DESY 11-061

arXiv:1104.2695v1 [hep-ph] 14 Apr 2011

Generating μ and $B\mu$ in models with Dirac Gauginos

Karim Benakli♣*, Mark D. Goodsell◇† and Ann-Kathrin Maier♠‡

♣*Laboratoire de Physique Théorique et Hautes Energies, CNRS, UPMC Univ Paris 06 Boite 126, 4 Place Jussieu, 75252 Paris cedex 05, France*

◇*Deutsches Elektronen-Synchrotron, DESY, Notkestraße 85, 22607 Hamburg, Germany*

♠ *Laboratoire de physique de la matière complexe, SB-EPFL, CH-1015 Lausanne, Switzerland*

Abstract

We consider the extension of the Minimal Supersymmetric Standard Model by Dirac masses for the gauginos. We study the possibility that the same singlet $\mathbf{S}$ that pairs up with the bino, to form a Dirac fermion, is used to generate μ and $B\mu$ terms through its vacuum expectation value. For this purpose, we assume that, in the Higgs potential, the necessary R-symmetry breaking originates entirely from a superpotential term $\frac{\kappa}{3}\mathbf{S}^3$ and discuss the implications for the spectrum of the model.

*kbenakli@lpthe.jussieu.fr

†mark.goodsell@desy.de

‡ann-kathrin.maier@epfl.ch

1 Introduction

The supersymmetric extension of the standard model introduces new charged particles that need to a acquire a large mass in order to explain the absence of evidence in present collider experiments.In particular, the predicted gauginos are fermionic states that can obtain (after supersymmetry breaking) either Majorana or Dirac masses. Here we are interested in this latter case [1–27].

An important feature in models with Dirac gauginos is the fate of R-symmetry. In the global supersymmetric models considered here, it appears as a continuous symmetry (which can be broken to a discrete subgroup). It can not be spontaneously broken at the electroweak scale with a generic vacuum expectation value (vev), as this would lead to a massless R-axion with a coupling insufficently suppressed to have evaded early discovery. There remain two options: either it is conserved or explicitly broken. In order to quantify the required size of R-symmetry breaking, one needs to identify the minimal set of operators that violate the symmetry.

First, it is usual to consider that R-symmetry is broken in the Minimal Supersymmetric extension of the Standard Model (MSSM) by Majorana gaugino masses. We can however use instead Dirac masses for the gauginos, pairing them with additional states in adjoint representations (henceforth DG-adjoints): a singlet $\mathbf{S}$ for $U(1)_Y$, a triplet $\mathbf{T}$ for $SU(2)_w$ and an octet $\mathbf{O_g}$ for $SU(3)_c$.

Second, the gravitino mass required in flat space-time breaks R-symmetry. Again, this can be avoided by taking a Dirac mass for the gravitino. Such masses require the gravitational multiplet to be in $N = 2$ representations. To illustrate such a scenario, consider that the $N = 1$ gauge and matter fields appear on 3-branes. These are localised in a bulk having one flat extra dimension of radius R. Then a Dirac gravitino mass of size $1/2R$, and preserving R-symmetry, is obtained when the $N = 2$ supergravity is broken through a Scherk-Schwarz mechanism. Alternatively, the effect of R-symmetry being broken by minimal coupling to supergravity may be estimated and used as a minimum estimate for Majorana masses induced in the model [26, 27].

Finally, the simultaneous presence of μ and $B\mu$ terms in the Higgs sector is incompatible with R-symmetry. It is difficult to arrange a viable electroweak symmetry breaking which preserves R-symmetry and satisfies the LEP bound. There are some interesting possibilities, such as adding extra fields [10] or interactions with the supersymmetry breaking sector to generate new Higgs couplings [27]. However, we shall take the philosophy that, being a chiral symmetry, it is natural for R-symmetry to be broken in the Higgs sector. Our approach has the advantage that we do not need to introduce any interactions between the Higgs and supersymmetry-breaking sectors, and no additional mass scales.

It is popular to use the adjunction of a singlet to the MSSM as a way to address some its issues. One is the so-called μ-problem and it is the main motivation of the NMSSM (see [29] and references therein). There one starts by a vanishing tree-level μ-term, as it is forbidden by a Peccei-Quinn symmetry, and adds to the MSSM a singlet with coupling $\lambda \mathbf{S}\mathbf{H_u} \cdot \mathbf{H_d}$. The potential of the singlet is arranged such that the breaking of supersymmetry induces a vacuum expectation value to S naturally of the order of electroweak scale.

Another issue is the possibility to increase the Higgs mass. In the MSSM, the tree level quartic Higgs coupling is given by the $SU(2)_W \times U(1)_Y$ D-terms. Therefore it is governed by the strength of associated gauge couplings. This implies that the lightest Higgs boson mass is bounded to be smaller than M_Z. Agreement with present collider experiments is then obtained by radiative corrections. With the presence of the coupling $\lambda \mathbf{S}\mathbf{H_u} \cdot \mathbf{H_d}$, the lightest Higgs mass is now bounded, at tree level, to be:

$$m_h^2 \leq M_Z^2 (s_W^2 c_{2\beta}^2 + \frac{2\lambda^2}{(g')^2 + g^2} s_{2\beta}^2) \tag{1.1}$$

where g' and g are the gauge couplings of the hypercharge $U(1)_Y$, and the weak $SU(2)_W$ respectively, while β is defined by the ratio of the two Higgs vevs, $\tan\beta = \frac{<H_u>}{<H_d>}$. It is clear from (1.1) that a tree-level bound larger than M_Z can be obtained for a large value of λ, and $s_{2\beta}^2 \to 1$. It can be shown that such a possibility is in agreement with bounds from the electroweak precision measurements [31].

We wish to study here if the singlet $\mathbf{S}$ that pairs up with the bino, allowing a Dirac soft mass, can be used in similar ways. We will show how the addition of a cubic superpotential coupling $\frac{\kappa}{3}\mathbf{S}^3$ may indeed allow the generation of both μ and $B\mu$ terms, and to push the tree-level Higgs mass above the LEP bound. It is important to stress that although it is in many ways similar to the NMSSM, there are additional particles and couplings, and therefore a separate study is required.

The study here will be performed from an entirely low-energy perspective. However, if we consider the UV completion of the model, one might ask what symmetries allow the generation of the cubic superpotential coupling and not other similar polynomials in $\mathbf{S}$. We could suppose that the model contains different sectors obeying different symmetries, the singlet being a "bulk" state that belongs to and interacts with all of them. These sectors are classified as

- a hidden or secluded sector that breaks supersymmetry respects $U(1)_R$ symmetry. Supersymmetry breaking appears as the vev of a D-term or an F-term with zero R-charge. We include supersymmetry breaking messengers in this sector.
- another hidden sector which instead respects a $\mathbb{Z}_3$ symmetry under which the singlet transforms, but violates the $U(1)_R$ one.
- the visible sector contains the MSSM fields, as well as the $SU(2)_w$ triplet and $SU(3)_c$ octet DG-adjoints. This sector respects the $U(1)_R$ symmetry, but violates the $\mathbb{Z}_3$ in its couplings to $\mathbf{S}$.

The interactions between the different sectors will lead to a collective breaking of supersymmetry, the $U(1)_R$ and the $\mathbb{Z}_3$ symmetries.

The singlet $\mathbf{S}$ may appear either as an elementary or composite state, in which case its coupling to visible matter can grow large rapidly with energy. Its coupling to the observable sector breaks the $\mathbb{Z}_3$ symmetry by the $\lambda \mathbf{S}\mathbf{H_u} \cdot \mathbf{H_d}$ and Dirac gaugino couplings. On the other hand, its coupling to the $\mathbb{Z}_3$ symmetry preserving sector is assumed to give rise to a superpotential term $\frac{\kappa}{3}\mathbf{S}^3$ which represents the only operator violating R-symmetry (at tree level) relevant to the visible sector.

This note is organized as follows. Section 2 introduces the notations and the general framework for the model. The patterns of electroweak symmetry breaking partially discussed in [8, 14, 23], are reviewed in section 3 in order to include the effects the coupling κ. An important constraint on such an extension of the MSSM comes from the ρ parameter due to the vev of the $SU(2)_W$ triplet and is discussed in section 4. The possibility of allowing a heavy Higgs is mentioned in section 5. In section 6 the possible generation of μ and/or $B\mu$ terms through the vev of the singlet S is studied, and the main patterns of the resulting spectrum are discussed, with a set of numerical examples. Conclusions are drawn in section 7.

2 The Model

In order to allow Dirac masses for the gauginos, the particle content of the MSSM is extended by states in the adjoint representations:

$$\mathbf{S} = S + \sqrt{2}\theta\chi_S + \cdots \tag{2.1}$$
$$\mathbf{T} = T + \sqrt{2}\theta\chi_T + \cdots \tag{2.2}$$
$$\mathbf{O_g} = O_g + \sqrt{2}\theta\chi_g + \cdots \tag{2.3}$$

where S is a singlet, $\mathbf{O_g}$ a color octet and $T = \sum_{a=1,2,3} T^{(a)}$ an $SU(2)$ triplet. The latter can be written as:

$$T^{(1)} = T_1\frac{\sigma^1}{2}, \qquad T^{(2)} = T_2\frac{\sigma^2}{2}, \qquad T^{(3)} = T_0\frac{\sigma^3}{2},$$
$$T = \frac{1}{2}\begin{pmatrix} T_0 & \sqrt{2}T_+ \\ \sqrt{2}T_- & -T_0 \end{pmatrix},$$
$$T_0 = \frac{1}{\sqrt{2}}(T_R + iT_I), \quad T_+ = \frac{1}{\sqrt{2}}(T_{+R} + iT_{+I}), \qquad T_- = \frac{1}{\sqrt{2}}(T_{-R} + iT_{-I}), \tag{2.4}$$

and σ^a are the Pauli matrices.

The Dirac gaugino masses are described with superfields by the Lagrangian:

$$\mathcal{L}^{Dirac}_{gaugino} = \int d^2\theta \left[\sqrt{2}\mathbf{m}^{\alpha}_{1D}\mathbf{W}_{1\alpha}\mathbf{S} + 2\sqrt{2}\mathbf{m}^{\alpha}_{2D}\mathrm{tr}(\mathbf{W}_{2\alpha}\mathbf{T}) + 2\sqrt{2}\mathbf{m}^{\alpha}_{3D}\mathrm{tr}(\mathbf{W}_{3\alpha}\mathbf{O_g})\right] + h.c. \tag{2.5}$$

where we have introduced spurion superfields

$$\mathbf{m}_{\alpha iD} = \theta_\alpha m_{iD}. \tag{2.6}$$

The integration over the Grasmannian coordinates leads to

$$\int d^2\theta 2\sqrt{2}m_D\theta^\alpha \mathrm{tr}(W_\alpha\Sigma) \supset -m_D(\lambda_a\psi_a) + \sqrt{2}m_D\Sigma_a D_a \tag{2.7}$$

Then with $D^a_b = -g_b\phi^\dagger_i R^a_b(i)\phi_i$ (where $R^a_b(i)$ is the a^{th} generator of the group b in the representation of field i, and $R^b_Y(i) = Y(i)$ for the hypercharge) we find

$$\mathcal{L} \supset -m_{bD}\sqrt{2}g_b\Sigma_a\phi^\dagger R^a_b\phi \tag{2.8}$$

The Higgs sector of the model is described by the superpotential:

$$W^{(s)} = \mu \mathbf{H_u}\cdot\mathbf{H_d} + \lambda_S \mathbf{S}\mathbf{H_u}\cdot\mathbf{H_d} + 2\lambda_T \mathbf{H_d}\cdot\mathbf{T}\mathbf{H_u} + \frac{M_S}{2}\mathbf{S}^2 + \frac{\kappa}{3}\mathbf{S}^3 + M_T \mathrm{tr}(\mathbf{T}\mathbf{T}) + M_O \mathrm{tr}(\mathbf{O_g}\mathbf{O_g}), \tag{2.9}$$

the Higgs soft masses

$$m^2_{H_u}|H_u|^2 + m^2_{H_d}|H_d|^2 + [\tilde{B}\mu H_u \cdot H_d + h.c.] \tag{2.10}$$

as well as soft terms involving the DG-adjoint fields

$$\begin{aligned} -\Delta\mathcal{L}^{soft} &= m_S^2|S|^2 + \frac{1}{2}B_S(S^2 + h.c.) + 2m_T^2 \mathrm{tr}(T^\dagger T) + B_T(\mathrm{tr}(TT) + h.c.) \\ &+ A_S\lambda_S S H_u \cdot H_d + 2A_T\lambda_T H_d \cdot T H_u + \frac{1}{3}\kappa A_\kappa S^3 \\ &+ 2m_O^2 \mathrm{tr}(O_g^\dagger O_g) + B_O(\mathrm{tr}(O_g O_g) + h.c.) \end{aligned} \tag{2.11}$$

with the definition $H_u \cdot H_d = H_u^+ H_d^- - H_u^0 H_d^0$.

In this work we will restrict for simplicity to the above terms, while the most general renormalisable Lagrangian includes additional superpotential interactions[1]

$$W_2^{(s)} = L\mathbf{S} + \lambda_{ST}\mathbf{S}\mathrm{tr}(\mathbf{T}\mathbf{T}) + \lambda_{SO}\mathbf{S}\mathrm{tr}(\mathbf{O_g}\mathbf{O_g}) + \frac{\kappa_O}{3}\mathrm{tr}(\mathbf{O_g}\mathbf{O_g}\mathbf{O_g}). \tag{2.12}$$

as well as adjoint scalar A-terms (including the possible scalar tadpole) are given by

$$-\Delta\mathcal{L}_2^{soft} = t^S S + \lambda_{ST}A_{ST}S\mathrm{tr}(TT) + \lambda_{SO}A_{SO}S\mathrm{tr}(O_g O_g) + \frac{1}{3}\kappa_O A_{\kappa_O}\mathrm{tr}(O_g O_g O_g) + h.c. \tag{2.13}$$

and we require $L = t^S = \lambda_{ST} = \lambda_{SO} = \kappa_O = 0$.

3 Electroweak symmetry breaking potential

We define for the neutral components [14, 23]:

$$H_u^0 = \frac{H_{uR}^0 + iH_{uI}^0}{\sqrt{2}} = \frac{1}{\sqrt{2}}[s_\beta(v+h) + c_\beta H + i(c_\beta A - s_\beta G^0)], \tag{3.1}$$

$$H_d^0 = \frac{H_{dR}^0 + iH_{dI}^0}{\sqrt{2}} = \frac{1}{\sqrt{2}}[c_\beta(v+h) - s_\beta H + i(s_\beta A + c_\beta G^0)], \tag{3.2}$$

where G^0 is the would-be Goldstone boson (traded for the longitudinal component of the Z-boson).

We shall use the compact notation:

$$\begin{aligned} c_\beta &\equiv \cos\beta, \quad s_\beta \equiv \sin\beta, \quad t_\beta \equiv \tan\beta \\ c_{2\beta} &\equiv \cos 2\beta, \quad s_{2\beta} \equiv \sin 2\beta \end{aligned} \tag{3.3}$$

[1] Note there are no terms $\mathrm{tr}(\mathbf{T}), \mathrm{tr}(\mathbf{O_g}), \mathrm{tr}(\mathbf{TTT})$ since these vanish by gauge invariance.

Within these conventions, $v \simeq 246$ GeV, $\frac{(g')^2+g^2}{4}v^2 = M_Z^2$. We are interested by the case of CP neutral vacuum, i.e. $H_{dI}^0 = H_{uI}^0 = 0$ which implies $S_I = T_I = 0$ [14]. The CP-even singlet and neutral component of the triplet may also acquire a vacuum expectation value, so we define

$$S = \frac{1}{\sqrt{2}}((v_s + S_R + iS_I) \qquad T^0 = \frac{1}{\sqrt{2}}(v_T + T_R + iT_I) \tag{3.4}$$

It will also be useful to introduce the following effective mass parameters:

$$\begin{aligned} \tilde{\mu} &= \mu + \frac{1}{\sqrt{2}}(\lambda_S\, v_S + \lambda_T\, v_T) \\ \tilde{B}\mu &= B\mu + \frac{\lambda_S}{\sqrt{2}}(M_S + A_S)v_S + \frac{\lambda_T}{\sqrt{2}}(M_T + A_T)v_T + \frac{1}{2}\lambda_S\, \kappa\, v_S^2 \end{aligned} \tag{3.5}$$

3.1 Equations of motion for the CP-even neutral fields

The scalar potential for the CP-even neutral fields with is given by:

$$\begin{aligned} V_{EW} &= \left[\frac{g^2+g'^2}{4}c_{2\beta}^2 \;+\; \frac{\lambda_S^2+\lambda_T^2}{2}s_{2\beta}^2\right]\frac{v^4}{8} \\ &+ \left[m_{H_u}^2 s_\beta^2 + m_{H_d}^2 c_\beta^2 \;+\tilde{\mu}^2 - \tilde{B}\mu\, s_{2\beta} \;+(g\, m_{2D}\, v_T - g' m_{1D}\, v_S)c_{2\beta}\right]\frac{v^2}{2} \\ &+ \frac{\kappa^2}{4}v_S^4 + \frac{\kappa}{\sqrt{2}}\frac{(3M_S + A_S)}{3}\,v_S^3 + \frac{1}{2}\tilde{m}_{SR}^2\, v_S^2 + \;\frac{1}{2}\tilde{m}_{TR}^2\, v_T^2 \end{aligned} \tag{3.6}$$

where the effective masses for the real parts of the S and T fields read:

$$\tilde{m}_{SR}^2 \;=\; M_S^2 + m_S^2 + 4m_{1D}^2 + B_S, \qquad \tilde{m}_{TR}^2 = M_T^2 + m_T^2 + 4m_{2D}^2 + B_T \tag{3.7}$$

There is no restriction on the sign of the different mass parameters m_S^2 and B_S at this stage.

The imaginary parts of the fields have been dropped as their vevs are vanishing due to the assumed CP conservation [14]. The coefficients of the corresponding quadratic terms:

$$\tilde{m}_{SI}^2 = M_S^2 + m_S^2 - B_S, \qquad \tilde{m}_{TI}^2 = M_T^2 + m_T^2 - B_T \tag{3.8}$$

do not, in contrast to the CP-even partners, receive contributions from D-terms proportional to the Dirac masses.

As it is customarily done for the (N)MSSM, the minimization of the scalar potential allows here also to express $\tilde{\mu}$ and $\tilde{B}\mu$ as a function of the other parameters:

$$\tilde{\mu}^2 + \frac{M_Z^2}{2} = \frac{m_{H_d}^2 - t_\beta^2\, m_{H_u}^2}{t_\beta^2 - 1} + \left[\frac{t_\beta^2+1}{t_\beta^2-1}\right]\left(g\, m_{2D}\, v_T - g' m_{1D}\, v_S\right) \tag{3.9}$$

and

$$M_A^2 \;\equiv\; \frac{2\tilde{B}\mu}{s_{2\beta}} = 2\tilde{\mu}^2 + m_{H_u}^2 + m_{H_d}^2 - \frac{\lambda_S^2+\lambda_T^2}{2}v^2 c_\beta^2 \tag{3.10}$$

where:

$$v_T \;\simeq\; \frac{v^2}{2\tilde{m}_{TR}^2}\;\left[-gm_{2D}c_{2\beta} - \sqrt{2}\tilde{\mu}\lambda_T + \frac{\lambda_T}{\sqrt{2}}(M_T + A_T)s_{2\beta}\right], \tag{3.11}$$

while v_S is determined as a solution for the cubic equation:

$$0 \;=\; \kappa^2 v_S^3 + \frac{\kappa}{\sqrt{2}}(A_\kappa + 3M_S)v_S^2 - \tilde{m}_{S0}^2 v_S + v_0^3 \tag{3.12}$$

under the assumption $\tilde{m}_{S0}^2, \tilde{m}_{S0}^2 \gg \lambda_S\lambda_T v^2$ [14] with

$$\tilde{m}_{S0}^2 = -\tilde{m}_{SR}^2 - \lambda_S^2\,\frac{v^2}{2} + \kappa\,\lambda_S\,\frac{v^2}{2}s_{2\beta} \tag{3.13}$$

and

$$v_0^3 \;=\; -\frac{v^2}{2}\left[g'm_{1D}c_{2\beta} - \lambda_S\left(\sqrt{2}\mu - \frac{(A_\kappa + M_S)}{\sqrt{2}}s_{2\beta} + \lambda_T v_T\right)\right]. \tag{3.14}$$

3.2 Masses of the CP even neutral scalars

The coefficient of the quadratic term for the scalar singlet S_R is given by the effective mass:

$$\begin{aligned}
\tilde{m}_S^2 \;&=\; -\tilde{m}_{S0}^2 + 3\kappa^2 v_S^2 + \frac{\sqrt{2}}{3}\kappa\, v_S\,(A_\kappa + 3M_S) && (3.15)\\
&=\; \tilde{m}_{SR}^2 + \lambda_S^2\,\frac{v^2}{2} - \kappa\,\lambda_S\,\frac{v^2}{2}s_{2\beta} + 3\kappa^2 v_S^2 + \frac{\sqrt{2}}{3}\kappa\, v_S\,(A_\kappa + 3M_S) && (3.16)
\end{aligned}$$

while it is

$$\tilde{m}_T^2 = \tilde{m}_{TR}^2 + \lambda_T^2\frac{v^2}{2} \tag{3.17}$$

for the neutral component of the triplet T_R^0.

The resulting CP even scalars mass matrix takes, in the basis $\{h, H, S_R, T_R^0\}$ the form:

$$\begin{pmatrix}
M_Z^2 + \Delta_h s_{2\beta}^2 & \Delta_h s_{2\beta}c_{2\beta} & \Delta_{hs} & \Delta_{ht}\\
\Delta_h s_{2\beta}c_{2\beta} & M_A^2 - \Delta_h s_{2\beta}^2 & \Delta_{Hs} & \Delta_{Ht}\\
\Delta_{hs} & \Delta_{Hs} & \tilde{m}_S^2 & \lambda_S\lambda_T\frac{v^2}{2}\\
\Delta_{ht} & \Delta_{Ht} & \lambda_S\lambda_T\frac{v^2}{2} & \tilde{m}_T^2
\end{pmatrix} \tag{3.18}$$

where we have introduced the compact notation:

$$\Delta_h \;=\; \frac{v^2}{2}(\lambda_S^2 + \lambda_T^2) - M_Z^2 \tag{3.19}$$

which vanishes when λ_S and λ_T take their $N = 2$ values [8]. We denote non-diagonal elements describing the mixing of S_R and T_R^0 states with the light Higgs h:

$$\Delta_{hs} = -2\frac{v_S}{v}\tilde{m}_{SR}^2 - \sqrt{2}\kappa\frac{v_S^2}{v}(A_\kappa + 3M_S) - 2\kappa^2\frac{v_S^3}{v}, \qquad \Delta_{ht} = -2\frac{v_T}{v}\tilde{m}_{TR}^2 \tag{3.20}$$

while

$$\Delta_{Hs} = g' m_{1D} v s_{2\beta} - \lambda_S \frac{v(A_s + M_s)}{\sqrt{2}} c_{2\beta}, \qquad \Delta_{Ht} = -g m_{2D} v s_{2\beta} - \lambda_T \frac{v(A_T + M_T)}{\sqrt{2}} c_{2\beta} \tag{3.21}$$

stand for the corresponding mixing with heavier Higgs, H.

From this mass matrix, we see that the lightest Higgs scalar mass is bounded to be:

$$m_h^2 \leq M_Z^2 c_{2\beta}^2 + \frac{v^2}{2}(\lambda_S^2 + \lambda_T^2) s_{2\beta}^2. \tag{3.22}$$

3.3 Masses of the CP odd neutral scalars

The CP odd neutral scalars mass matrix is given, in the basis $\{A, S_I, T_I^0\}$, by

$$\begin{pmatrix} M_A^2 & -\lambda_S v[\frac{(M_S - A_S)}{\sqrt{2}} + k v_S] & -\lambda_S v \frac{(A_T - M_T)}{\sqrt{2}} \\ -\lambda_S v[\frac{(M_S - A_S)}{\sqrt{2}} + k v_S] & \tilde{m}_{aS}^2 & \lambda_S \lambda_T \frac{v^2}{2} \\ -\lambda_S v \frac{(A_T - M_T)}{\sqrt{2}} & \lambda_S \lambda_T \frac{v^2}{2} & \tilde{m}_{TI}^2 + \lambda_T \frac{v^2}{2} \end{pmatrix} \tag{3.23}$$

where

$$\tilde{m}_{aS}^2 = \tilde{m}_{SI}^2 + \lambda_S^2 \frac{v^2}{2} + \kappa \lambda_S \frac{v^2}{2} s_{2\beta} + \kappa^2 v_s^2 + \sqrt{2} \kappa v_s (M_S - A_\kappa) \tag{3.24}$$

is the coefficient of quadratic term in the effective Lagrangian for the imaginary part of the singlet S.

3.4 Masses of charged scalars

The charged would-be-Goldstone bosons, traded for the $W_\pm$ longitudinal modes, take now the form:

$$G_T^\pm = \frac{1}{\sqrt{\rho}}(G^\pm + \sqrt{2}\frac{v_T}{v}(T_\pm + (T_\mp)^*)) \tag{3.25}$$

where the erstwhile Goldstone boson (before the triplet coupling is switched on) is

$$\begin{aligned} G^+ &\equiv c_\beta \overline{H}_d^- - s_\beta H_u^+ \\ G^- &\equiv c_\beta H_d^- - s_\beta \overline{H}_u^+, \end{aligned} \tag{3.26}$$

the tree-level ρ-parameter given by:

$$\rho = 1 + 4\frac{v_t^2}{v^2}, \tag{3.27}$$

and the remaning orthogonal combinations of charged states can be written as

$$T_I^+ = \frac{1}{\sqrt{2}}(T_+ - (T_-)^*) \qquad T_I^- = (T_I^+)^* \tag{3.28}$$

$$T_{RG}^+ = \frac{1}{\sqrt{\rho}}(\frac{1}{\sqrt{2}}(T_\pm + (T_\mp)^*) - 2\frac{v_T}{v}G^+) \qquad T_{RG}^- = (T_{RG}^+)^*. \tag{3.29}$$

The mass matrix in the Lagrangian $\mathcal{L}^{Dirac}_{gaugino} = -(\Sigma^+)^T M_{Ch} \Sigma^-$ with $\Sigma^+ = \{H^+, T_I^+, T_{RG}^+\}$ and $\Sigma^- = \{H^-, T_I^-, T_{RG}^-\}$ reads then:

$$\begin{pmatrix} \tilde{m}_{H^\pm} & (\frac{g^2-2\lambda_T^2}{2} v_T s_{2\beta} - \frac{\lambda_T(M_T-A_T)}{\sqrt{2}})v & -\sqrt{\rho}\Delta_t \\ (\frac{g^2-2\lambda_T^2}{2} v_T s_{2\beta} - \frac{\lambda_T(M_T-A_T)}{\sqrt{2}})v & \tilde{m}^2_{TI} + \lambda_T^2 \frac{v^2}{2} + g^2 v_T^2 & -\sqrt{\rho}\frac{g^2-2\lambda_T^2}{4} v^2 c_{2\beta} \\ -\sqrt{\rho}\Delta_t & -\sqrt{\rho}\frac{g^2-2\lambda_T^2}{4} v^2 c_{2\beta} & \rho \tilde{m}_T^2 \end{pmatrix} \tag{3.30}$$

where the charged Higgs $H^\pm$ quadratic term reads:

$$\begin{aligned} \tilde{m}^2_{H^\pm} = \quad & M_Z^2 c_W^2 + M_A^2 + (\lambda_T^2 - \lambda_S^2)\frac{v^2}{2} + 2g m_{2D} c_{2\beta} v_T \\ & +2\lambda_T^2 v_T^2 - 2\sqrt{2}\lambda_T v_T \tilde{\mu} - \sqrt{2}\lambda_T (M_T + A_T) v_T. \end{aligned} \tag{3.31}$$

This result is given here for completeness, but can be obtained from [23] upon the use of the new M_A^2, which contains all the dependance on κ.

4 The $SU(2)_W$ triplet vacuum expectation value

The neutral components of the DG-adjoint scalars acquire non-vanishing expectation values at the minimum of the electroweak scalar potential. The minimization with respect to the neutral component of the triplet leads to:

$$v_T = \frac{v^2}{2(M_T^2 + m_T^2 + 4m_{2D}^2 + B_T)} \left[-g m_{2D} c_{2\beta} - \sqrt{2}\tilde{\mu}\lambda_T + \frac{\lambda_T}{\sqrt{2}}(M_T + A_T) s_{2\beta} \right] \tag{4.1}$$

As this contributes to the W boson mass, the electroweak precision data give important bounds on the parameters of the model. For instance, using $\rho \simeq 1 + \alpha T = 1.0004^{+0.0008}_{-0.0004}$ [30], we require:

$$\Delta\rho \simeq 4\frac{v_t^2}{v^2} \lesssim 8 \cdot 10^{-4} \tag{4.2}$$

which is satisfied for $v_t \lesssim 3$ GeV.

Given our assumption on having the A-term parameters, such as A_T, small, there are three different ways to satisfy the bound on v_t:

- One is to have a large supersymmetric triplet mass. In the limit $M_T \to \infty$, the v_T vanishes and the full triplet superfield decouples. A Majorana mass for the winos is then required in order to avoid the charginos being too light.
- A second possibility is to satisfy the bound by taking instead m_{2D} large, of the order of $\gtrsim 2$ TeV. It is also meaningful to take simultaneously m_T to be large as it is anyway expected to be in models of gauge mediation. This makes not only the triplet heavy, but also the wino. In fact, in the limit $m_{2D} \to \infty$, $m_T \to \infty$

and $m_T/m_{2D} \to 0$, the weak $SU(2)$ D-term contribution to the effective Higgs quartic coupling cancels, and the scalar potential becomes:

$$\begin{aligned} V_{EW} &= \left[\frac{g'^2}{4}c_{2\beta}^2 \;+\; \frac{\lambda_S^2+\lambda_T^2}{2}s_{2\beta}^2\right]\frac{v^4}{8} \\ &\quad + \left[m_{H_u}^2 s_\beta^2 + m_{H_d}^2 c_\beta^2 \;+\tilde{\mu}^2 - \tilde{B}\mu\, s_{2\beta} \;- g' m_{1D}\, v_S c_{2\beta}\right]\frac{v^2}{2} \\ &\quad + \frac{\kappa^2}{4}v_S^4 + \frac{\kappa}{\sqrt{2}}\frac{(3M_S+A_S)}{3}v_S^3 + \frac{1}{2}\tilde{m}_{sR}^2 v_S^2 \end{aligned} \tag{4.3}$$

with subsequent modification of the scalar mass matrices discussed in the previous section.

- Finally, the limit on v_t can be satisfied just by taking m_T large enough, keeping $m_{2D}/m_T \to 0$. In which case, the electroweak neutral fields' scalar potential becomes:

$$\begin{aligned} V_{EW} &= \left[\frac{g^2+g'^2}{4}c_{2\beta}^2 \;+\; \frac{\lambda_S^2+\lambda_T^2}{2}s_{2\beta}^2\right]\frac{v^4}{8} \\ &\quad + \left[m_{H_u}^2 s_\beta^2 + m_{H_d}^2 c_\beta^2 \;+\tilde{\mu}^2 - \tilde{B}\mu\, s_{2\beta} \;- g' m_{1D}\, v_S c_{2\beta}\right]\frac{v^2}{2} \\ &\quad + \frac{\kappa^2}{4}v_S^4 + \frac{\kappa}{\sqrt{2}}\frac{(3M_S+A_S)}{3}v_S^3 + \frac{1}{2}\tilde{m}_{SR}^2 v_S^2, \end{aligned} \tag{4.4}$$

where we see that both contributions from the $SU(2)_W$ D-term and λ_T remain.

We shall focus in the following on the last two cases.

5 Allowing a heavy Higgs

From the CP-even scalar mass matrix, we see that the lightest Higgs scalar mass is bounded by:

$$m_h^2 \leq M_Z^2 c_{2\beta}^2 + \frac{v^2}{2}(\lambda_S^2+\lambda_T^2)s_{2\beta}^2. \tag{5.1}$$

In order to obtain a spectrum with $m_h > M_Z$, we need both a large $\lambda_S^2 + \lambda_T^2$ and $s_{2\beta}^2 \to 1$. While the second condition can be satisfied quite easily, the first one requires the presence of new physics at energies Λ_{NP} of a few tens of TeV, and may be subject to constraints from electroweak precision tests. In fact, as the couplings λ_T, λ_S grow with energy, and now start with a large value at the electroweak scale, they become non-pertubative in the ultraviolet very rapidly at the scale Λ_{NP}. We expect this scale to be, at the lowest, in the range 10–50 TeV in order to keep harmless any contributions from the new physics to electroweak precision tests. Following the one-loop renormalization equations:

$$\begin{aligned}16\pi^2\frac{d}{dt}\lambda_S =&\lambda_S\left[4\lambda_S^2+6\lambda_T^2+2|\kappa_S|^2+\cdots\right]\\16\pi^2\frac{d}{dt}\lambda_T =&\lambda_T\left[2\lambda_S^2+8\lambda_T^2+\cdots\right]\\16\pi^2\frac{d}{dt}\kappa_S =&\kappa_S\left[6\lambda_S^2+6|\kappa_S|^2\right].\end{aligned} \tag{5.2}$$

at scale $t = \log E$, with $E > m_{2D}, m_T$, we can consider the implications for the three ways to obtain a heavy Higgs scenario:

1. The **S**-way where the coupling λ_S is chosen to be large enough while $\lambda_T \ll \lambda_S$. A bound on the size of the λ_S^2 coupling arise from the requirement that the coupling remains perturbative up to an energy of the order of 10 TeV. In such a case, the maximal value of λ_S at low scale ranges from approximatively $\lambda_S \sim 1.4$ for $\Lambda_{NP} \sim 50$ TeV, to $\lambda_S \sim 2$ for $\Lambda_{NP} \sim 10$ TeV.
2. The **T**-way, where the maximal value is somewhat lower since the coefficient in the RGE is bigger, due to the fact that above a few TeVs all the triplet fields contribute. The maximal value is now smaller of order $\lambda_T \sim$ 1–1.5 for $\Lambda_{NP} \sim$ 10–50 TeV.
3. A combination of the two, the **S**/**T**-way with smaller value for each as they both run to a non-perturbative regime, with a maximal value of $\lambda_{S,T} \sim$ 0.9–1.2 for $\Lambda_{NP} \sim$ 10–50 TeV.

Note that in order to allow the largest values of λ_S, one needs to keep κ small as it contribute to the corresponding RGEs.

First, let us assume the presence of both μ and $B\mu$ terms not generated by the singlet vev v_S. The alternative possibility shall be investigated in the next section. The μ term by itself does not break R-symmetry, but then B_μ does. Expecting the same source of R-symmetry breaking to provide contributions to other soft terms, requires then to assume B_μ and thus M_A to be small.

We can consider the simplest case of a very massive singlet $\tilde{m}_S \gg v$, so that at low energies we remain with effectively two Higgs doublets. The resulting CP even scalar mass matrix becomes:

$$\begin{pmatrix} M_Z^2 s_W^2 + \Delta_h s_{2\beta}^2 & \Delta_h s_{2\beta} c_{2\beta} \\ \Delta_h s_{2\beta} c_{2\beta} & M_A^2 - \Delta_h s_{2\beta}^2 \end{pmatrix} \tag{5.3}$$

which has as eigenvalues:

$$\frac{1}{2}\left(M_A^2 + M_Z^2 s_W^2 \pm \left[\Delta_h^2 + (\Delta_h - M_A^2 + M_Z^2 s_W^2)^2 - 2\Delta_h(\Delta_h - M_A^2 + M_Z^2 s_W^2)c_{4\beta}\right]^{1/2}\right). \tag{5.4}$$

Note that the case of λSUSY considered in [31] corresponds to neglecting the terms $M_Z^2 s_W^2$ and leads to (for large M_A^2):

$$m_{H,h} \simeq \frac{1}{2}\left(M_A^2 \pm \left[M_A^4 - 2\Delta_h(M_A^2 - \Delta_h)c_{2\beta}^2\right]^{1/2}\right). \tag{5.5}$$

6 The singlet vacuum expectation value

A scalar component $\mathbb{S}$ with masses of a few hundred GeV does not play an important phenomenological role, so one can assume a large $\tilde{m}_{S0}$ (as in [31]) . For simplicity (and in line with our philosophy of κ being the source of R-symmetry breaking) we will consider $M_S = 0$ and negligible A-terms: $A_S \simeq A_\kappa \simeq 0$, which can be obtained in gauge mediated models with an approximate R-symmetric supersymmetry breaking.

For the case with $\kappa = 0$, as motivated for example some $N = 2$ origin of the DG-adjoints [14], one gets:

$$
\begin{aligned}
v_S &\simeq \frac{v_0^3}{\tilde{m}_{S0}^2} \\
&\simeq \frac{v^2}{2(m_S^2 + 4m_{1D}^2 + B_S + \lambda_S^2 \frac{v^2}{2})} \left[\sqrt{2}\mu\lambda_S - g' m_{1D} c_{2\beta}\right]. \qquad (6.1)
\end{aligned}
$$

6.1 The generation of μ and $B\mu$-terms

We would like to switch on the trilinear coupling κ for the singlet S as the only R-symmetry breaking parameter in our model. We proceed to investigate if a $\tilde{\mu}$ and/or $\tilde{B}\mu$ can be generated in a way similar to the case in the NMSSM. For simplicity, we take A_κ smaller than the other relevant mass parameters.

Let us first consider the case of a large $\tilde{m}_{S0}$, i.e.$\tilde{m}_{S0}^2 v_s \gg v_0^3$. The term $\tilde{m}_{S0}^2 v_s$ and $\kappa^2 v_S^3$ dominate over the others in the equation determining the singlet. The solution can be approximated as:

$$
\begin{aligned}
|v_S| &\simeq \frac{\tilde{m}_{S0}}{\kappa} + \frac{v_0^3}{2\tilde{m}_{S0}^2} + \cdots \qquad (6.2) \\
&\simeq \frac{v}{\sqrt{2}\kappa}\left(-\frac{2\tilde{m}_{SR}^2}{v^2} - \lambda_S^2 + \kappa\,\lambda_S\, s_{2\beta}\right)^{1/2} - \frac{g' m_{1D} c_{2\beta}}{2(\frac{2\tilde{m}_{SR}^2}{v^2} + \lambda_S^2 - \kappa\,\lambda_S\, s_{2\beta})} + \cdots \qquad (6.3)
\end{aligned}
$$

(6.4)

The validity of the approximation (6.4) requires $\tilde{m}_{S0}/\kappa$ to be larger than other masses, taking the effective quadratic term $\tilde{m}_{SR}^2$ to be governed by a sufficiently large $|B_S|$,

$$
- B_S = |B_S| \gg m_{1D}^2 \qquad (6.5)
$$

This hierarchy is not unnatural, it is even quite generic in models of gauge mediation with Dirac gauginos [17]. It is easy to find a generic set of messengers with couplings to the singlet that leads to large contribution to $|B_S|$ with the desired sign, following the results of [13]. Therefore:

$$
|v_S| \simeq \tfrac{|\tilde{m}_{SR}|}{\kappa}. \qquad (6.6)
$$

This singlet vev (6.4) induces both a μ-term

$$
\tilde{\mu}^2 = \frac{\lambda_S^2}{2} v_S^2 \simeq \frac{\lambda_S^2}{2\kappa^2}\left(-\frac{2\tilde{m}_{SR}^2}{v^2}\right)v^2 \simeq \frac{\lambda_S^2}{\kappa^2}(|B_S| - m_S^2 - 4m_{1D}^2) \qquad (6.7)
$$

and a $B\mu$-term

$$\tilde{B}\mu \simeq \frac{\kappa}{\lambda_S}\tilde{\mu}^2 \simeq \frac{\lambda_S}{\kappa}(|B_S| - m_S^2 - 4m_{1D}^2) \tag{6.8}$$

with sizes controlled by λ_S, κ and $(|B_S|-m_S^2)$. With four parameters we can obviously fit the desired values for $\tilde{\mu}$. When $\kappa \lesssim \lambda_S$, this implies that $B\mu \lesssim \mu^2$. A small R-symmetry breaking corresponds then to a small $B\mu$, and a hierarchy between $B\mu$ and μ^2 would require some amount of tuning between the parameters in (3.10).

Before discussing the implications for the scalar spectrum, it is important to notice that the term $\frac{\kappa}{3}\mathbf{S}^3$ induces through the vev of S a Majorana mass for the singlino χ_S. The neutralino mass matrix, in the basis $\tilde{S}, \tilde{B}, \tilde{T}^0, \tilde{W}^0, \tilde{H}_d^0, \tilde{H}_u^0$ reads:

$$\begin{pmatrix} M_S + \sqrt{2}\kappa v_S & m_{1D} & 0 & 0 & -\frac{\sqrt{2}\lambda_S}{g_Y} M_Z s_W s_\beta & -\frac{\sqrt{2}\lambda_S}{g_Y} M_Z s_W c_\beta \\ m_{1D} & M_1 & 0 & 0 & -M_Z s_W c_\beta & M_Z s_W s_\beta \\ 0 & 0 & M_T & m_{2D} & -\frac{\sqrt{2}\lambda_T}{g_2} M_Z c_W s_\beta & -\frac{\sqrt{2}\lambda_T}{g_2} M_Z c_W c_\beta \\ 0 & 0 & m_{2D} & M_2 & M_Z c_W c_\beta & -M_Z c_W s_\beta \\ -\frac{\sqrt{2}\lambda_S}{g_Y} M_Z s_W s_\beta & -M_Z s_W c_\beta & -\frac{\sqrt{2}\lambda_T}{g_2} M_Z c_W s_\beta & M_Z c_W c_\beta & 0 & -\tilde{\mu} \\ -\frac{\sqrt{2}\lambda_S}{g'} M_Z s_W c_\beta & M_Z s_W s_\beta & -\frac{\sqrt{2}\lambda_T}{g_2} M_Z c_W c_\beta & -M_Z c_W s_\beta & -\tilde{\mu} & 0 \end{pmatrix} \tag{6.9}$$

Therefore a Majorana mass of order:

$$M_1' = \sqrt{2}\kappa v_S \simeq \frac{2\kappa}{\lambda_S}\tilde{\mu} \tag{6.10}$$

spoils the pseudo-Dirac nature of the bino, unless we have

$$\frac{\kappa}{\lambda_S}\frac{\tilde{\mu}}{m_{1D}} \ll 1. \tag{6.11}$$

Noting (6.7) we see that this requires

$$\left(\frac{|B_S| - m_S^2}{m_{1D}^2}\right) - 4 \ll 1, \tag{6.12}$$

which represents a degree of fine-tuning. Moreover, it requires a hierarchy between the two couplings $\kappa \ll \lambda_S$ to avoid the neutralino being mostly Higgsino (and rendering the pseudo-Dirac mass for the Bino irrelevant) which can in turn cause the pseudoscalar Higgs to be light. Hence the most generic situation would be that the neutralino contains a Majorana mix of the Bino and singlino.

The chargino masses, $-\frac{1}{2}((v^-)^T M_{Ch} v^+ + (v^+)^T M_{Ch}^T v^- + h.c)$ in the basis $v^+ = (\tilde{T}^+, \tilde{W}^+, \tilde{H}_u^+)$, $v^- = (\tilde{T}^-, \tilde{W}^-, \tilde{H}_d^-)$, are given by

$$M_{Ch} = \begin{pmatrix} M_T & m_{2D} + g v_T & \frac{2\lambda_T}{g} M_Z c_W c_\beta \\ m_{2D} - g v_T & M_2 & \sqrt{2} M_Z c_W s_\beta \\ -\frac{2\lambda_T}{g} M_Z c_W s_\beta & \sqrt{2} M_Z c_W c_\beta & \tilde{\mu} - \sqrt{2}\lambda_T v_T \end{pmatrix} \tag{6.13}$$

and do not depend on κ.

6.2 The scalar spectrum

The quadratic term coefficient for CP even scalar singlet S_R is given by the effective mass:

$$\tilde{m}_S^2 \simeq 2\tilde{m}_{S0}^2 \simeq -2\tilde{m}_{SR}^2 \simeq 2(|B_S| - m_S^2 - 4m_{1D}^2) > 0 \tag{6.14}$$

while the one for CP-odd S_I reads:

$$\tilde{m}_{aS}^2 \quad = \quad 2|B_S| - 4m_{1D}^2 + 2\kappa\lambda_S \frac{v^2}{2} s_{2\beta}. \tag{6.15}$$

On the other hand the coefficients of the quadratic terms of T_R and T_I^0 are given by:

$$\tilde{m}_T^2 \simeq m_T^2 + 4m_{2D}^2 + B_T \tag{6.16}$$

$$\tilde{m}_{aT}^2 \simeq m_T^2 - B_T. \tag{6.17}$$

The CP odd neutral scalar mass is given by

$$M_A^2 = \frac{2\tilde{B}\mu}{s_{2\beta}} \quad \simeq \quad \frac{2\lambda_S}{s_{2\beta}\kappa}(|B_S| - m_S^2 - 4m_{1D}^2). \tag{6.18}$$

The charged Higgs mass is

$$\tilde{m}_{H^\pm}^2 \quad = \quad M_W^2 {s_{2\beta}}^2 + M_A^2 + (\lambda_T^2 - \lambda_S^2)\frac{v^2}{2}. \tag{6.19}$$

The off diagonal elements describing the mixing of S_R, T_R^0 components with the Higgs scalars h, H can be approximated as:

$$\Delta_{hs} \simeq -\sqrt{2}(\lambda_S + \kappa s_{2\beta})\tilde{\mu} v, \qquad \Delta_{ht} \simeq \sqrt{2}\lambda_T \tilde{\mu} v + g m_{2D} v c_{2\beta} \tag{6.20}$$

$$\Delta_s \simeq g' m_{1D} v s_{2\beta}, \qquad \Delta_t \simeq -g m_{2D} v s_{2\beta} \tag{6.21}$$

which means that neglecting mixings requires working with approximations of order $\mathcal{O}(v\tilde{\mu}/\tilde{m}_S^2)$ for S_R and $\mathcal{O}(v m_{2D}/\tilde{m}_T^2)$ for T_R^0 .

With both μ and B_μ generated by the singlet vev, we can look for illustrative examples, and we would in particular like to exhibit the possibility of a larger tree-level mass for the lightest Higgs. A set of such examples is given in table 1. The first two models have a large Higgs mass thanks to a larger value of λ_S. To avoid a light pseudo-scalar Higgs we required $\kappa \sim \lambda_S$. The list of gauge and Higgs sector soft masses can be read in the table, notably with a lightest CP-even neutral Higgs mass evading the LEP bound already at tree-level. For model I, $\lambda_S = \kappa = 1.2, \lambda_T = 0.1, \tan\beta = 1$ lead to the lightest Higgs mass of 202 GeV at tree level, while $\lambda_S = 0.8, \kappa = 06, \lambda_T = 0.1, \tan\beta = 1.38$ lead to a tree-level mass of 116 GeV. We note that the spectrum is sensitive to the relative values of the couplings $\lambda_S, \kappa, \lambda_T, \tan\beta$. In model III, we give an example with large λ_T and small λ_S.

Among noticeable features that can be seen in the presented spectra is that, due to the choice of a large M_A, we have $M_{H^\pm} \sim M_H \sim M_A$. We have also chosen in these examples to suppress the triplet expectation value by a large soft mass for the scalar triplet field. In model I, the lightest neutralino is made of approximatively 72 percent of bino, 18 percent of singlino, while they become approximatively 75 percent bino, 21 percent singlino in model II and nearly all the rest Higgsino. Instead, in model III, the lightest neutralino is made of Higgsinos for about 87 percent, 7 percent wino and 4 percent bino.

An alternative approach, more easily compatible with models perturbative up to the GUT scale such as those of [17], is to take the couplings $\lambda_S \sim \kappa, \lambda_T$ to be small so that λ_S, λ_T will not drastically contribute to the running of the Higgs masses. However in this case we can no longer rely on the tree level potential, as after electroweak symmetry breaking the Higgs mass will not exceed the LEP bounds. Instead we must include, as usual, loop corrections to the Higgs potential and work at relatively large $\tan\beta$. We give an example of this in table 2, taking squarks at 1100 GeV. In this case, we have an almost entirely bino–singlino neutralino, with composition 81%–18%.

7 Conclusions

In any extension of the MSSM, the presence of a singlet S immediately raises a question about the possibility to use it to generate μ and $B\mu$ parameters. The S vev, easily computed, was found in all previous studies to be typically too small. As the combination of μ and $B\mu$ break the $U(1)_R$ R-symmetry, their generation is related to the way this breaking is implemented in the potential of S. We have argued for the adjunction to the model of a superpotential term $\frac{1}{3}\kappa\mathbf{S}^3$, and shown that it is sufficient. Both μ and $B\mu$ are induced with a ratio $B\mu/\mu \sim \kappa/\lambda_S$, which should be close to unity in order to obtain a successful electroweak symmetry breaking without excessive tuning of the parameters. We have also verified that this allows to raise the tree-level mass for the lightest CP-even Higgs above the LEP bound. This can be achieved by moderate values of either of λ_S, λ_T as given in the explicit examples.

While this work exhibited the main features of the model, some important issues need to be investigated before deriving phenomenological implications of it for the LHC and dark matter search experiments. Foremost among these is a thorough study of the impact of the presence of the singlet and the triplet on electroweak precision tests, such as calculating the S and T parameters. We believe that this merits an independent careful study by itself in the light of the disagreements in the literature on the correct treatment and the result of such an analysis already at the level of the (non-supersymmetric) extension of the standard model with triplets (see for example [32–34]). Such an analysis will allow the determination of the regions allowed for the parameters λ_S, λ_T, m_{2D} and m_T beyond the tree level considerations discussed here. On the other hand, top-down models for hidden sectors that allow the generation of the $\kappa\mathbf{S}^3$, and investigation of the resulting pattern of the soft-terms are needed. We will return to these important issues in future work.

Input parameter	Model I	Model II	Model III
λ_S	1.2	0.8	0.1
λ_T	0.1	0.1	0.7
κ	1.2	0.6	0.2
$\tan\beta$	1	1.38	1.38
m_S^2	10^5 GeV^2	10^5 GeV^2	10^5 GeV^2
B_S	-10^6 GeV^2	-10^6 GeV^2	-10^6 GeV^2
m_T^2	4 10^6 GeV^2	4 10^6 GeV^2	4 10^6 GeV^2
B_T	0	0	0
A_κ	0	0	0
m_{H_u}	197 GeV	479 GeV	596 i GeV
m_{H_d}	287 i GeV	339 i GeV	642 GeV
m_{1D}	400 GeV	400 GeV	400 GeV
m_{2D}	600 GeV	600 GeV	800 GeV
Output parameter	**Model I**	**Model II**	**Model III**
v_S	-425 GeV	838 GeV	2548 GeV
v_T	0.3 GeV	0.3 GeV	-0,08 GeV
$\Delta\rho$	5.3 $\times 10^{-6}$	5.9 $\times 10^{-6}$	4 $\times 10^{-7}$
$\tilde{\mu}$	-361 GeV	474 GeV	180 GeV

Model I	Chargino masses:	612, 604, 352 GeV
	Neutralino masses:	740, 613, 606, 388, 352, 203 GeV
	Neutral scalars:	2332, 723, 467, 208 GeV
	Neutral pseudoscalars:	2001, 1211, 491 GeV
	Charged scalars:	2333, 2000, 471 GeV
Model II	Chargino masses:	622, 602, 455 GeV
	Neutralino masses:	732, 619, 605, 484, 456, 215 GeV
	Neutral scalars:	2333, 718, 580, 116 GeV
	Neutral pseudoscalars:	2001, 1181, 588 GeV
	Charged scalars:	2333, 2000, 583 GeV
Model III	Chargino masses:	812, 808, 178 GeV
	Neutralino masses:	842, 830, 730, 226, 189, 171 GeV
	Neutral scalars:	2564, 722, 354, 120 GeV
	Neutral pseudoscalars:	2005, 1166, 369 GeV
	Charged scalars:	2565, 2004, 394 GeV

Table 1: Examples of model parameters.

Acknowledgments

This work is supported in part by the European contract "UNILHC" PITN-GA-2009-237920. MDG is supported by the German Science Foundation (DFG) under SFB 676. AKM would like to thank the Laboratoire de Physique Théorique des hautes

Input parameter	Model IV
λ_S	0.13
λ_T	0.13
κ	0.13
$\tan\beta$	31.8
m_S^2	400000 GeV^2
B_S	-1500000 GeV^2
m_T^2	6000000 GeV^2
B_T	0
A_κ	0
m_{H_u}	718 GeV
m_{H_d}	2520 GeV
m_{1D}	400 GeV
m_{2D}	500 GeV
Output parameter	Model IV
v_S	-5220 GeV
v_T	1.78 GeV
$\Delta\rho$	0.00021
$\tilde{\mu}$	-480 GeV

Chargino masses:	559.828, 500.624, 423.868 GeV
Neutralino masses:	864.145, 553.39, 552.476, 430.067, 428.46, 184.964 GeV
Neutral scalars:	2707.7, 2646.11, 960.048, 118.035 GeV
Neutral pseudoscalars:	2707.68, 2449.59, 1536.49 GeV
Charged scalars:	2708.7, 2646.12, 2449.59 GeV

Table 2: An example with small $\lambda_S, \lambda_T, \kappa$, including top-stop loop corrections.

Energies (LPTHE), Paris for financial support and hospitality during her masters degree when she worked on this project, MDG would like to thank the Centre de Physique Théorique de l'Ecole Polytechnique, where part of this work was completed, for hospitality; KB thanks Roberto Franceschini and Pietro Slavich for discussions.

References

[1] P. Fayet, Phys. Lett. B **78**, 417 (1978).

[2] J. Polchinski and L. Susskind, Phys. Rev. D **26**, 3661 (1982).

[3] L. J. Hall and L. Randall, Nucl. Phys. B **352**, 289 (1991).

[4] P. J. Fox, A. E. Nelson and N. Weiner, JHEP **0208**, 035 (2002) [arXiv:hep-ph/0206096].

[5] A. E. Nelson, N. Rius, V. Sanz and M. Unsal, JHEP **0208**, 039 (2002) [arXiv:hep-ph/0206102].

[6] I. Antoniadis, A. Delgado, K. Benakli, M. Quiros and M. Tuckmantel, Phys. Lett. B **634**, 302 (2006) [arXiv:hep-ph/0507192].

[7] I. Antoniadis, K. Benakli, A. Delgado, M. Quiros and M. Tuckmantel, Nucl. Phys. B **744**, 156 (2006) [arXiv:hep-th/0601003].

[8] I. Antoniadis, K. Benakli, A. Delgado and M. Quiros, Adv. Stud. Theor. Phys. **2**, 645 (2008) [arXiv:hep-ph/0610265].

[9] K. Hsieh, Phys. Rev. D **77**, 015004 (2008) [arXiv:0708.3970 [hep-ph]].

[10] G. D. Kribs, E. Poppitz and N. Weiner, Phys. Rev. D **78**, 055010 (2008) [arXiv:0712.2039 [hep-ph]].

[11] S. D. L. Amigo, A. E. Blechman, P. J. Fox and E. Poppitz, JHEP **0901**, 018 (2009) [arXiv:0809.1112 [hep-ph]].

[12] A. E. Blechman, "R-symmetric Gauge Mediation and the Minimal R-Symmetric Supersymmetric Mod. Phys. Lett. A **24**, 633 (2009) [arXiv:0903.2822 [hep-ph]].

[13] K. Benakli and M. D. Goodsell, Nucl. Phys. B **816**, 185 (2009) [arXiv:0811.4409 [hep-ph]].

[14] G. Belanger, K. Benakli, M. Goodsell, C. Moura and A. Pukhov, JCAP **0908**, 027 (2009) [arXiv:0905.1043 [hep-ph]].

[15] K. Benakli and M. D. Goodsell, Nucl. Phys. B **830**, 315 (2010) [arXiv:0909.0017 [hep-ph]].

[16] E. J. Chun, J. C. Park and S. Scopel, JCAP **1002**, 015 (2010) [arXiv:0911.5273 [hep-ph]].

[17] K. Benakli and M. D. Goodsell, Nucl. Phys. B **840**, 1 (2010) [arXiv:1003.4957 [hep-ph]].

[18] L. M. Carpenter, arXiv:1007.0017 [hep-th].

[19] Z. Chacko, P. J. Fox and H. Murayama, Nucl. Phys. B **706**, 53 (2005) [arXiv:hep-ph/0406142].

[20] S. Y. Choi, M. Drees, J. Kalinowski, J. M. Kim, E. Popenda and P. M. Zerwas, Phys. Lett. B **672**, 246 (2009) [arXiv:0812.3586 [hep-ph]].

[21] S. Y. Choi, M. Drees, J. Kalinowski, J. M. Kim, E. Popenda and P. M. Zerwas, Acta Phys. Polon. B **40**, 1947 (2009) [arXiv:0902.4706 [hep-ph]].

[22] S. Y. Choi, D. Choudhury, A. Freitas, J. Kalinowski and P. M. Zerwas, Phys. Lett. B **697**, 215 (2011) [arXiv:1012.2688 [hep-ph]].

[23] S. Y. Choi, D. Choudhury, A. Freitas, J. Kalinowski, J. M. Kim and P. M. Zerwas, "Dirac Neutralinos and Electroweak Scalar Bosons of N=1/N=2 Hybrid JHEP **1008**, 025 (2010) [arXiv:1005.0818 [hep-ph]].

[24] A. De Simone, V. Sanz and H. P. Sato, Phys. Rev. Lett. **105**, 121802 (2010) [arXiv:1004.1567 [hep-ph]].

[25] G. D. Kribs, A. Martin and T. S. Roy, JHEP **0901**, 023 (2009) [arXiv:0807.4936 [hep-ph]].

[26] S. Abel and M. Goodsell, arXiv:1102.0014 [hep-th].

[27] R. Davies, J. March-Russell and M. McCullough, arXiv:1103.1647 [hep-ph].

[28] S. P. Martin, arXiv:hep-ph/9709356.

[29] U. Ellwanger, C. Hugonie and A. M. Teixeira, Phys. Rept. **496**, 1 (2010) [arXiv:0910.1785 [hep-ph]].

[30] C. Amsler *et al.* [Particle Data Group], Phys. Lett. B **667** (2008) 1.

[31] R. Barbieri, L. J. Hall, Y. Nomura and V. S. Rychkov, Phys. Rev. D **75**, 035007 (2007) [arXiv:hep-ph/0607332]. L. Cavicchia, R. Franceschini and V. S. Rychkov, Phys. Rev. D **77**, 055006 (2008) [arXiv:0710.5750 [hep-ph]]. R. Franceschini and S. Gori, arXiv:1005.1070 [hep-ph].

[32] P. H. Chankowski, S. Pokorski and J. Wagner, Eur. Phys. J. C **50** (2007) 919 [arXiv:hep-ph/0605302].

[33] M. C. Chen, S. Dawson and C. B. Jackson, Phys. Rev. D **78** (2008) 093001 [arXiv:0809.4185 [hep-ph]].

[34] W. Skiba, arXiv:1006.2142 [hep-ph].