generated in $S$ and $S̄$ but we prefer to be as general as possible.) The singlets must then decouple from the quarks before enough interactions have come into equilibrium to wash out all the asymmetries present. Finally the singlet must decay after the electroweak phase transition, and sufficiently before nucleosynthesis to restore a homogeneous and isotropic radiation dominated Universe with the baryon asymmetry stored in the quarks, and $\Omega_0 \simeq 1$. We will see that there is room for our model to sit comfortably between these bounds.

The first requirement is easy to meet. The couplings $\lambda$ and $\lambda'$ from equation (1) must be sufficiently large such that $S$ is in chemical equilibrium with the quarks at some time in the early Universe. We assume that the Universe has a reheat temperature, $T_{rh} \simeq 10^8$ GeV, to avoid the gravitino problem [18]. (If it is hotter, then the lower bound we compute would be decreased.) If the triplet mass is less than $10^8$ GeV, it will be present in the thermal bath, and $S$ will have to be brought into chemical equilibrium by the non-renormalizable interaction in eq. (2). Requiring this interaction to be in thermal equilibrium at $T \simeq m_T$, or [9]

$$10^{-2}\lambda^2 m_T, 10^{-2}\lambda'^2 m_T \lesssim H \simeq \frac{20m_T^2}{m_{pl}} ,$$

(3)

neglecting generations and color factors. This implies

$$\lambda^2, \lambda'^2 > 2 \times 10^2 \frac{m_T}{m_{pl}}$$

(4)

for $m_T \lesssim T_{rh}$

or $\lambda^2, \lambda'^2 > 2 \times 10^{-13} m_T / \text{TeV}$.

If $m_T \gg T_{rh} \simeq 10^5$ TeV, then $T$ and $T̄$ will not be thermally produced, and $S$ will have to be brought into chemical equilibrium by the non-renormalizable interaction in eq. (2). Requiring this interaction to be in thermal equilibrium at $T \simeq T_{rh}$ gives

$$\Gamma_{NR} \sim \frac{[\lambda, \lambda']^2 T_{rh}^3}{4\pi m_T^2} \frac{3}{4} \simeq 10^{-2}[\lambda, \lambda']^2 T_{rh}^3 \frac{m_T^2}{m_{pl}^2} > H \simeq \frac{20T_{rh}^2}{m_{pl}}$$

(5)

3
This implies
\[ [\lambda, \lambda'] > 10^{-9} \frac{m_T}{\text{TeV}} \quad \text{for} \quad m_T \gg T_{\text{rh}} \] (6)

For \( m_T \gtrsim T_{\text{rh}} = 10^8 \text{ GeV} \), the triplets \( T \) and \( \bar{T} \) will be present in the thermal bath with a Boltzmann suppressed number density. The decays and inverse decays of \( T \) can then keep \( S \) in chemical equilibrium with the quarks. We estimate the rate for these processes to be of order
\[ \Gamma \simeq \frac{\lambda^2}{16\pi} m_T e^{-m_T/T_{\text{rh}}} \] (7)

Requiring this to be greater than the expansion rate at \( T = T_{\text{rh}} \) gives
\[ \lambda^2 > 10^{-3} \left( \frac{\text{TeV}}{m_T} \right) e^{m_T/T_{\text{rh}}} \quad m_T \simeq T_{\text{rh}} \] (8)

which can be used to join the constraints from eqs. (4) and (6) at \( m_T \simeq T_{\text{rh}} \).

The next step is to determine the lower bound on \([\lambda, \lambda']/m_T\) from requiring that \( S \) and \( \bar{S} \) decouple from the MSSM while there are still asymmetries present in the plasma. Thus, we must first determine the temperature at which the asymmetry gets washed out. At very high temperatures in the early Universe, most Yukawa couplings are too weak to be in thermal equilibrium; as the temperature drops, more and more of them come into equilibrium. For the electron Yukawa coupling this happens at \( T \simeq 1 - 10 \text{ TeV} \). An asymmetry in the right-handed electron, \( e_R \), carries net electric charge, which has to be compensated by asymmetries carried by other particles to ensure charge neutrality of the Universe. Thus, in the MSSM the baryon asymmetry will remain until this temperature \([19]\).

In our model with broken \( R_p \) there are other couplings involving \( e_R \). Hence, the temperature at which the baryon asymmetry gets washed out depends on the strength of the lepton flavor violating interactions. Let
\[ W_R = \mu_i \bar{H} L_i + \frac{1}{2} y_{ijk}^L L_i L_j E_k^c + y_{ab}^D L_i Q_a D_b^c + \frac{1}{2} y_{ijk}^D U_i^c D_j^c D_k^c \] (9)

be the \( R \)-parity violating part of the superpotential. Clearly, it is quite complicated to determine the region in parameter space where \( e_R \) is out of chemical equilibrium at \( T \simeq T_{\text{B}} \equiv 10 \text{ TeV} \) due to the large number of parameters and because any constraint on an individual coupling has to be formulated in a particular basis. A sufficient but not necessary condition is to assume that the \( R_p \) violating couplings involving the \( i \)th lepton generation are smaller than the Yukawa coupling of that generation. In the MSSM thermal mass eigenstate basis this condition on the parameters reads
\[ \frac{\mu_i}{\mu_0} y_{33}^D, y_{ab}^D, y_{ijk}^L \lesssim \frac{g m_\ell}{\sqrt{2 m_W}}, \quad \text{for} \quad i, j, k \leq \ell, \] (10)

where \( m_1 < m_2 < m_3 \) denote the three charged lepton masses and \( y_{33}^D \) is the bottom Yukawa \([20]\). In this case, the baryon asymmetry will generically be preserved until the electron Yukawa comes into chemical equilibrium at \( T \simeq 1 - 10 \text{ TeV} \). We therefore require \( S \) to be out of chemical equilibrium by \( T_{\text{B}} \simeq 10 \text{ TeV} \).
Figure 1: The area in the $[\lambda, \lambda'] - m_T$ plane allowed by the requirements that the singlet $S$ be in equilibrium early in the Universe (this rules out area below the lower diagonal curve), and out of equilibrium by $T \simeq 10$ TeV (this rules out the area above the upper line, and to the left of the left-hand curve). We set $\lambda^2 = \lambda'^2 = [\lambda, \lambda']$.

We get a lower bound on $[\lambda, \lambda']/m_T$ by requiring that the dimension 5 operator of eq. (2) to be out of equilibrium before the temperature drops below $T_\beta$. Requiring the rate $\Gamma_{NR}$ from eq. (5) to be less than the Hubble expansion at $T_\beta$ gives

$$\Lambda_{NR} \equiv \frac{[\lambda, \lambda']}{m_T} < 9 \times 10^{-8} \text{TeV}^{-1}. \quad (11)$$

This constraint applies to all the coupling constant combinations $[\lambda, \lambda']_{ijk}$, because the singlet $S$ has to be out of chemical equilibrium with all the MSSM particles.

The constraint of eq. (11) was estimated from the zero temperature scattering cross-section for processes where a heavy off-shell $T$ is exchanged. At finite temperature in the early Universe, there will also be some number of $T$ particles present in the thermal bath, and their interaction rates below $T_\beta$ must also be small enough to keep $S$ out of chemical equilibrium with the quarks. Using the estimate [eq. (7)] for the interaction rate of $S$ with the quarks via decays and inverse decays of $T$s present in the thermal bath, and requiring this to be less than the expansion rate at $T_\beta$ gives

$$\lambda^2 < 8 \times 10^{-12} \left(\frac{\text{TeV}}{m_T}\right) e^{m_T/T_\beta} \quad (12)$$

In figure 1, we plot the bounds on $[\lambda, \lambda']$ as a function of $m_T$. For simplicity we assume $[\lambda, \lambda'] = \lambda^2 = \lambda'^2$. Note, that the antisymmetrization of $[\lambda, \lambda']$ below eq. (2) only projects out the flavor changing components. In a realistic model with approximate flavor symmetries these
are expected to be suppressed. Hence, the lower bounds on \([\lambda, \lambda']\) which are actually constraints on \(\lambda^2\) and \(\lambda'^2\) from eq. (4) and eq. (8) are very conservative. They are obtained by requiring that the singlet be initially in equilibrium with the MSSM, and then out of equilibrium by \(T \simeq T_{\beta} = 10\) TeV [eq. (11)]. In a realistic three generation model, the lower line corresponding to eqs. (4), (8) and (8), applies to the largest component of \([\lambda, \lambda'\vert_{ijk}, |\lambda_i|^2\) or \(|\lambda_{jk}'|^2\) that couples \(S\) to quarks \(U_i^c, D_i^c\) or \(D_i^c\), one of which carries a primordial asymmetry. We expect that in a generic GUT baryogenesis model, some asymmetry would be created in the second or third generation leptons or quarks. At \(T \simeq 10^8\) GeV, the sphalerons and these particles Yukawas are in equilibrium, so the asymmetry would be shared among them. If \(S\) couples sufficiently strongly to at least one second or third generation particle, it will also acquire an asymmetry.\(^2\)

Masses to the left of the left line are ruled out, because the particles \(T\) would be present in the thermal bath in sufficient numbers to mediate \(S\) to quark transitions [eq. (12)]. This bound applies to all the \(\lambda_j\), where \(j\) is a quark generation index. Clearly, there is a substantial region in the \(m_T - [\lambda, \lambda']\) plane consistent with all the bounds in eqs. (4), (6) and (8)–(12).

We now turn to the third requirement which guarantees that the fermionic and scalar singlets decay in such a way as to preserve a BAU without disrupting primordial nucleosynthesis or other cosmological observations. Thus, we have to look at the mass spectrum and decay pattern of \(S\) and \(\tilde{S}\). We assume that SUSY is broken explicitly by soft SUSY breaking terms [21]. Furthermore, we keep all relevant soft SUSY breaking terms degenerate with the exception of the mass of the lightest SUSY partner (LSP), \(m_{\text{LSP}} \ll M_{\text{SUSY}}\). The mass spectrum then looks as follows: there are two complex scalars \(S_{1,2} = (S \pm S^*)/\sqrt{2}\) with mass \(m_{S_{1,2}}^2 = m_{S_{\text{SUSY}}^2} + Bm_{S}\) (\(B\) is the soft SUSY breaking term multiplying the bilinear term in \(W\); in our numerical work we set \(B = M_{\text{SUSY}}\)) and one four component Dirac fermion \(s = (\psi_S, \bar{\psi}_{\tilde{S}})\) with mass \(m_S\).

Now we have to find the range of parameters where the longest-lived singlet, \(s\) or \(S_i (1 = 1, 2)\), decays after the EPT (so that the baryon asymmetry it carries is not washed out by the sphaleron transitions), but sufficiently before nucleosynthesis to not disrupt the process of light element formation. The constraints on low temperature baryogenesis models from nucleosynthesis were studied in [22], where it was shown that nucleosynthesis will proceed in the standard fashion if the post-baryogenesis Universe reheats to look like a standard Big Bang model with \(T \gtrsim 3\) MeV. We take “sufficiently before nucleosynthesis” to mean that the (instantaneous) reheat temperature after the decays of \(S_i\) and \(s\) should exceed 100 MeV; this means that all but \(e^{-1000}\) of the singlets will have decayed by \(T \approx 3\) MeV. If the Universe is radiation dominated when \(S_i\) and \(s\) decay, then we require

\[
10^{-10} \text{sec} \simeq H^{-1}(T_{\text{EPT}}) < \Gamma_S^{-1} < H^{-1}(T \approx 100\,\text{MeV}) \approx 10^{-4}\,\text{sec} \quad (13)
\]

where \(\Gamma_S \equiv \min(\Gamma_{S_i}, \Gamma_s)\) is the decay rate of the particles \(S_i\) and \(s\) with the longest life-time. For a large range of parameter space, \(S_i\) or \(s\) will dominate the energy density of the Universe before they decay. In this case, we want \(\rho_S\), the energy density in \(S_i\) (or \(s\) if they decay last) to be greater than \(\rho_{\text{rad}}(T \approx 100\,\text{MeV}) = g_{\text{eff}}(T)\pi^2T^4/30\). This means \(\Gamma_S > H(T \approx 100\,\text{MeV})\), as above.

\[\text{From the widths for the dominant decay modes}
\]

\[\Gamma(S_i \to 2q + \tilde{q}) \simeq \frac{\kappa m_{S_i}^3 \Lambda_{NR}^2}{f(x)},\]

\[\Gamma(S_i \to 3q + \text{LSP}) \simeq \frac{\alpha_\text{em} m_{S_i}^3 \Lambda_{NR}^2}{4\pi},\]

\[\text{This constraint does not need to be respected if a GUT-scale asymmetry of the right magnitude was created directly in } S.\]
\[ \Gamma(s \to 2\bar{q} + q) \simeq \kappa m_S^3 \Lambda^2_{NR} g(x), \]
\[ \Gamma(s \to 3q) \simeq \left(\frac{\alpha_s}{3\pi}\right)^2 \kappa m_S^3 \Lambda^2_{NR} \min\{1, x^{-1}\}, \tag{14} \]

we obtain
\[ \Gamma_S = \kappa \Lambda^2_{NR} \min\left\{ m_S^3 \left[f(\bar{x}_i) + \frac{\alpha_{em}}{4\pi}\right], m_S^3 \left[g(x) + \left(\frac{\alpha_s}{4\pi}\right)^2 \min\{1, x^{-1}\}\right]\right\}. \tag{15} \]

Here, we set the fine structure constant \(\alpha_{em} = 1/137\) the strong coupling \(\alpha_s = 0.11\) and \(\kappa = 1/(6144\pi^3)\). Furthermore, we have defined \(\bar{x}_i = (m_{\bar{q}}/m_S)^2, x = (m_\bar{q}/m_S)^2\) and
\[ f(x) = 6x(1 + x) \ln(x) + (1 - x)(1 + 10x + x^2), \]
\[ g(x) = 3(2x^2 - x - 2x^3) \left\{ \frac{1}{2} \ln(4x) - \ln(1 - \sqrt{1 - 4x}) - \frac{1}{2} \ln\left[\frac{(1 - 3x) - (1 - x)\sqrt{1 - 4x}}{(1 - 3x) + (1 - x)\sqrt{1 - 4x}}\right] \right\} \]
\[ + (5x - 6x^2 + 1)\sqrt{1 - 4x}, \tag{16} \]

The decay \(s \to 3q\) proceeds at the one-loop level where we have again assumed that squarks and gluinos are mass-degenerate.

In figure 2, we plot the constraints in the \(m_S - \Lambda_{NR}\) plane for \(M_{SUSY} = 300\) GeV. The two horizontal lines correspond to bounds of eqs. (6) and (11); the area between the two lines is allowed. The constraints in eq. (13) rule out anything outside the diagonal curves. In the upper right-hand corner \(S_i\) and \(s\) would decay before the EPT. The lower left-hand corner is ruled out by our conservative requirement that only \(\simeq 10^{-400}\) of the \(s\) and \(S_i\) should be left to decay during nucleosynthesis.

Finally, we must also require that the decay products of \(s\) and \(S_i\) thermalize efficiently before nucleosynthesis, so that they do not disassociate any light elements. With a baryon to photon ratio of \(\eta \simeq 2 - 4 \times 10^{-10}\) having one overly energetic particle in \(10^{10}\) means that there is one dangerous particle per baryon. Thus, we not only have to guarantee that the singlets have decayed sufficiently before nucleosynthesis [eq. (13)] but also that their decay products have sufficiently thermalized by \(T \simeq 3\) MeV. Thus, let us consider the universe shortly after the decay of \(s\) where we expect a soup of squarks, quarks or nucleons, and LSPs with energies of order \(m_S\). The squarks will rapidly decay. The quarks will scatter and emit lower energy quarks and gluons at a rate of order
\[ \Gamma_{\text{scat}} \simeq \frac{\alpha^2_s}{m_S^2} n_s \tag{17} \]
where \(n_s\) was the pre-decay number density of \(s\):
\[ n_s \sim \frac{\rho_{\text{dec}}}{m_S^2} = \frac{3m_S^2 \Gamma_S^2}{8\pi} \tag{18} \]

If the particles to be thermalized are nucleons, the scattering cross-section is much larger, and we expect the time-scale we compute for quark thermalizations to be more than long enough.

The time-scale to scatter once is of order the thermalization time-scale \(\langle \mathcal{E} \rangle\), so we require that \(\Gamma_{\text{scat}}\) (scaled to \(T \simeq 3\) MeV) be more than a factor of a hundred larger than the expansion rate at \(T_n \simeq 3\) MeV:
\[ \frac{\alpha^2_s}{m_S^2} n_s \frac{R_d}{R_n} > 10^2 \times H(T_n) \tag{19} \]
Figure 2: The area in $\Lambda_{NR}$ (in TeV$^{-1}$) vs $m_S$ space allowed after requiring 1) that $S$ be in equilibrium early (rules out below the lower horizontal line) and out of equilibrium by 10 TeV (rules out above the upper horizontal line), 2) that the fermion $s$ decay after the electroweak phase transition (excludes to the right of the right jagged curve), but before nucleosynthesis (excludes to the left of upper left curve), and 3) that the decay products of $s$ and $S_i$ have time to thermalize before nucleosynthesis (this rules out below the lower left curve). These bounds assume $M_{SUSY} = 300$ GeV.
where $R_d$ ($R_n$) is the scale factor when the singlet decays (at nucleosynthesis). This means that at most a fraction $\lesssim e^{-100}$ of the decay products could be unthermalized at $T \simeq 3$ MeV. Expressing $H$ and the ratio of scale factors in terms of the energy density at decay [or equivalently, the decay rate of eq. (14)], we get the bound

$$\Lambda_{NR} = \frac{[\lambda, \bar{\lambda}]}{m_T} > 3 \times 10^{-11} \sqrt{\frac{\text{TeV}}{m_S}} \text{ TeV}^{-1}$$ \hspace{1cm} (20)$$

This is the lower diagonal line on the left in figures 2. As expected, the thermalization bound is only relevant for weakly interacting (small $\Lambda_{NR}$) heavy singlets.

The decay of $s$ can produce particles other than quarks, who must also decay and/or thermalize. If these are Standard Model fermions or gauge bosons, this should be no problem. If they are heavy superpartners, they will also decay soon enough (the Universe at $T \simeq 100$ MeV is $\simeq 3 \times 10^{-5}$ seconds old). The only potential problem is the LSP, which in our model is assumed to be unstable and has a life-time of

$$\tau_{LSP} \simeq 4 \times 10^{-18} s \left( \frac{M_{\text{SUSY}}}{300 \text{ GeV}} \right)^4 \left( \frac{50 \text{ GeV}}{m_{\text{LSP}}} \right)^5,$$ \hspace{1cm} (21)

where $\lambda_R$ stands generically for the dominant $R$-parity violating Yukawa coupling in eq. (9). Requiring that $\tau_{LSP} \lesssim H^{-1}(100 \text{ MeV})$ yields the constraint

$$\lambda_R \gtrsim 2 \times 10^{-7} \left( \frac{M_{\text{SUSY}}}{300 \text{ GeV}} \right)^2 \left( \frac{50 \text{ GeV}}{m_{\text{LSP}}} \right)^{5/2},$$ \hspace{1cm} (22)

which may not be necessary but is certainly sufficient.

We note that for most of the allowed parameter space for this model, the singlet $s$ will dominate the energy density of the Universe before it decays, and can therefore generate substantial entropy. This means that the baryon excess stored in the singlet must be somewhat larger than in the standard scenario. This is straightforward to quantify: if

$$\epsilon \simeq \frac{n_s - n_{\bar{s}}}{n_s + n_{\bar{s}}} \simeq \frac{\mu_s}{T},$$ \hspace{1cm} (23)

where $\mu_s$ is the singlet chemical potential when it was in chemical equilibrium, and $n_s$ is the sum of the number densities of $S_1$, $S_2$ and $s$, then the baryon to entropy ratio today is

$$\frac{n_B}{\sigma} \simeq \frac{\epsilon}{\sigma}$$ \hspace{1cm} (24)

where $\sigma$ is the entropy density (to avoid confusion with the singlet fermion $s$). If $s$ dominates the energy density of the Universe before it decays, then

$$\frac{n_s}{\sigma} \simeq \frac{m_s}{T_{\text{rh},S}}$$ \hspace{1cm} (25)

where $T_{\text{rh},S}$ is the temperature at which the Universe becomes radiation-dominated after the singlets decay (“reheat temperature”). We assumed this was at least 100 MeV, so for singlet masses in the TeV range, this gives $n_B/\sigma \simeq 10^{-4} \epsilon$, or $\epsilon \simeq 10^{-6}$. This is not excessively large.
A mechanism similar to the one we have outlined here can be used to protect an asymmetry with $B - L = 0$ in the Standard Model. The difference in this case, is that the non-renormalizable operator induced by the triplet is suppressed with respect to eq. (2) by an extra power of $m_T$: the triplet is a scalar, and generates the following four fermion operator:

$$V = \frac{[\lambda, \lambda']}{m_T^2} \text{sudd} + \text{H.c.} \quad (26)$$

For $m_T$ less than the reheat temperature after inflation, and for relatively large values of $\lambda$ and $\bar{\lambda}$ ($\gtrsim 10^{-2}$), the non-SUSY version can satisfy the same constraints as its supersymmetric counterpart.

To summarize, we have presented a mechanism to protect the primordial baryon asymmetry from being washed out before the electroweak phase transition. A mechanism of this kind is necessary in the Standard Model if $B - L = 0$ or in supersymmetric models with $R$-parity violation. Our mechanism is quite economical in that it only requires the existence of a pair of singlets and a pair of Higgs triplets and no exotic representations. It is also very generic in that it works in a sizable region of the parameter space and does not require any additional assumptions about the generation of the primordial baryon asymmetry. Some mild constraints on $R$-parity violating parameters remain: the LSP has to decay sufficiently before nucleosynthesis [eq. (22)]; the asymmetry in $\epsilon_R$ should not be washed out before the singlet $S$ decouples from the thermal bath (a sufficient but not necessary set of constraints is presented in eq. (10)]. Note, that the $B$ violating couplings are unconstrained by the baryon asymmetry in our model.

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