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Spontaneous Breaking of R parity in the Minimal Supersymmetric Standard Model Revisited.

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Abstract

We reconsider the possibility of spontaneous breaking of R parity in the Minimal Supersymmetric Standard Model. By a renormalization group analysis we find the parameter space in which a sneutrino gets a vacuum expectation value, leading to the spontaneous breaking of the lepton number and to the appearance of a phenomenologically unacceptable massless Goldstone boson. We then analyze the effect of operators giving rise to a tiny amount of explicit violation of lepton number, which could emerge as remnants of physics at some superheavy (Planck or GUT) scale in the low energy effective theory. We show that the conspiracy between the spontaneous and the explicit breaking scales can provide a mass to the Goldstone boson larger than the Z^0 boson mass, hence allowing for a non vanishing sneutrino vacuum expectation value without increasing the invisible width of the Z^0 .

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2

In any supersymmetric (SUSY) extension of the standard model [1] there exist electrically neutral scalar particles which carry lepton (L) number (the so-called sneutrinos) and are also odd under a discrete symmetry, called R parity, which distinguishes between the usual particles of the standard model and their supersymmetric partners. Hence, even taking the minimal SUSY standard model (MSSM), one may conceive the possibility of spontaneous breaking of L , and then of R , through the vacuum expectation value (VEV) of sneutrinos, $\langle \bar{\nu} \rangle$. That such a potentiality indeed can be realized was shown both for the cases of the tree level [2] and radiative [3] spontaneous breaking of L . A major phenomenological consequence is the existence of a massless Goldstone boson (majoron) J , and of its scalar partner ρ , whose mass is of $\mathcal{O}(\langle \bar{\nu} \rangle)$. The LEP impressive results on the Z^0 invisible decay width [4] definitely exclude the process $Z^0 \rightarrow J \rho$, which would be allowed if $\langle \bar{\nu} \rangle \neq 0$, so spontaneous breaking of L through a non vanishing VEV of the sneutrino field is forbidden in the MSSM. On the other hand, one could question the exact conservation of L in any low energy effective lagrangian. Indeed, Planck scale physics could spoil any global symmetry and the net effect would be the appearance of extremely tiny violations of otherwise exact global symmetries in the lagrangians describing physics at the electroweak scale [5]. The problem is even more acute for low energy SUSY lagrangians given that baryon and lepton numbers are no longer automatic symmetries as they are in the standard model.

In this paper we show that even a very tiny amount of L violation in the low energy lagrangian of the MSSM might make points of the parameter space leading to $\langle \bar{\nu} \rangle \neq 0$ phenomenologically allowed. The key point is that the mass of the pseudo-majoron certainly goes to zero in the limit where the explicit breaking of L vanishes, however it is in general not true that its mass is of the same order of magnitude as the typical scale of explicit L violation. There can be quite a large difference, the decay of the Z^0 into the pseudo-majoron becoming kinematically forbidden even though the explicit breaking of L involves

an effective scale orders of magnitude smaller than m_{Z^0} . There already exist examples in the literature of pseudo-Goldstone bosons which acquire large masses even though the explicit violation of some global symmetry is very tiny [6]. The mechanism relies on the presence of more physical mass scales. In our case, in addition to the electroweak and the explicit L violation scales, we have also the sneutrino VEV and it is the interplay among these different scales which gives rise to the above-described phenomenon.

We start by reviewing the conditions for the spontaneous breaking of L in the MSSM. We consider the Higgs potential of the MSSM which is relevant for the correct breaking pattern of $SU(2)_L \otimes U(1)_Y \rightarrow U(1)_{em}$ and L . It reads

$$V_0 = m_1^2 v_1^2 + m_2^2 v_2^2 - 2m_3^2 v_1 v_2 + m_L^2 v_L^2 + \frac{1}{8} g^2 (v_1^2 - v_2^2 + v_L^2)^2, \quad (1)$$

where $v_1 \equiv \langle H_1^0 \rangle$, $v_2 \equiv \langle H_2^0 \rangle$, $v_L \equiv \langle \bar{\nu} \rangle$, and $g^2 \equiv g_2^2 + g_1^2$.

We study the conditions under which the three VEV's are non-vanishing. The condition for the electroweak breaking is

$$\lambda \equiv m_1^2 m_2^2 - m_3^4 < 0, \quad (2)$$

while having $v_L \neq 0$ at the minimum requires

$$\rho \equiv (m_1^2 - m_L^2) (m_2^2 + m_L^2) - m_3^4 = 0. \quad (3)$$

In particular notice that it is eq. (3) to yield the correct condition for $\langle \bar{\nu} \rangle \neq 0$, and not the often quoted relation $m_L^2 + (g^2/4)(v_1^2 - v_2^2) < 0$. This latter inequality implies the presence of an instability of the scalar potential, i.e. of an unbounded direction, but, in general, it is not sufficient to ensure the presence of a correct L violating minimum.

At first sight, eq. (3) seems to imply a fine-tuning among the parameters of the potential (1). However, ρ is actually a scale-dependent quantity and, hence, the condition (3) may be fulfilled at a certain scale Q_0 through the running of the parameters in ρ starting with the usual initial conditions on them at some superlarge scale of supergravity breaking. As

clearly stated in ref. [3], if eq. (3) is satisfied at the scale Q_0 , smaller than the scale at which the electroweak breaking takes place, i.e. $Q_0 < Q_c$, with Q_c such that $\lambda(Q_c) = 0$, then the minimum of eq. (1) can break L through $\langle \tilde{\nu} \rangle \neq 0$. Whether this actually occurs or not depends on another crucial scale of our problem, the supersymmetry breaking scale M_{SUSY} . Here M_{SUSY} denotes the scale at which the SUSY thresholds start appearing and, hence, the scale where the running down from the supergravity breaking scale should be stopped [7]. Indeed, if $M_{\text{SUSY}} > Q_0$, then the lowest value of Q to which it makes sense to run to the potential is $Q = M_{\text{SUSY}}$ and, hence, the lepton number is not spontaneously broken. On the contrary, if the hierarchy $M_{\text{SUSY}} < Q_0 < Q_c$ takes place, then we encounter Q_0 in the running down to M_{SUSY} and, in fact, one has to stop the running itself at Q_0 since at that scale a flat direction appears and the minimum is fixed by radiative correction. Hence, in this latter situation, R parity is broken. For a detailed discussion on these points see ref. [3].

We study the region of the MSSM parameter space where R is broken by considering the renormalization group improved potential (1) plus the one-loop correction to it

$$V = V_0 + \Delta V^{1\text{-loop}}, \quad (4)$$

where we have included in $\Delta V^{1\text{-loop}}$ the contribution of the top-stop sector. The condition (3) now reads

$$\rho(Q) + \Delta\rho(Q) \leq 0, \quad (5)$$

with $M_{\text{SUSY}} < Q < Q_c$. It is easy to see that the quantity $\Delta\rho$ in (5) is related to the one-loop corrections $\Delta V^{1\text{-loop}}$ in (4) and is proportional to logarithms of ratios Q^2/M_{SUSY}^2 . The l.h.s. of (5), being the second derivative of (4) with respect of the field corresponding to the flat direction [3], depends on Q only at $\mathcal{O}(\hbar^2)$ and also through the wave function renormalization. As was discussed for instance in ref. [7], it can be safely considered scale independent if the scale Q is not too far from M_{SUSY} . Notice that, being the top mass

$m_t > 110$ GeV [8], the additional condition on m_t^2 which derives from requiring $v_L^2 > 0$ does not add any new significant restrictions (see the detailed discussions ref. [3] and [9] on this point).

The region of the parameter space where the condition (5) is implemented, hence leading to the spontaneous breaking of R parity, is shown in figures 1 and 2. We remind that in the MSSM there are five independent parameters in addition to those of the standard model. Once the electroweak radiative breaking is imposed this number decreases to four. In our case we choose these four independent parameters to be A , B (related to the trilinear and bilinear scalar terms of the SUSY soft breaking sector), m_0 and $m_{1/2}$, where m_0 denotes the common universal mass for scalars in the soft breaking sector of the theory and $m_{1/2}$ is the gluino mass (we are assuming a common gaugino mass as a boundary condition). In fig. 1 we show the region of the $m_0 - m_{1/2}$ plane in which the condition (5) is fulfilled (region II); we set $m_t = 150$ GeV and plot the two curves corresponding to $\rho + \Delta\rho = 0$ for $A = B = 0$ (curve a)), and $A = B + 1 = 3$ (curve b)). Fig. 2 is the same as fig. 1 but for $m_t = 190$ GeV.

Spontaneous breaking of R parity (and L number) through a non vanishing VEV of the sneutrino is nowadays phenomenologically unacceptable, since in a model with spontaneous breaking of L through the VEV of a scalar field carrying $SU(2)_L \otimes U(1)_Y$ quantum numbers the Z^0 boson can decay into the Goldstone boson of L and the real component of the scalar field, which has the mass of the order of the VEV causing the breaking of L . This decay would contribute to the Z^0 invisible decay width and is by now excluded by LEP measurements [4]. Moreover the VEV of the sneutrino is severely limited by the usual arguments of the rapid energy loss of red giants through majoron emission [10] ($\langle \tilde{\nu} \rangle < \mathcal{O}(1)$ MeV) and by the experimental bound on the tau neutrino mass, $m_\nu < 31$ MeV, roughly speaking ($\langle \tilde{\nu} \rangle < \mathcal{O}(10)$ GeV [11]). We conclude that in the MSSM the points in the parameter space driving a spontaneous breaking of L through $\langle \tilde{\nu} \rangle \neq 0$ must be excluded

(region II in figures 1 and 2). This is an additional constraint to impose in selecting the phenomenologically acceptable areas of the MSSM parameter space. Our results are in substantial agreement with previous analyses which were excluding the region in the SUSY parameter space where an eigenvalue of the sneutrino squared masses matrix was negative in the vacuum in which $\langle \tilde{\nu} \rangle = 0$ (see for instance the analysis of Ellis *et al.* in ref. [12]).

It has been argued recently that global symmetries are generally spoiled by gravitational effects [5]. Although we find the argument leading to this conclusion not so compelling and, in any case, certainly defensible only at a qualitative level, we think that it is anyway physically relevant to tackle the following pragmatical question concerning the fate of the lepton number symmetry in SUSY: what are the effects of a possibly very tiny explicit breaking of L showing up in a low energy SUSY lagrangian which would otherwise conserve L exactly?

Let us first state the framework for our discussion. We consider the MSSM (hence a model with conserved R matter parity) and, as an effect of some L violation taking place at some superlarge scale, we supplement the standard scalar potential of the MSSM with the term $m_\tau^2 \tilde{L} H_2$ where $\tilde{L} = (\tilde{\nu}, \tilde{l})$ denotes a leptonic scalar iso-doublet. We consider $m_\tau \ll m_W$. The presence of this L violating term in the scalar potential of the low energy theory is unavoidable if one assumes some L violation was produced in the superpotential of the theory at some superheavy (such as Planck or grand unification) scale. Indeed, even if one rotates the F term $L H_2$ away, the other L violating F terms which are generated give rise to $\tilde{L} H_2$ in the scalar potential through a renormalization effect when proceeding from the superheavy scale down to the low energy scale (for a thorough discussion of this point and similar considerations in grand unified SUSY theories see Hall and Suzuki in ref. [13]). Anyway, for the sake of our discussion, it is completely irrelevant how actually the term $m_\tau^2 \tilde{L} H_2$ arises in the effective low energy theory; suffice it to say that if L symmetry

is broken at the superheavy scale, the net effect in the low energy scalar potential would be the appearance of such a bilinear term. Higher dimensional operators are suppressed by powers of the superheavy mass and, in any case, are not relevant for the effect discussed below.

In the presence of the additional (tiny) explicit breaking of L the previously phenomenologically unwelcome majoron is no longer true Goldstone boson but rather it becomes a pseudo-Goldstone boson acquiring a mass which must vanish in the limit of $m_\tau^2 \rightarrow 0$. In fact, one could naively expect the mass of the pseudo majoron to be $\mathcal{O}(m_\tau^2)$. Hence for $m_\tau \ll m_W$ and $\langle \tilde{\nu} \rangle \ll m_W$ one would not avoid the above-mentioned problem of a too large contribution to the invisible width of the Z^0 by the decay of the Z^0 into the pseudo majoron and its scalar particle. We show that there exist general cases where this naive expectation is not fulfilled since the pseudo majoron can obtain a mass larger than m_{Z^0} even though the hierarchy $m_\tau \ll m_{Z^0}$ holds. Adding the term $m_\tau^2 \tilde{\nu} H_2^0$ to the usual scalar potential of the MSSM we obtain the following squared mass matrices in the (ν_1, ν_2, ν_L)

$$\text{basis} \quad \mathcal{M}_S^2 = \begin{pmatrix} m_{3\nu_1}^2 + \frac{g^2}{2} \nu_1^2 & -m_3^2 - \frac{g^2}{2} \nu_1 \nu_2 & \frac{g^2}{2} \nu_1 \nu_L \\ -m_3^2 - \frac{g^2}{2} \nu_1 \nu_2 & m_{3\nu_2}^2 + \frac{g^2}{2} \nu_2^2 + m_\tau^2 \frac{\nu_L}{\nu_2} & \frac{g^2}{2} \nu_2 \nu_L - m_\tau^2 \\ \frac{g^2}{2} \nu_1 \nu_L & \frac{g^2}{2} \nu_2 \nu_L - m_\tau^2 & \frac{g^2}{2} \nu_L^2 + m_\tau^2 \frac{\nu_2}{\nu_L} \end{pmatrix}, \quad (6)$$

and

$$\mathcal{M}_P^2 = \begin{pmatrix} m_{3\nu_1}^2 & m_3^2 & 0 \\ m_3^2 & m_{3\nu_2}^2 + m_\tau^2 \frac{\nu_L}{\nu_2} & m_\tau^2 \frac{\nu_2}{\nu_L} \\ 0 & m_\tau^2 \frac{\nu_L}{\nu_2} & m_\tau^2 \frac{\nu_2}{\nu_L} \end{pmatrix}, \quad (7)$$

for the CP-even and CP-odd neutral scalars, respectively.

In the limit $\nu_L \ll \nu_1, \nu_2$, the pseudo majoron J and its scalar partner acquire a mass of the same order of magnitude

$$m_J \simeq m_\tau^2 \frac{\nu_2}{\nu_L}. \quad (8)$$

From eq. (8) we see that m_J can be larger than $m_{Z^0}/2$ even though $m_\tau \ll m_{Z^0}$ provided that ν_L is small enough. Indeed, $\langle \tilde{\nu} \rangle$ has to be rather small since now the tau neutrino is

cosmologically stable due to the fact that with a heavy pseudo majoron it can no longer rapidly decay into a lighter neutrino plus a majoron. To avoid the overclosure of the Universe by ν_i 's, we ask for $m_{\nu_i} \lesssim 100$ eV [14] which yields

$$\langle \tilde{\nu} \rangle < \mathcal{O}(1) \text{ MeV}. \quad (9)$$

Taking v_2 of order of 100 GeV and asking (8) we see that for

$$m_i^2 > 10^{-1} \left(\frac{v_L}{1 \text{ MeV}} \right) \text{GeV}^2, \quad (10)$$

one obtains $m_J \gtrsim 100$ GeV. Given the bound (9) we conclude that in the worst case, $v_L = 1$ MeV, a value of m_i of order of 1 GeV is enough to provide a pseudomajoron mass large enough to prevent new invisible decays of the Z^0 boson. This represents the major point of our paper: once a hierarchy $v_L < m_i \ll m_{Z^0}$ is present the would-be majoron acquires a mass which can be easily larger than m_{Z^0} .

We now come to a brief discussions about the cosmological consequences of an explicit violation of lepton number. The presence of a term which violates L (and $B - L$) explicitly may cause severe problems to the survival of the matter-antimatter asymmetry in the Universe. It is well-known that the simultaneous presence in equilibrium of sphaleron mediated ($B + L$ violating) processes and R violating phenomena would lead to the complete wash out of any cosmic baryon asymmetry generated at the grand unification scale [15]. Imposing the out-of-equilibrium condition on L violating processes generated by the $m_i^2 \tilde{L} H_2$ term leads to the bound $m_i \lesssim 10^{-6}$ GeV [16]. Clearly in this case $\langle \tilde{\nu} \rangle$ should be extremely small ($\langle \tilde{\nu} \rangle \lesssim 10^{-5}$ eV) to obtain $m_J \gtrsim m_{Z^0}$ according to eq. (10). However, one can envisage different mechanisms to avoid such strong constraints. First, the term $m_i^2 \tilde{L} H_2$ need not violate all the three partial lepton numbers (at least by comparable amount). If, for instance, the L_e number is (almost) conserved the initial ΔB is not washed out [17]. However, in view of the possible flavour blindness of L violating gravitational effects, maybe one should focus on the much more appealing possibility of generating the baryon

asymmetry of the Universe during the electroweak phase transition in the framework of the MSSM [18, 19].

In conclusion we have shown that the presence of an explicit tiny breaking of L in the scalar potential may provide the would-be majoron a mass which is much larger than the typical size of such breaking. In particular, it is not difficult to obtain $m_J \gtrsim m_{Z^0}$ hence allowing for a non vanishing $\langle \tilde{\nu} \rangle$ without increasing the invisible width of the Z^0 boson.

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Figure Caption

Fig. 1 Lines corresponding to the condition $\rho + \Delta\rho = 0$ (see text, eq. (5)) in the $m_u - m_g$ plane. We set $m_t = 150$ GeV and $A = B = 0$ (line a)), $A = B + 1 = 3$ (line b)). Region II corresponds to $\rho + \Delta\rho < 0$.

Fig. 2 The same as in Fig.1) but with $m_t = 190$ GeV.

fig.I)

fig.II)