R-parity breaking phase transition in the susy singlet majoron model

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Abstract

We consider thermal properties of the susy singlet majoron model. We compute the critical temperature $T_c$ and the subsequent reheating temperature $T_{RH}$ for $R$-parity breaking. Successful baryogenesis constrains the parameter space via the requirements that $T_c$ and $T_{RH}$ are lower than the electroweak phase transition temperature. A further constraint is provided by requiring that the gauge singlet should be kinematically allowed to decay, in order not to have a matter dominated universe at the time of nucleosynthesis. We have made a detailed study of the parameter space and find an upper limit for the susy breaking scalar mass $m_0 < 750$ (900) GeV if $m_{\text{gluino}} = 100$ (1000) GeV, which is valid except for certain special values of the singlet sector parameters.

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In the minimal supersymmetric standard model (MSSM), gauge invariance and supersymmetry allow terms in the superpotential which violate baryon ($B$) or lepton ($L$) number. The absence of these dangerous terms is ensured by assuming a conserved discrete symmetry, called $R$-parity, $R_p = (-1)^{3B+L+2S}$, where $S$ is the spin of the particle. Because of its connection to lepton number, it is interesting to consider situations where $R$-parity is slightly broken, either explicitly or spontaneously. Spontaneous breaking of lepton number is possible in MSSM without introducing additional fields, since the scalar partner of the neutrino may acquire a non-vanishing vacuum expectation value, giving rise to a majoron, which in this case is mainly the supersymmetric partner of the neutrino. Since such a majoron should be detected in $Z^0$ decay, the measurements of the $Z^0$ width rule out this possibility, unless one assumes that there is small explicit $L$ violation, in addition to the spontaneous one, in the MSSM. On the other hand it is possible to break $R$-parity spontaneously in the MSSM with the inclusion of additional singlet superfields, such that the resulting majoron, which is now dominantly a singlet under the gauge group, is not in conflict with the measurements of the $Z^0$ width. A particularly attractive scheme which implements such a spontaneous violation of $R$-parity is the supersymmetric version of the singlet majoron model. This model is the simplest extension which incorporates successfully the spontaneous $L$-(and $R$-parity-)violation in the MSSM. In this model there are, besides the chiral superfields of the MSSM, right-handed neutrino chiral superfields $N_i$ (one for each generation) and an additional gauge singlet superfield $\Phi$, having two units of $L$. The superpotential of the susy singlet majoron model, which is invariant under gauge symmetry and $L$ is, in the standard notation,

$$W = h^U QU^cH_2 + h^D QD^cH_1 + h^E LE^cH_1 + h^\nu LNH_2 + \mu H_1H_2 + \lambda NN\Phi ,$$

(1)

where we have suppressed the generation indices. Eq. (1) contains the usual terms of MSSM together with the Yukawa interactions for the right-handed neutrinos $N_i$ and an interaction term for the gauge singlet $\Phi$. It has been shown, through an analysis of the renormalization group equations (RGE), that for a wide range of parameters one can obtain radiative breaking of $R$-parity in this model. It has also been argued that the $L$-(and $R$-parity-)violating transition may take place after the electroweak phase transition. In that case the sphaleron-induced $B$-violating transitions are frozen out, so that any pre-existing baryon asymmetry is unaffected by this phase transition.

In this paper we extend and improve the analysis of ref. and study the $L$-breaking cosmological phase transition of the susy singlet majoron model in detail. In particular, we compute the critical temperature $T_c$ and the reheating temperature $T_{\text{RH}}$ and compare them with the critical temperature of the electroweak phase transition. We also point out that the gauge singlet sector should decay fast enough in order not to clash with primordial nucleosynthesis. All the constraints, when taken together,
imply that successful baryogenesis in the susy singlet Majoron model requires an upper limit on the susy breaking scalar mass $m_0$.

To start, we recall the tree level potential for the susy singlet majoron model \cite{9, 10}. It consists of three pieces, and can be written along the neutral directions as

$$V = V_H(H_1^0, H_2^0) + V_{\tilde{N}_\Phi}(\tilde{N}_i, \Phi) + V_\nu(\tilde{\nu}_i, \tilde{N}_i, \Phi, H_1^0, H_2^0).$$

(2)

Here, $V_H$ is the Higgs potential of the MSSM. We assume that the couplings $\lambda_{ij}$ are real and work in the basis where they are diagonal, in which case we can write the potential ($\lambda_{ij} = \lambda_i \delta_{ij}$)

$$V_{\tilde{N}_\Phi} = \sum_i m_{\tilde{N}_i}^2 |\tilde{N}_i|^2 + m_\Phi^2 |\Phi|^2 - \left( \sum_i A_i \lambda_i \bar{N}_i^2 \Phi + h.c. \right) + 4 \sum_i |\lambda_i \bar{N}_i \Phi|^2 + 4 \sum_i \lambda_i \bar{N}_i^2 |^2, \quad (3)$$

where $m_{\tilde{N}_i}^2$ and $m_\Phi^2$ are soft susy breaking masses, and $A_i$ are the trilinear couplings. Finally, if we assume that the coupling constants $h^\nu$ are small, and retain only the leading terms in $h^\nu$, then we may write the third piece of the potential Eq. (2) as

$$V_\nu = \sum_i m_{\tilde{\nu}_i}^2 |\tilde{\nu}_i|^2 + \left[ \sum_{ij} h_{ij}^\nu \tilde{\nu}_i (2 \lambda_{ij} \bar{N}_j^* \Phi + h.c.) - \mu \bar{N}_j H_1^{0*} - A_{ij}^{(h)} \bar{N}_j H_2^{0*} \right] + h.c \right]$$

$$+ \frac{1}{8} (g^2 + g') \left[ \sum_i |\tilde{\nu}_i|^2 + 2 \sum_i |\tilde{\nu}_i|^2 (|H_1^{0*}|^2 - |H_2^{0*}|^2) \right], \quad (4)$$

and

$$V_\nu = \frac{1}{8} (g^2 + g') \left[ \sum_i |\tilde{\nu}_i|^2 + 2 \sum_i |\tilde{\nu}_i|^2 (|H_1^{0*}|^2 - |H_2^{0*}|^2) \right], \quad (5)$$

where $m_{\tilde{\nu}_i}^2$ are the soft susy breaking masses, $A_{ij}^{(h)}$ are the trilinear couplings and $g$ and $g'$ are the SU(2) and U(1) gauge couplings, respectively. Assuming, as in the case of MSSM, that all susy breaking masses and trilinear couplings are equal to a universal mass $m_0 \sim 10^2 - 10^3$ GeV and a universal coupling $A$, respectively, at some GUT scale $M_U \sim 10^{16}$ GeV, the values of the parameters in the scalar potential can be obtained by solving the appropriate RGEs. The effect of running of these parameters is to drive SU(2)×U(1) breaking through $V_H(H_1^0, H_2^0)$ with $\langle H_1^0 \rangle \equiv v_1 = v \cos \beta$ and $\langle H_2^0 \rangle \equiv v_2 = v \sin \beta$, where $v = v_1^2 + v_2^2 = (174 \text{ GeV})^2$. Furthermore, for a wide range of parameters the nontrivial global minimum of $V_{\tilde{N}_\Phi}(\tilde{N}_i, \Phi)$ is realized in such a manner as to break the global lepton number and $R$-parity:

$$\langle \Phi \rangle = \phi, \quad \langle \tilde{N}_i \rangle = y_i, \quad (6)$$

where typically $\phi \sim y_i \sim m_0$. Then, nonzero vevs are induced for the sneutrinos through the $h^\nu$-coupling which connects the ordinary doublet Higgs and lepton sectors to the singlet sector. If $m_{\bar{\nu}_i} \sim m_0$, $\phi \sim y_i \sim \mu \sim m_0$, $v_1 \sim v_2 \sim M_W$ and $A^{(h)} \sim m_0$, then $\langle \tilde{\nu}_i \rangle \sim h^\nu M_W$. The vevs $\phi$ and $y_i$ are determined from the minimum of the potential $V_{\tilde{N}_\Phi}$. We assume that $\lambda_3 \gg \lambda_1, \lambda_2$ (a minimum with $\lambda_3 = \lambda_1 = \lambda_2$ is not favoured by the
solution of RGEs \[9, 10\]). The minimum at scale \(Q\) is defined by
\[
\phi = \frac{x}{4\lambda_3}, \quad y_3^2 = \frac{m_{\tilde{N}_3}^2}{4\lambda_3^2 (A_3 - x)}, \quad y_2 = y_1 = 0,
\]
\[
0 = x^3 - 3A_3x^2 + 2(m_{\tilde{N}_3}^2 - m_{\phi}^2 + A_3^2)x - 4A_3m_{\tilde{N}_3}^2, \quad A_3/x > 1.
\]

Lepton number is broken spontaneously if Eq. (7) yields the absolute minimum, which is obtained when
\[
A_3 - m_{\phi} < x < A_3 + m_{\phi}, \quad m_{\phi}^2 > 0.
\]

Assuming universal boundary conditions at the GUT scale \(M_U\), the parameters of the potential \(V_{N\Phi}\) at the low scale \(Q\) can be written as \[9, 10\]
\[
m_{\phi}^2(Q) = \frac{1}{5}m_0^2[2 + (3 - A^2)K^2 + A^2K^4],
\]
\[
m_{\tilde{N}_1,2}^2 = m_0^2, \quad m_{\tilde{N}_3}^2(Q) = 2m_{\phi}^2(Q) - m_0^2,
\]
\[
A_{1,2}(Q) = AK, \quad A_3(Q) = AK^2,
\]
\[
\lambda_{1,2}(Q) = 0, \quad \lambda_3(Q) = \lambda_0 K,
\]
where
\[
K = \left[1 + \frac{5}{4\pi^2}\lambda_0^3 \ln \left(\frac{M_U}{Q}\right)\right]^{-\frac{1}{2}}.
\]

From the scalar potential, Eq. (2), one can derive the mass squared matrix for the scalar bosons \[11\] in a straightforward manner. In this paper we shall only consider the situation when \(\lambda_3\) is the largest of the non-zero values of \(\lambda_i\), so that Eq. (4) holds. In the limit \(h^\nu \to 0\) the mass squared matrix is of a block diagonal form
\[
M_2^2 = \text{diag}(M_H^2, M_{\tilde{\nu}}^2, M_{\tilde{N}_3}^2),
\]
where \(M_H^2\) is effectively the \(2 \times 2\) mass squared matrix of the scalar Higgs bosons in MSSM \[12\], \(M_{\tilde{\nu}}^2\) is the \(3 \times 3\) mass squared matrix of left-handed sneutrinos, and
\[
M_{\tilde{N}_3}^2 = \begin{pmatrix}
A_3\lambda_3y_3^2 & 8\lambda_3^2y_3\phi - 2A_3\lambda_3y_3^2 \\
8\lambda_3^2y_3\phi - 2A_3\lambda_3y_3 & 4\lambda_3^2y_3^2
\end{pmatrix}.
\]

To study the \(L\)-breaking phase transition, we shall focus on the \((\tilde{N}_3, \Phi)\)-sector only. At high temperatures \(\tilde{N}_3\) and \(\Phi\) are brought into equilibrium by decays, inverse decays and ordinary scattering processes. The thermally averaged rate for a given process is, neglecting final state blocking, given by
\[
\Gamma = \int \prod_k d\Pi_k (2\pi)^4 \delta(P_i - P_f) \prod_{n_i} \frac{f_i}{n_i} |M|^2
\]
\[
= 1.4 \times 10^{-2}T^{-1}|M|^2 \quad (1 \to 2 + 3),
\]
\[
= 4 \times 10^{-4}T|M|^2 \quad (2 \to 2),
\]
where \( d\Pi_k = d^3p_k/(2\pi)^32E_k \), \( f_i \) is the momentum distribution of the particles in the initial state, and \( n_i \) is their number density. In Eqs. (17) and (18) a constant matrix element is assumed. These rates are to be compared with the Hubble rate \( H \simeq 23T^2/M_{Pl} \). From Eq. (16) one easily finds that for \( \lambda \) not too small, \( \tilde{N}_3 \) and \( \Phi \) are in equilibrium already at temperatures much above 1 TeV, because of the reactions \( \tilde{N}_3\Phi \rightarrow \tilde{N}_3\Phi \) and \( \Phi \rightarrow \tilde{N}_3\tilde{N}_3 \). The fields of the MSSM can also be assumed to be in equilibrium already at high temperatures.

On the other hand, it is not obvious that the \((\tilde{N}_3, \Phi)\)-system is in thermal contact with \( H^0_1, H^0_2 \) and \( \tilde{\nu} \). This is because the coupling is only through \( h\nu \), which is small. Here the most important reaction is the decay \( H^0_1 \rightarrow \tilde{\nu}_i\tilde{N}_j \), which induces chemical equilibrium at temperatures

\[
T \lesssim 194|\frac{\mu}{\text{GeV}} h_{ij}^{\nu}|^{2/3} \text{ TeV} .
\] (19)

If \( \mu \simeq 100 \text{ GeV} \), there is no thermal contact at the electroweak phase transition if \( h\nu \lesssim 10^{-7} \). This would mean that the \((N_3, \Phi)\)-system and the MSSM would experience different temperatures. In what follows we tacitly assume that this is not the case. For light \( \Phi \), \( \tilde{N}_3 \) and \( H^\pm_2 \) also the scattering process \( \Phi H^+_2 \rightarrow \tilde{N}_3e_L \) is important. For instance, at \( m/T \simeq 0.1 \) the particles are kept in equilibrium at temperatures \( T \lesssim 10^{22}\lambda^2 h^{\nu^2} \). However, for increasing particle masses the thermalization temperature decreases rapidly.

At high temperatures the plasma masses of \( \tilde{N}_3 \) and \( \Phi \) are given by [I]

\[
M^2_\Phi(T, Q) = m^2_\Phi(Q) + \frac{1}{2}\lambda^2(Q)T^2 ; \quad M^2_{\tilde{N}}(T, Q) = m^2_{\tilde{N}}(Q) + \lambda^2(Q)T^2,
\]

where the RGE masses \( m^2(Q) \) are given by Eq. (11), and we have neglected terms of order \( O(h\nu) \). Here \( Q \) is the renormalization point, which in principle is independent of \( T \). However, as in thermal bath the average momenta of the particles are peaked about \( 3T \), it makes sense to choose the renormalization point to depend on temperature, so that \( Q \simeq T \). Setting \( Q \simeq m_{\text{susy}} \simeq m_0 \) as usual would yield masses which would differ by terms of the order of \( O(\lambda) \); for our purposes this difference is inessential. With this choice, we may use Eq. (20) to search for the critical temperature \( T_c \), at which the high \( T \) minimum \( (\tilde{N}_3, \Phi) = (0, 0) \) becomes unstable.

Note that the expressions Eq. (21) are valid if the temperature is much bigger than the masses circulating in the loops. There is a subtlety here in that at some point \( m^2_{\tilde{N}}(Q) \) becomes negative, and the finite \( T \) perturbation expansion, leading to Eq. (20), becomes unstable. This can be remedied by performing resummation in the graph involving \( \tilde{N}_3 \) loop, which effectively replaces \( m^2_{\tilde{N}}(Q) \) by the plasma mass \( M^2_{\tilde{N}} \), but does not change the potential to lowest order in \( \lambda \). As \( M^2_{\tilde{N}} \rightarrow 0 \) for \( T \rightarrow T_c \), the plasma mass of \( \tilde{N}_3 \) may always be assumed to be less than \( T \). Hence loops of \( \tilde{N}_3 \) always contribute to Eq. (20).
The resummed loop involving Φ becomes Boltzmann-suppressed when \( M_{\phi} \gtrsim T \). This implies that if \( m_0^2 \gtrsim 2T_c^2 \), loops of Φ should be neglected. Effectively, this means that \( \lambda^2T^2 \rightarrow \frac{1}{2} \lambda^2T^2 \) in \( M_N^2 \) in Eq. (20). The critical temperature for \( L \)-breaking is, however, in both cases defined through \( M_N^2(T_c, T_c) = 0 \).

For a given \( A \) and \( \lambda \), it is then possible to run the RGEs to find \( T_c \) using Eq. (20). In addition, one has to impose the conditions of Eq. (7) to make sure that a zero temperature global minimum exists for the chosen set of parameters. We have done this numerically, and the result, displaying the allowed region for fixed \( T_c/m_0 \), is shown in Fig. 1.

The parameter space can be constrained by noting that the physical eigenstates of the \( (\tilde{N}_3, \Phi) \)-sector should be unstable. Otherwise the energy density \( \rho_{osc} \) associated with coherent field oscillations about the vacuum would soon after the \( R \)-breaking phase transition start to dominate the energy density of the universe, with disastrous consequences for e.g. primordial nucleosynthesis. Roughly, \( \rho_{osc} \approx m^2v^2(T/T_c)^3 \) where \( m \) and \( v \) are a generic mass and vev, respectively. Thus, in the absence of dissipation, \( \rho_{osc} \) would soon become larger than radiation energy density \( \rho_{rad} \sim T^4 \).

The eigenstates can easily be found from Eq. (15) for a given value of \( \lambda \) and \( A \). It turns out that the heavier eigenstate, with mass \( m_2 \), is predominantly Φ. Because Φ does not couple directly to the MSSM states, it decays either via the small mixing with \( \tilde{N}_3 \) or via an \( \tilde{N}_3 \) virtual state. Here we shall assume that the latter is the case. Requiring that the process \( \Phi \rightarrow \tilde{N}_3\tilde{N}_3 \rightarrow \tilde{N}_3\nu_L\tilde{\chi}_i^0 \), where \( \chi_1^0 \) is the lightest neutralino, is not kinematically forbidden implies \( m_\Phi - m_{\tilde{N}_3} > m_{\tilde{\chi}_1^0} > 18 \text{ GeV} \), where the last figure is the experimental lower bound on LSP [6]. The mass difference \( m_\Phi - m_{\tilde{N}} \) is a function of \( m_0, A \) and \( \lambda \), and the different contours are displayed in Fig. 1. If the gluino is light, \( m_{\tilde{g}} \sim 100 \text{ GeV} \), the experimental lower limit on sneutrino, \( m_\nu > 37.1 \text{ GeV} \) [8], rules out part of the parameter space as indicated in Fig. 1.

Note that the actual decay rate depends on the \( \Phi - \tilde{N}_3 \) mixing and the mixing of the gauge singlets with the MSSM sparticles, which is characterized by the unknown (but small) coupling \( h^\nu \). We have checked that \( \Phi - \tilde{N}_3 \) mixing is indeed small in the physical large \( \lambda \) region. It is conceivable, though, that for some values of \( h^\nu \) the heaviest eigestate could decay directly to MSSM particles already before the onset of nucleosynthesis. In any case there always is some kinematic constraint that must be satisfied in order that the decay is possible.

The critical temperature for \( L \)-breaking can further be constrained by baryogenesis considerations. \( L \)-violating interactions, if in equilibrium with the anomalous electroweak \( B + L \)-violating interactions, will wash out any pre-existing baryon asymmetry. This can be avoided if \( R \)-parity breaking couplings are very small, so that these interactions are never in equilibrium. In this case \( R \)-parity breaking would not be relevant for experiments. Another possibility, as suggested in [9], is to impose the
condition that $L$-breaking takes place after electroweak phase transition, so that the sphaleron induced transitions have already dropped out of equilibrium.

In the small $h\nu$ approximation adopted in this paper, the critical temperature for the electroweak phase transition may be assumed to be the same as in the MSSM. Computing the critical temperature is, however, notoriously difficult task, and it is conceivable that it is dominated by non-perturbative effects like in the non-supersymmetric case [13]. Because the transition is presumably only weakly first order, and the latent heat release is small, for our purposes it suffices to estimate the electroweak critical temperature by the spinoidal instability temperature $T_0$, determined by $\det(M^2_H(T_0)) = 0$. The resummed one-loop mass matrix $M^2_H$ has been presented in [14]. The expressions are lengthy, and we do not reproduce them here. Using these, we have varied the parameters of the MSSM in the following ranges: $1 < \tan\beta < 50$, the trilinear coupling associated with the stop-sector $0 < A_t < 1 \text{ TeV}$ and the Higgs mixing $0 < \mu < 1 \text{ TeV}$. We have explicitly studied two cases, the case of a light gluino with $m_{\text{gluino}} = 100 \text{ GeV}$, and the case of a heavy gluino, $m_{\text{gluino}} = 1 \text{ TeV}$. The resulting $T_0$ is shown in Fig. 2.

Successful baryogenesis requires that $T_c < T_0$. Moreover, we should also require that the reheating temperature $T_{RH}$ after $R$-parity breaking does not exceed $T_0$ so as to make sphaleron interactions operative again [9]. The reheating temperature is given by

$$T_{RH} \simeq \left( \frac{30|V_{\text{min}}|}{\pi^2 g_*} \right)^{\frac{1}{4}} \quad (21)$$

where $V_{\text{min}}$ is the global minimum and we have taken $g_* \simeq 100$. In Fig. 2 we show $T_{RH}$ for a choice of parameter values. In Fig. 3 we have taken all the constraints together to find the allowed parameter space. If $A > \sim 2$, we find the upper limits

$$m_0 \lesssim 750 \text{ GeV} \quad (m_{\text{gluino}} = 100 \text{ GeV}) \quad m_0 \lesssim 900 \text{ GeV} \quad (m_{\text{gluino}} = 1 \text{ TeV}) \quad (22)$$

where the figures refer to the case $\tan\beta = 1.5$; if $\tan\beta = 50$, the limits are slightly relaxed, as can be seen in Fig. 3. However, the limits Eq. (22) are not completely general. If it happens that $A \lesssim 2$ and $\lambda \simeq 1$, $m_0$ can be as high as $4 \text{ TeV}$. Note that for $\lambda \ll 1$ one gets $A \simeq 3$ so that Eq. (22) always holds. We wish to emphasize that in all the cases an upper limit on $m_0$ exists.

To conclude, we have studied in detail the cosmological $R$-parity breaking phase transition in the susy singlet majoron model. In particular, we focussed on the constraints imposed on the model by the requirements of a succesful baryogenesis. The necessary conditions are that $T_c$ and $T_{RH}$ are lower than the electroweak phase transition temperature $T_0$, which depends in a complicated way on the parameters of MSSM. The connection to the gauge singlet sector of the majoron model is provided by the common scalar mass parameter $m_0$. An additional cosmological constraint is that the
gauge singlet sector should be allowed to decay. This is a kinematical constraint that restricts the range of $m_0$. Because susy sparticle spectrum is essentially given by $m_0$, plus D-term and radiative corrections, baryogenesis requires a sparticle spectrum in the susy singlet majoron model which should be observable at LHC.

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Figure 1: The allowed values for $A$ and $\lambda_0$. The solid lines indicate the parameter space for which R-parity is not broken at temperatures $T > T_c$ with $T_c/m_0$ values shown in the figure. The dashed line indicates the parameter space allowed by the requirement of global minimum. The dotted lines are contours of $m_2 - m_1/m_0$, where $m_{1,2}$ are the physical eigenstates of $\Phi$ and $\tilde{N}_3$. From experimental lower limit for $\tilde{\nu}$, the area above the starred-dashed line is not allowed if $m_{\text{gluino}} \sim 100$ GeV, $\mu = A_t = 0$ and $\tan \beta = 1.5$. 
Figure 2: Spinoidal instability temperature $T_0$ (solid and dashed lines) and reheating temperature $T_{RH}$ (dotted lines) as functions of $m_0$. The solid lines correspond to $m_{gluino} = 1$ TeV, $\mu = A_t = 1$ TeV, and the dashed ones to $m_{gluino} = 100$ GeV, $\mu = A_t = 0$. As indicated, tan $\beta = 1.5$ or 50.
Figure 3: The allowed values of $m_0$ and $A$ for $\lambda_0$ values 0.5, 1, 2, 5, and 100. The allowed range for $m_0$, which is the area between the lines, is shown for the cases $m_{\text{gluino}} = 100$ GeV, $\mu = A_t = 0$ with a) $\tan \beta = 1.5$, b) $\tan \beta = 50$, and for $m_{\text{gluino}} = 1$ TeV, $\mu = A_t = 1$ TeV with c) $\tan \beta = 1.5$ and d) $\tan \beta = 50$. 