Proposal for a Supersymmetric Standard Model

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The fact that neutrinos are massive suggests that the minimal supersymmetric standard model (MSSM) might be extended in order to include three gauge-singlet neutrino superfields with Yukawa couplings of the type $H_2 L \nu^c$. We propose to use these superfields to solve the μ problem of the MSSM without having to introduce an extra singlet superfield as in the case of the next-to-MSSM (NMSSM). In particular, terms of the type $\nu^c H_1 H_2$ in the superpotential may carry out this task spontaneously through sneutrino vacuum expectation values. In addition, terms of the type $(\nu^c)^3$ avoid the presence of axions and generate effective Majorana masses for neutrinos at the electroweak scale. On the other hand, these terms break lepton number and R parity explicitly. For Dirac masses of the neutrinos of order 10^{-4} GeV, eigenvalues reproducing the correct scale of neutrino masses are obtained.

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Neutrino experiments have confirmed during the last years that neutrinos are massive [1]. As a consequence, all theoretical models must be modified in order to reproduce this result. In particular, it is natural in the context of the minimal supersymmetric standard model (MSSM) [2] to supplement the ordinary neutrino superfields, $\hat{\nu}_i$, $i = 1, 2, 3$ contained in the $SU(2)_L$ doublet, $\hat{L}_i$, with gauge-singlet neutrino superfields, $\hat{\nu}^c_i$. Once experiments induce us to introduce these new superfields, and given the fact that sneutrinos are allowed to get vacuum expectation values (VEVs), we may wonder why not to use terms of the type $\hat{\nu}^c \hat{H}_1 \hat{H}_2$ to produce an effective μ term. This would allow us to solve the naturalness problem of the MSSM, the so-called μ problem [3], without having to introduce an extra singlet superfield as in case of the next-to-minimal supersymmetric standard model (NMSSM) [4]. It is true that in the model with bilinear R-parity violation (BRpV) [5], the bilinear terms $\hat{H}_2 \hat{L}_i$ induce neutrino masses through the mixing with the neutralinos (actually only one mass at tree level and the other two at one loop) without using the superfields $\hat{\nu}^c_i$; however, the μ problem is augmented with the three new bilinear terms.

Thus the aim of this Letter is to analyze the "μ from ν" supersymmetric standard model ($\mu\nu$SSM) arising from this proposal: natural particle content without μ problem.

In addition to the MSSM Yukawa couplings for quarks and charged leptons, the $\mu\nu$SSM superpotential contains Yukawa couplings for neutrinos, and two additional type of terms involving the Higgs doublet superfields, $\hat{H}_1$ and $\hat{H}_2$, and the three neutrino superfields, $\hat{\nu}^c_i$,

$$
\begin{aligned}
W = \epsilon_{ab}(Y_u^{ij} \hat{H}_2^b \hat{Q}_i^a \hat{u}_j^c + Y_d^{ij} \hat{H}_1^a \hat{Q}_i^b \hat{d}_j^c + Y_e^{ij} \hat{H}_1^a \hat{L}_i^b \hat{e}_j^c \\
+ Y_\nu^{ij} \hat{H}_2^b \hat{L}_i^a \hat{\nu}_j^c) - \epsilon_{ab} \lambda^i \hat{\nu}_i^c \hat{H}_1^a \hat{H}_2^b + \tfrac{1}{3} \kappa^{ijk} \hat{\nu}_i^c \hat{\nu}_j^c \hat{\nu}_k^c, \quad (1)
\end{aligned}
$$

where we take $\hat{H}_1^T = (\hat{H}_1^0, \hat{H}_1^-)$, $\hat{H}_2^T = (\hat{H}_2^+, \hat{H}_2^0)$, $\hat{Q}_i^T = (\hat{u}_i, \hat{d}_i)$, $\hat{L}_i^T = (\hat{\nu}_i, \hat{e}_i)$, a, b are $SU(2)$ indices, and $\epsilon_{12} = 1$. In this model, the usual MSSM bilinear μ term is absent from the superpotential, and only dimensionless trilinear couplings are present in W. For this to happen we can invoke a Z_3 symmetry as is usually done in the NMSSM. On the other hand, let us recall that this is actually what happens in the low-energy limit of string constructions: only trilinear couplings are present in the superpotential. Since string theory seems to be relevant for the unification of interactions, including gravity, this argument in favor of the absence of a bare μ term in the superpotential is robust. When the scalar components of the superfields $\hat{\nu}^c_i$, denoted by $\tilde{\nu}^c_i$, acquire VEVs of order the electroweak scale, an effective interaction $\mu \hat{H}_1 \hat{H}_2$ is generated through the fifth term in (1), with $\mu \equiv \lambda^i \langle \tilde{\nu}^c_i \rangle$. The last type of terms in (1) is allowed by all symmetries, and avoids the presence of an unacceptable axion associated to a global $U(1)$ symmetry. In addition, it generates effective Majorana masses for neutrinos at the electroweak scale. These two types of terms replace the two NMSSM terms $\hat{S}\hat{H}_1\hat{H}_2$, $\hat{S}\hat{S}\hat{S}$, with $\hat{S}$ an extra singlet superfield.

It is worth noticing that these terms break explicitly lepton number, and therefore, after spontaneous symmetry breaking, a massless Goldstone boson (Majoron) does not appear. On the other hand, R parity ($+1$ for particles and -1 for superpartners) is also explicitly broken and this means that the phenomenology of the $\mu\nu$SSM is going to be very different from the one of the MSSM. Needless to mention, the lightest R-odd particle is not stable. Obviously, the neutralino is no longer a candidate for dark matter. Nevertheless, other candidates can be found in the literature, such as the gravitino [6], the well-known axion, and many other (exotic) particles [7]. It is also interesting to realize that the Yukawa couplings producing Dirac masses for neutrinos, the fourth term in (1), generate

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through the VEVs of $\tilde{\nu}_i^c$, three effective bilinear terms $\hat{H}_2\hat{L}_i$. As mentioned above these characterize the BRpV.

Let us finally remark that the superpotential (1) has a Z_3 symmetry, just like the NMSSM. Therefore, one expects to have also a cosmological domain wall problem [8,9] in this model. Nevertheless, the usual solutions to this problem [10] will also work in this case: nonrenormalizable operators [8] in the superpotential can break explicitly the dangerous Z_3 symmetry, lifting the degeneracy of the three original vacua, and this can be done without introducing hierarchy problems. In addition, these operators can be chosen small enough as not to alter the low-energy phenomenology.

Working in the framework of gravity mediated supersymmetry breaking, we will discuss now in more detail the phenomenology of the $\mu\nu$SSM. Let us write first the soft terms appearing in the Lagrangian $\mathcal{L}_{\rm soft}$, which in our conventions is given by

$$\begin{aligned}-\mathcal{L}_{\rm soft} = {}& (m^2_{\tilde{Q}})^{ij}\tilde{Q}_i^{a*}\tilde{Q}_j^a + (m^2_{\tilde{u}^c})^{ij}\tilde{u}_i^{c*}\tilde{u}_j^c + (m^2_{\tilde{d}^c})^{ij}\tilde{d}_i^{c*}\tilde{d}_j^c + (m^2_{\tilde{L}})^{ij}\tilde{L}_i^{a*}\tilde{L}_j^a + (m^2_{\tilde{e}^c})^{ij}\tilde{e}_i^{c*}\tilde{e}_j^c + m^2_{H_1}H_1^{a*}H_1^a + m^2_{H_2}H_2^{a*}H_2^a \\ & + (m^2_{\tilde{\nu}^c})^{ij}\tilde{\nu}_i^{c*}\tilde{\nu}_j^c + \epsilon_{ab}[(A_uY_u)^{ij}H_2^b\tilde{Q}_i^a\tilde{u}_j^c + (A_dY_d)^{ij}H_1^a\tilde{Q}_i^b\tilde{d}_j^c + (A_eY_e)^{ij}H_1^a\tilde{L}_i^b\tilde{e}_j^c + (A_\nu Y_\nu)^{ij}H_2^b\tilde{L}_i^a\tilde{\nu}_j^c + \text{H.c.}] \\ & + [-\epsilon_{ab}(A_\lambda\lambda)^i\tilde{\nu}_i^cH_1^aH_2^b + \tfrac{1}{3}(A_\kappa\kappa)^{ijk}\tilde{\nu}_i^c\tilde{\nu}_j^c\tilde{\nu}_k^c + \text{H.c.}] - \tfrac{1}{2}(M_3\tilde{\lambda}_3\tilde{\lambda}_3 + M_2\tilde{\lambda}_2\tilde{\lambda}_2 + M_1\tilde{\lambda}_1\tilde{\lambda}_1 + \text{H.c.}). \end{aligned} \quad (2)$$

In addition to terms from $\mathcal{L}_{\rm soft}$, the tree-level scalar potential receives the usual D and F term contributions. Once the electroweak symmetry is spontaneously broken, the neutral scalars develop in general the following VEVs:

$$\langle H_1^0\rangle = v_1, \qquad \langle H_2^0\rangle = v_2, \qquad \langle\tilde{\nu}_i\rangle = \nu_i, \qquad \langle\tilde{\nu}_i^c\rangle = \nu_i^c. \quad (3)$$

In what follows it will be enough for our purposes to neglect mixing between generations in (1) and (2), and to assume that only one generation of sneutrinos gets VEVs, ν, ν^c. The extension of the analysis to all generations is straightforward, and the conclusions are similar. We then obtain for the tree-level neutral scalar potential:

$$\begin{aligned}\langle V_{\rm neutral}\rangle = {}& \frac{g_1^2+g_2^2}{8}(|\nu|^2+|v_1|^2-|v_2|^2)^2 + |\lambda|^2(|\nu^c|^2|v_1|^2+|\nu^c|^2|v_2|^2+|v_1|^2|v_2|^2) + |\kappa|^2|\nu^c|^4 + |Y_\nu|^2(|\nu^c|^2|v_2|^2 \\ & + |\nu^c|^2|\nu|^2 + |\nu|^2|v_2|^2) + m^2_{H_1}|v_1|^2 + m^2_{H_2}|v_2|^2 + m^2_{\tilde{\nu}^c}|\nu^c|^2 + m^2_{\tilde{\nu}}|\nu|^2 + (-\lambda\kappa^* v_1v_2\nu^{c*2} - \lambda Y_\nu^*|\nu^c|^2v_1\nu^* \\ & - \lambda Y_\nu^*|v_2|^2v_1\nu^* + kY_\nu^*v_2^*\nu^*\nu^{c2} - \lambda A_\lambda\nu^c v_1v_2 + Y_\nu A_\nu\nu^c\nu v_2 + \tfrac{1}{3}\kappa A_\kappa\nu^{c3} + \text{H.c.}). \end{aligned} \quad (4)$$

In the following, we assume for simplicity that all parameters in the potential are real. One can derive the four minimization conditions with respect to the VEVs v_1, v_2, ν^c, ν, with the result

$$\begin{aligned} & \tfrac{1}{4}(g_1^2+g_2^2)(\nu^2+v_1^2-v_2^2)v_1 + \lambda^2v_1(\nu^{c2}+v_2^2) \\ & \quad + m^2_{H_1}v_1 - \lambda\nu^cv_2(\kappa\nu^c+A_\lambda) - \lambda Y_\nu\nu(\nu^{c2}+v_2^2) = 0, \\ & -\tfrac{1}{4}(g_1^2+g_2^2)(\nu^2+v_1^2-v_2^2)v_2 + \lambda^2v_2(\nu^{c2}+v_1^2) \\ & + m^2_{H_2}v_2 - \lambda\nu^cv_1(\kappa\nu^c+A_\lambda) + Y_\nu^2v_2(\nu^{c2}+\nu^2) \\ & \quad + Y_\nu\nu(-2\lambda v_1v_2 + \kappa\nu^{c2} + A_\nu\nu^c) = 0, \\ & \lambda^2(v_1^2+v_2^2)\nu^c + 2\kappa^2\nu^{c3} + m^2_{\tilde{\nu}^c}\nu^c - 2\lambda\kappa v_1v_2\nu^c \\ & - \lambda A_\lambda v_1v_2 + \kappa A_\kappa\nu^{c2} + Y_\nu^2\nu^c(v_2^2+\nu^2) \\ & \quad + Y_\nu\nu(-2\lambda\nu^cv_1 + 2\kappa v_2\nu^c + A_\nu v_2) = 0, \\ & \tfrac{1}{4}(g_1^2+g_2^2)(\nu^2+v_1^2-v_2^2)\nu + m^2_{\tilde{\nu}}\nu + Y_\nu^2\nu(v_2^2+\nu^{c2}) \\ & \quad + Y_\nu(-\lambda\nu^{c2}v_1 - \lambda v_2^2v_1 + \kappa v_2\nu^{c2} + A_\nu\nu^cv_2) = 0. \end{aligned} \quad (5)$$

As discussed in the context of R-parity breaking models with extra singlets [11], the VEV of the left-handed sneutrino, ν, is in general small. Here we can use the same argument. Notice that in the last equation in (5) $\nu \to 0$ as $Y_\nu \to 0$, and since the coupling Y_ν determines the Dirac mass for the neutrinos, $m_D \equiv Y_\nu v_2$, ν has to be very small. Using this rough argument we can also get an estimate of the value, $\nu \lesssim m_D$. This also implies that we can approximate the other three equations as follows:

$$\begin{aligned} & \tfrac{1}{2}M_Z^2\cos2\beta + \lambda^2(\nu^{c2}+v^2\sin^2\beta) \\ & \quad + m^2_{H_1} - \lambda\nu^c\tan\beta(\kappa\nu^c+A_\lambda) = 0, \\ & -\tfrac{1}{2}M_Z^2\cos2\beta + \lambda^2(\nu^{c2}+v^2\cos^2\beta) \\ & \quad + m^2_{H_2} - \lambda\nu^c\cot\beta(\kappa\nu^c+A_\lambda) = 0, \\ & \lambda^2v^2 + 2\kappa^2\nu^{c2} + m^2_{\tilde{\nu}^c} - \lambda\kappa v^2\sin2\beta \\ & \quad - \frac{\lambda A_\lambda v^2}{2\nu^c}\sin2\beta + \kappa A_\kappa\nu^c = 0, \end{aligned} \quad (6)$$

where $\tan\beta \equiv v_2/v_1$, $2M_W^2/g_2^2 = v_1^2+v_2^2+\nu^2 \approx v_1^2+v_2^2 \equiv v^2$, and we have neglected terms proportional to Y_ν. It is worth noticing that these equations are the same as the ones defining the minimization conditions for the NMSSM, with the substitution $\nu^c \leftrightarrow s$. Thus one can carry out the analysis of the model similarly to the NMSSM case, where many solutions in the parameter space λ, κ, $\mu(\equiv \lambda s)$, $\tan\beta$, A_λ, A_κ, can be found [12].

Once we know that solutions are available in this model, we have to discuss in some detail the important issue of mass matrices. Concerning this point, the breaking of R parity makes the $\mu\nu$SSM very different from the MSSM and the NMSSM. In particular, neutral gauginos and

Higgsinos are now mixed with the neutrinos. Not only the fermionic component of $\tilde{\nu}^c$ mixes with the neutral Higgsinos (similarly to the fermionic component of S in the NMSSM), but also the fermionic component of $\tilde{\nu}$ enters in the game, giving rise to a sixth state. Of course, now we have to be sure that one eigenvalue of this matrix is very small, reproducing the experimental results about neutrino masses. In the weak interaction basis defined by $\Psi^{0T} \equiv (\tilde{B}^0 = -i\tilde{\lambda}', \tilde{W}_3^0 = -i\tilde{\lambda}_3, \tilde{H}_1^0, \tilde{H}_2^0, \nu^c, \nu)$, the neutral fermion mass terms in the Lagrangian are $\mathcal{L}_{\text{neutral}}^{\text{mass}} = -\frac{1}{2} \times (\Psi^0)^T \mathcal{M}_{\text{n}} \Psi^0 + \text{H.c.}$, with $\mathcal{M}_{\text{n}}$ a 6×6 (10×10 if we include all generations of neutrinos) matrix,

$$\mathcal{M}_{\text{n}} = \begin{pmatrix} M & m \\ m^T & 0 \end{pmatrix}, \tag{7}$$

where

$$M = \begin{pmatrix} M_1 & 0 & -M_Z \sin\theta_W \cos\beta & M_Z \sin\theta_W \sin\beta & 0 \\ 0 & M_2 & M_Z \cos\theta_W \cos\beta & -M_Z \cos\theta_W \sin\beta & 0 \\ -M_Z \sin\theta_W \cos\beta & M_Z \cos\theta_W \cos\beta & 0 & -\lambda\nu^c & -\lambda v_2 \\ M_Z \sin\theta_W \sin\beta & -M_Z \cos\theta_W \sin\beta & -\lambda\nu^c & 0 & -\lambda v_1 + Y_\nu \nu \\ 0 & 0 & -\lambda v_2 & -\lambda v_1 + Y_\nu \nu & 2\kappa\nu^c \end{pmatrix}, \tag{8}$$

is very similar to the neutralino mass matrix of the NMSSM (substituting $\nu^c \leftrightarrow s$ and neglecting the contributions $Y_\nu \nu$), and

$$m^T = \left(-\frac{g_1 \nu}{\sqrt{2}} \frac{g_2 \nu}{\sqrt{2}} 0 Y_\nu \nu^c Y_\nu v_2 \right). \tag{9}$$

Matrix (7) is a matrix of the seesaw type that will give rise to a very light eigenvalue if the entries of the matrix M are much larger than the entries of the matrix m. This is generically the case since the entries of M are of order the electroweak scale, but for the entries of m, ν is small and $Y_\nu v_2$ is the Dirac mass for the neutrinos m_D as discussed above ($Y_\nu \nu^c$ has the same order of magnitude of m_D). We have checked numerically that correct neutrino masses can easily be obtained. For example, using typical electroweak-scale values in (8), and a Dirac mass of order 10^{-4} GeV in (9), one obtains that the lightest eigenvalue of (7) is of order 10^{-2} eV. Including the three generations in the analysis we can obtain different neutrino mass hierarchies playing with the hierarchies in the Dirac masses.

The possibility of using a seesaw at the electroweak scale has not been considered in much detail in the literature [For a recent work see Ref. [13], where an extension of the NMSSM is considered with Majorana masses for neutrinos generated dynamically through the VEV of the singlet S. R parity may be broken in this extension, although spontaneously], although this avoids the introduction of *ad-hoc* high energy scales. Of course, with a seesaw at the scale of a grand unified theory (GUT), one can have Yukawa couplings of order one for neutrinos. However, since we know that the Yukawa coupling of the electron has to be of order 10^{-6}, why should the one of the neutrino be 6 orders of magnitude larger? As mentioned above, with the electroweak-scale seesaw a Yukawa coupling of order of the one of the electron is sufficient to reproduce the neutrino mass. Notice also that a purely Dirac mass for the neutrino would imply a Yukawa coupling of order 10^{-13}, i.e., 7 orders of magnitude smaller than the one we need with a electroweak-scale seesaw. It is worth mentioning here that in some string constructions, where supersymmetric standardlike models can be obtained without the necessity of a GUT, and Yukawa couplings can be explicitly computed, those for neutrinos cannot be as small as 10^{-13}, and therefore the presence of a seesaw at the electroweak scale is helpful [14]. In any case, let us remark that in our model the seesaw is dynamical and unavoidable, since the matrix of Eq. (7) producing such a seesaw is always present.

It has been noted in the literature that the sneutrino-antisneutrino mixing effect generates a loop correction to the neutrino mass, which depends on the mass-splitting of the sneutrino mass eigenstates [15]. In the case of assuming a large Majorana mass this correction is negligible if all parameters are of order the supersymmetric scale. We have checked that the same result is obtained in our model with a seesaw at the electroweak scale, unless a fine tune of the parameters is forced producing a too large sneutrino mass difference.

On the other hand, the charginos mix with the charged leptons and therefore in a basis where $\Psi^{+T} \equiv (-i\tilde{\lambda}^+, \tilde{H}_2^+, e_R^+)$ and $\Psi^{-T} \equiv (-i\tilde{\lambda}^-, \tilde{H}_1^-, e_L^-)$, one obtains the matrix

$$\begin{pmatrix} M_2 & g_2 v_2 & 0 \\ g_2 v_1 & \lambda\nu^c & -Y_e \nu \\ g_2 \nu & -Y_\nu \nu^c & Y_e v_1 \end{pmatrix}. \tag{10}$$

Here we can distinguish the 2×2 submatrix which is similar to the chargino mass matrix of the NMSSM (substituting $\nu^c \leftrightarrow s$). Clearly, given the vanishing value of the 13 element of the matrix (10), and the extremely small absolute value of the 23 element, there will always be a light eigenvalue corresponding to the electron mass $Y_e v_1$. The extension of the analysis to three generations is again straightforward.

Of course, other mass matrices are also modified. This is the case for example of the Higgs boson mass matrices. The presence of the VEVs ν, ν^c, leads to mixing of the neutral Higgs bosons with the sneutrinos. Concerning the

Higgs phenomenology, since basically the ν^c plays the role of the singlet S, this will be similar to the one of the NMSSM [12]. For example, two CP-odd Higgs bosons are present, and we have checked that one of them can in principle be light. Likewise the charged Higgs bosons will be mixed with the charged sleptons. On the other hand, when compared to the MSSM case, the structure of squark mass terms is essentially unaffected, provided that one uses $\mu = \lambda \nu^c$, and neglects the contribution of the fourth term in (1).

Obviously, the phenomenology of the $\mu\nu$SSM is very rich and different from other models, and therefore many more issues might have been addressed, such as possible experimental constraints, implications for accelerator physics, analysis of the (modified) renormalization group equations, study of the neutrino masses in detail, etc. However, these are beyond the scope of this Letter, and we leave this necessary task for a future work [16]. Our main interest here was to introduce the characteristics of this new model, and sketch some important points concerning its phenomenology.

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