Minimal Nonminimal Supersymmetric Standard Model

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Abstract. We review the basic field-theoretic and phenomenological features of the recently introduced Minimal Nonminimal Supersymmetric Standard Model (MNSSM). The introduced model is the simplest and most economic version among the proposed nonminimal supersymmetric models, in which the so-called $\mu$-problem can be successfully addressed. As opposed to the MSSM and the frequently-discussed NMSSM, the MNSSM can naturally predict the existence of a light charged Higgs boson with a mass smaller than 100 GeV. Such a possible realization of the Higgs sector can be soon be tested at the upgraded Run II phase of the Tevatron collider.

It is known that Minimal Supersymmetric Standard Model (MSSM) suffers from the so called $\mu$-problem. The superpotential of the MSSM contains a bilinear term $-\mu \hat{H}_1 \hat{H}_2$ involving the two Higgs-doublet superfields $\hat{H}_1$ and $\hat{H}_2$, known as the $\mu$-term. Naive implementation of the $\mu$-parameter within supergravity theories would lead to a $\mu$ value of the order of the Planck scale $M_P$. However, for a successful Higgs mechanism at the electroweak scale, the $\mu$-parameter is actually required to be many orders of magnitude smaller of order $M_{\text{SUSY}}$. Many scenarios, all based on extensions of the MSSM, have been proposed in the existing literature \cite{1} to provide a natural explanation for the origin of the $\mu$-term.

Recently, a minimal extension of the MSSM has been presented\cite{2,3,4,5}, called the Minimal Nonminimal Supersymmetric Standard Model (MNSSM)\cite{3,5}, in which the $\mu$-problem can be successfully addressed in a rather minimal way. In the MNSSM the $\mu$-parameter is promoted to a chiral singlet superfield $\hat{S}$, and all linear, quadratic and cubic operators involving only $\hat{S}$ are absent from the renormalizable superpotential; $\hat{S}$ enters through the single term $\lambda \hat{S} \hat{H}_1 \hat{H}_2$: \begin{equation} W_{\text{MSSM}} = \tilde{W}_{\text{MSSM}} + \lambda \hat{S} \hat{H}_1^T i\tau_2 \hat{H}_2, \end{equation}

where $\tilde{W}_{\text{MSSM}}$ is the superpotential of the MSSM without the presence of the $\mu$ term. The crucial difference between the MNSSM and the frequently-discussed Next-to-Minimal Supersymmetric Standard Model (NMSSM) \cite{3} lies in the fact that the cubic term $\frac{1}{3} \kappa \hat{S}^3$ does not appear in the renormalizable superpotential of the former.

The key point in the construction of the renormalizable MNSSM superpotential is that the simple form (1) may be enforced by discrete $R$-symmetries, such as $Z_R^5$ \cite{2,3,4,5} and $Z_R^7$ \cite{3,5}. These discrete $R$-symmetries, however, must be extended to the gravity-induced non-renormalizable superpotential and Kähler potential terms as well. To communicate

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the breaking of supersymmetry (SUSY), we consider the scenario of $N = 1$ supergravity spontaneously broken by a set of hidden-sector fields at an intermediate scale. Within this framework of SUSY-breaking, we have then been able to show [3] that the above $R$-symmetries are sufficient to postpone the appearance of the potentially dangerous tadpole [7].

$\kappa \neq 0$, $\mu \neq 0$ $\kappa \neq 0$, $\mu \neq 0$

$M_{\text{SUSY}}^{2} S S^* \sim M_{\text{SUSY}}^{2} S S^* S$ lead to a vacuum expectation value (VEV) for $S$, $\langle S \rangle = \frac{1}{\sqrt{2}} v_S$, of order $M_{\text{SUSY}} \sim 1$ TeV. The latter gives rise to a $\mu$-parameter at the required electroweak scale, i.e.

$$\mu = -\frac{1}{\sqrt{2}} \lambda v_S \sim M_{\text{SUSY}}.$$  

Thus, a natural explanation for the origin of the $\mu$-parameter can be obtained. Finally, since the effective tadpole term $t_S S$ explicitly breaks the continuous Peccei–Quinn symmetry governing the remaining renormalizable Lagrangian of the MNSSM, the theory naturally avoids the presence of a phenomenologically excluded weak-scale axion.

In addition to the tadpole $t_S$ of the physical scalar $S$, an effective tadpole for its auxiliary component $F_S$ is generated [8]. However, depending on the underlying mechanism of SUSY breaking, the effective tadpole proportional to $F_S$ could in principle be absent from the model. Such a reduction of the renormalizable operators does not thwart the renormalizability of the theory. The resulting renormalizable low-energy scenario has one parameter less than the frequently-discussed NMSSM with the cubic singlet-superfield term $\hat{S}^3$ present; it therefore represents the most economic, renormalizable version among the non-minimal supersymmetric models proposed in the literature.

As opposed to the NMSSM, the MNSSM satisfies the tree-level mass sum rule [3]:

$$M_{H_1}^2 + M_{H_2}^2 + M_{A_1}^2 = M_{Z}^2 + M_{A_2}^2,$$  

where $H_{1,2,3}$ and $A_{1,2}$ are the three CP-even and two CP-odd Higgs fields, respectively. The tree-level mass sum rule (4) is very analogous to the corresponding one of the MSSM [9], where the two heavier Higgs states $H_3$ and $A_2$ are absent in the latter. This striking analogy to the MSSM allows us to advocate that the Higgs sector of the MNSSM differs indeed minimally from the one of the MSSM, i.e. the introduced model truly constitutes the minimal supersymmetric extension of the MSSM. In the NMSSM, the violation of the mass sum rule (4) can become much larger than the one induced by the one-loop stop/top effects, especially for relatively large values of $|\kappa|$, $|\mu|$ and $|A_\kappa|$.

In the non-minimal supersymmetric standard models, the upper bound on the lightest CP-even Higgs-boson mass $M_{H_1}$ has a tree-level dependence on the coupling $\lambda$ [3, 10, 3, 4], i.e.

$$M_{H_1}^{2(0)} \leq M_{Z}^2 \left( \cos^2 2\beta + \frac{2 \lambda^2}{g_w^2 + g'^2} \sin^2 2\beta \right),$$  

where the angle $\beta$ is defined by means of $\tan \beta = v_2/v_1$, the ratio of the VEVs of the two Higgs doublets. Since in the MNSSM $\lambda$ can take its maximum allowed value naturally corresponding to the NMSSM with $\kappa = 0$ [10], the value of $M_{H_1}$ is predicted to be the highest. In particular, a renormalization-group-improved analysis [8] of the effective MNSSM Higgs potential leads to the upper bound: $M_{H_1} \lesssim 145$ GeV, for large stop mixing (see also Fig. 1).
Consequently, such a scenario can only be decisively tested by the upgraded Run II phase of the Tevatron collider at Fermilab and by the Large Hadron Collider (LHC) at CERN.

The MNSSM can comfortably predict viable scenarios, where the mass of the charged Higgs boson $H^+$ is in the range: $80 \text{ GeV} < M_{H^+} \lesssim 3 \text{ TeV}$, for phenomenologically relevant values of $|\mu| \gtrsim 100 \text{ GeV}$ [11]. In fact, as can be seen from Fig. 2, there is an absolute upper

Figure 1. Numerical values for $M_{H_1}$ versus $\mu$ in the MNSSM with $m_{T_2}^2 = 0$, for $M_{H^+} = 0.1$ (solid), 0.3 (dashed), 0.7 (dotted) and 1 (dash-dotted) TeV.
bound on $M_{H^+}$ for fixed values of $\lambda$ and $\tan \beta$.

On the other hand, charged Higgs-boson masses smaller than 100 GeV can naturally be obtained within the MNSSM, while the SM-like Higgs boson $H_{SM}$, with dominant coupling to the $Z$ boson, can be heavier than about 115 GeV \cite{3,5}. Instead, both in the MSSM and the NMSSM \cite{5}, such a Higgs-boson mass spectrum is theoretically inaccessible, if the phenomenologically favoured range $|\mu| \gtrsim 100$ GeV is considered. In Figs. 3 and 4, we display numerical values for the masses of the lightest and next-to-lightest Higgs bosons, $H_1$ and $H_2$, and their couplings to the $Z$ boson as functions of $\mu$, for a number of versions of the MNSSM that predict light charged Higgs bosons. In the MNSSM versions under study, the SM-like Higgs boson $H_{SM}$ (mainly $H_2$) can have a mass larger 110 GeV, compatible with the present experimental bound. The generic prediction is that the first CP-even Higgs boson $H_1$ is lighter than $H_2$ and has a suppressed coupling to the $Z$ boson in agreement with LEP2 data. From Figs. 3 and 4, we observe that the charged Higgs boson can be as light as the present experimental upper bound, i.e. $M_{H^+} \sim 80$ GeV. This is an important phenomenological feature of the MNSSM, which is very helpful to discriminate it from the NMSSM. It is a reflection of a new non-trivial decoupling limit due to a large tadpole $|t_s|$, which is only attainable in the MNSSM \cite{3}. In this limit, the heavier Higgs states $H_3$ and $A_3$ can both decouple from the Higgs spectrum as a heavy singlets. The upcoming upgraded Run II phase of the Tevatron collider has the physics potential to probe the viability of a
light-charged-Higgs-boson realization.

For scenarios with $M_{H^+} \gtrsim 200$ GeV, the distinction between the MNSSM and the NMSSM becomes more difficult. In this case, additional experimental information would be necessary to distinguish the two SUSY extensions of the MSSM, resulting from a precise determination of the masses, the widths, the branching ratios and the production cross sections.
of the CP-even and CP-odd Higgs bosons. Nevertheless, if the tadpole parameter $\lambda t_S/\mu$ becomes much larger than $M_{H^+}^2$, with the remaining kinematic parameters held fixed, the Higgs states $H_3$ and $A_2$ will be predominantly singlets. As an important phenomenological consequence of this, the complementarity relations between the $H_{1,2}ZZ$- and $H_{2,1}A_1Z$-
couplings will then hold approximately true in the MNSSM, i.e.
\[ g_{H_1ZZ}^2 = g_{H_2A_1Z}^2, \quad g_{H_2ZZ}^2 = g_{H_1A_1Z}^2. \]  \hspace{1cm} (6)
In addition, the couplings of the two heaviest states \( H_3 \) and \( A_2 \) to the gauge bosons will vanish. Here, we should stress that the relations (6) are not generically valid in the NMSSM. The latter is a consequence of the absence of the aforementioned large tadpole decoupling limit, such that the states \( H_3 \) and \( A_2 \) could decouple as singlets. Future next linear \( e^+e^- \) colliders have the capabilities to experimentally determine the \( H_{1,2}ZZ \)- and \( H_{2,1}A_1Z \)-couplings to an accuracy even up to 3% and so test, to a high degree, the complementarity relations (6) which are an essential phenomenological feature of the MNSSM.

As has been discussed in [3], the MNSSM also predicts the existence of a light neutralino, the axino. The axino is predominantly a singlet field, for \( |\mu| \gtrsim 120 \) GeV. LEP limits on the \( Z \)-boson invisible width lead to the additional constraint: \( 200 \lesssim |\mu| \lesssim 250 \) GeV, for \( \lambda \approx 0.65 \). However, such a constraint disappears completely for smaller values of \( \lambda \), namely for \( \lambda \lesssim 0.45 \). In fact, the axino may become the lightest supersymmetric particle, which is very long-lived in the MNSSM, and hence it potentially qualifies as a candidate for cold dark matter. We feel that a dedicated study in this direction needs to be done.

Let us summarize the basic field-theoretic and phenomenological features of the MNSSM: (i) The MNSSM minimally departs from the MSSM through the presence of a gauge-singlet superfield whose all self-couplings are absent. On the basis of discrete \( R \) symmetries, such as \( Z_5^R \) and \( Z_7^R \), the quadratically divergent harmful tadpoles first appear at the 6- and 7-loop levels, thereby avoiding to destabilize the gauge hierarchy. By the same token, the MNSSM can minimally account for the origin of \( \mu \)-term; (ii) Since the loop-induced tadpoles break any continuous or discrete symmetry, the model does not suffer from problems [12] related to visible weak-scale axions and domain walls; (iii) As a consequence of a new decoupling limit due to a large tadpole, the MNSSM can naturally predict viable scenarios in which the charged Higgs boson \( H^+ \) is much lighter than the neutral Higgs boson with a SM-type coupling to the \( Z \) boson. The planned colliders, i.e. the upgraded Tevatron collider [3] and the LHC, have the potential capabilities to test such interesting scenarios with a relatively light \( H^+ \); (iv) Unlike the frequently-discussed NMSSM, the Higgs sector of the MNSSM exhibits a much closer resemblance to the one of the MSSM, by means of the tree-level mass sum rule (4) and the complementarity relations (6) of the Higgs-boson couplings to the \( Z \) boson.

In conclusion, all the above facts point to a single perspective: the only truly minimal supersymmetric extension of the MSSM is the Minimal Nonminimal Supersymmetric Standard Model.

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