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

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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 $S$ that pairs up with the bino, to form a Dirac fermion, is used to generate $\mu$ 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 $\frac{\kappa^3}{3}S^3$ and discuss the implications for the spectrum of the model.

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1 Introduction

The supersymmetric extension of the standard model introduces new charged particles that need to 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 insufficiently 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 $S$ for $U(1)_Y$, a triplet $T$ for $SU(2)_W$ and an octet $O_g$ 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 $\mu$ 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 of its issues. One is the so-called $\mu$-problem and it is the main motivation of the NMSSM (see [29] and references therein). There one starts by a vanishing tree-level $\mu$-term, as it is forbidden by a Peccei-Quinn symmetry, and adds to the MSSM a singlet with coupling $\lambda S H_u \cdot H_d$. 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 $\lambda S H_u \cdot H_d$, the lightest Higgs mass is now bounded, at tree level, to be:

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

where $g'$ and $g$ are the gauge couplings of the hypercharge $U(1)_Y$, and the weak $SU(2)_W$ respectively, while $\beta$ 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 $\lambda$, 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 $S$ 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 $\kappa S^3$ may indeed allow the generation of both $\mu$ 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 $S$. 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 $Z_3$ 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 $Z_3$ in its couplings to $S$.

The interactions between the different sectors will lead to a collective breaking of supersymmetry, the $U(1)_R$ and the $Z_3$ symmetries. The singlet $S$ 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 $Z_3$ symmetry by the $\lambda S H_u \cdot H_d$ and Dirac gaugino couplings. On the other hand, its coupling to the $Z_3$ symmetry preserving sector is assumed to give rise to a superpotential term $\frac{4}{3}S^3$ 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 $\kappa$. An important constraint on such an extension of the MSSM comes from the $\rho$ 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 $\mu$ 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:

\begin{align}
S &= S + \sqrt{2}\theta \chi_S + \cdots \\
T &= T + \sqrt{2}\theta \chi_T + \cdots \\
O_\xi &= O_\xi + \sqrt{2}\theta \chi_\xi + \cdots
\end{align}

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

\begin{align}
T^{(1)} &= T_1 \sigma^1, \\
T^{(2)} &= T_2 \sigma^2, \\
T^{(3)} &= T_0 \sigma^3, \\
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), \\
T_+ &= \frac{1}{\sqrt{2}}(T_{+R} + iT_{+I}), \\
T_- &= \frac{1}{\sqrt{2}}(T_{-R} + iT_{-I}).
\end{align}

and $\sigma^a$ are the Pauli matrices.

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

\begin{equation}
L^{\text{Dirac}}_{\text{gaugino}} = \int d^2\theta \left[ \sqrt{2}m^{aD}_{1}\tilde{W}_{1a}S + 2\sqrt{2}m^{aD}_{2}\text{tr}(W_{2a}T) + 2\sqrt{2}m^{aD}_{3}\text{tr}(W_{3a}O_\xi) \right] + \text{h.c.}
\end{equation}

where we have introduced spurion superfields

\begin{equation}
m_{aiD} = \theta_{a}m_{iD}.
\end{equation}

The integration over the Grassmannian coordinates leads to

\begin{equation}
\int d^2\theta 2\sqrt{2}m_D \theta^a \text{tr}(W_a \Sigma) \supset -m_D(\lambda_a \psi_a) + \sqrt{2}m_D \Sigma_a D_a
\end{equation}

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

\begin{equation}
L \supset -m_{bD} \sqrt{2}g_b \Sigma_a \phi^a_i R_b^a \phi
\end{equation}
The Higgs sector of the model is described by the superpotential:

\[ W(s) = \mu H_u \cdot H_d + \lambda_S S H_u \cdot H_d + 2\lambda_T H_d \cdot T H_u + \frac{M_S}{2} S^2 + \frac{1}{3} S^3 + M_T \text{tr}(TT) \]
\[ + M_O \text{tr}(O_g O_g), \]  
(2.9)

the Higgs soft masses

\[ m_{H_u}^2 |H_u|^2 + m_{H_d}^2 |H_d|^2 + [\tilde{B}_\mu H_u \cdot H_d + h.c.] \]  
(2.10)

as well as soft terms involving the DG-adjoint fields

\[ -\Delta L_{\text{soft}} = m_S^2 |S|^2 + \frac{1}{2} B_S (S^2 + h.c.) + 2m_T^2 \text{tr}(T^\dagger T) + B_T \text{tr}(TT) + h.c. \]
\[ + A_S \lambda_S S H_u \cdot H_d + 2A_T \lambda_T H_d \cdot T H_u + \frac{1}{3} \kappa A_S S^3 \]
\[ + 2m_O^2 \text{tr}(O_g^4 O_g) + B_O \text{tr}(O_g O_g) + h.c. \]  
(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:

\[ W_2^{(s)} = L S + \lambda_{ST} \text{Str}(TT) + \lambda_{SO} \text{Str}(O_g O_g) + \frac{\kappa_O}{3} \text{tr}(O_g O_g O_g). \]  
(2.12)

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

\[ -\Delta L_{\text{soft}}^A = t^S S + \lambda_{ST} A_{ST} \text{Str}(TT) + \lambda_{SO} A_{SO} \text{Str}(O_g O_g) + \frac{1}{3} \kappa_O A_{SO} \text{tr}(O_g O_g O_g) + h.c. \]  
(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]:

\[ H_u^0 = \frac{H_{uR}^0 + iH_{ul}^0}{\sqrt{2}} = \frac{1}{\sqrt{2}} [c_\beta (v + h) + c_\beta H + i(c_\beta A - s_\beta G^0)], \]  
(3.1)
\[ H_d^0 = \frac{H_{dR}^0 + iH_{dl}^0}{\sqrt{2}} = \frac{1}{\sqrt{2}} [s_\beta (v + h) - s_\beta H + i(s_\beta A + c_\beta G^0)], \]  
(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:

\[ c_\beta \equiv \cos \beta, \quad s_\beta \equiv \sin \beta, \quad t_\beta \equiv \tan \beta \]
\[ c_{2\beta} \equiv \cos 2\beta, \quad s_{2\beta} \equiv \sin 2\beta \]  
(3.3)

\(^1\)Note there are no terms \( \text{tr}(T), \text{tr}(O_g), \text{tr}(TT) \) since these vanish by gauge invariance.
Within these conventions, $v \simeq 246$ GeV, $(g')^2 + g^2 v^2 = M_Z^2$. We are interested by the case of CP neutral vacuum, i.e. $H^0_{uI} = H^0_{aI} = 0$ which implies $S_I = T_I = 0$ \cite{14}. The CP-even singlet and neutral component of the triplet may also acquire a vacuum expectation value, so we define

$$S = \frac{1}{\sqrt{2}} (v_s + S_R + i S_I) \quad T^0 = \frac{1}{\sqrt{2}} (v_T + T_R + iT_I)$$  \hspace{1cm} (3.4)$$

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

$$\tilde{\mu} = \mu + \frac{1}{\sqrt{2}} (\lambda S v_S + \lambda_T v_T)$$
$$\tilde{B}_\mu = B_\mu + \frac{\lambda_S}{\sqrt{2}} (M_S + A_S) v_S + \frac{\lambda_T}{\sqrt{2}} (M_T + A_T) v_T + \frac{1}{2} \lambda_S \kappa v_S^2$$ \hspace{1cm} (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:

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

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

$$\tilde{m}_{SR}^2 = M_S^2 + m_S^2 + 4m_{1D}^2 + B_S$$
$$\tilde{m}_{TR}^2 = M_T^2 + m_T^2 + 4m_{2D}^2 + B_T$$ \hspace{1cm} (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 \cite{14}. The coefficients of the corresponding quadratic terms:

$$\tilde{m}_{SI}^2 = M_S^2 + m_S^2 - B_S$$
$$\tilde{m}_{TI}^2 = M_T^2 + m_T^2 - B_T$$ \hspace{1cm} (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:

$$\tilde{\mu}^2 + \frac{M_\lambda^2}{2} = \frac{m_{H_u}^2 - t_{2\beta} m_{H_u}^2}{t_{2\beta} - 1} + \left[ \frac{t_{2\beta}^2 + 1}{t_{2\beta}^2 - 1} \right] (gm_{2D} v_T - g'm_{1D} v_S)$$ \hspace{1cm} (3.9)$$

and

$$M_A^2 = \frac{2\tilde{B}_\mu}{s_{2\beta}} = 2\tilde{\mu}^2 + m_{H_u}^2 + m_{H_d}^2 \lambda_S^2 + \lambda_T^2 \frac{v^2}{2} c_{2\beta}^2$$ \hspace{1cm} (3.10)$$
where:
\[ v_T \simeq \frac{v^2}{2\tilde{m}_{T R}^2} \left[ -gm_{2D}c_{2\beta} - \sqrt{2}\tilde{\mu}_{T} + \frac{\lambda_T}{\sqrt{2}}(M_T + A_T)s_{2\beta} \right], \] (3.11)
while \( v_S \) is determined as a solution for the cubic equation:
\[ 0 = \kappa^2 v^3_S + \frac{\kappa}{\sqrt{2}}(A_\kappa + 3M_S) v^2_S - \tilde{m}_{S0}^2 v_S + v_0^3 \] (3.12)
under the assumption \( \tilde{m}_{S0}^2, \tilde{m}_{S0}^2 \gg \lambda_S \lambda_T v^2 \) with
\[ \tilde{m}_{S0}^2 = -\tilde{m}_{SR}^2 - \lambda_S^2 \frac{v^2}{2} + \kappa \lambda_S v^2 \] (3.13)
and
\[ \theta^3 = -\frac{v^2}{2} \left[ g'm_{1DC_{2\beta}} - \lambda_S \left( \sqrt{2}\mu - \frac{A_\kappa + M_S}{\sqrt{2}} s_{2\beta} + \lambda_T v T \right) \right]. \] (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:
\[ \tilde{m}_{S}^2 = -\tilde{m}_{S0}^2 + 3\kappa^2 v^2_S + \frac{\sqrt{2}}{3} \kappa v_S (A_\kappa + 3M_S) \] (3.15)
\[ = \tilde{m}_{S0}^2 + \lambda_S^2 \frac{v^2}{2} - \kappa \lambda_S v^2 s_{2\beta} + 3\kappa^2 v^2_S + \frac{\sqrt{2}}{3} \kappa v_S (A_\kappa + 3M_S) \] (3.16)
while it is
\[ \tilde{m}_{T}^2 = \tilde{m}_{TR}^2 + \lambda_T^2 \frac{v^2}{2} \] (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:
\[
\begin{pmatrix}
M^2_h + \Delta h s^2_{2\beta} & \Delta h s_{2\beta} c_{2\beta} & \Delta h s & \Delta h t \\
\Delta h s_{2\beta} c_{2\beta} & M^2_A - \Delta h s^2_{2\beta} & \Delta h s & \Delta h t \\
\Delta h s & \Delta h s & \tilde{m}_{S}^2 & \lambda_S \lambda_T \frac{v^2}{2} \\
\Delta h t & \Delta h t & \lambda_S \lambda_T \frac{v^2}{2} & \tilde{m}_{T}^2
\end{pmatrix}
\] (3.18)
where we have introduced the compact notation:
\[ \Delta h = \frac{v^2}{2}(\lambda_S^2 + \lambda_T^2) - M^2_Z \] (3.19)
which vanishes when \( \lambda_S \) and \( \lambda_T \) 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 \( h \):
\[ \Delta h_s = -2\frac{v_S}{v} \tilde{m}_{S}^2 - \sqrt{2}\kappa \frac{v^2}{v} (A_\kappa + 3M_S) - 2\kappa^2 \frac{v^3}{v}, \quad \Delta h_t = -2\frac{v_T}{v} \tilde{m}_{T}^2 \] (3.20)
while
\[ \Delta_{H_s} = g' m_1 v s_{2\beta} - \lambda_s v (A_s + M_s) \left( \frac{\sqrt{2}}{2} c_{2\beta} \right), \quad \Delta_{H_t} = -g m_2 v s_{2\beta} - \lambda_T v (A_T + M_T) \left( \frac{\sqrt{2}}{2} c_{2\beta} \right) \]

(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:
\[ m_{h}^2 \leq M_Z^2 + \frac{v^2}{2} \left( \frac{\lambda_S^2}{2} + \frac{\lambda_T^2}{2} - \lambda_S v \sqrt{2} \right), \]

(3.22)

### 3.3 Masses of the CP odd neutral scalars

The CP odd neutral scalars mass matrix is given, in the basis \( \{ A, S, T \} \), by
\[
\begin{pmatrix}
M_P^2 & -\lambda_S v \frac{(M_S - A_S)}{\sqrt{2}} + k v_s & -\lambda_T v \frac{(A_T - M_T)}{\sqrt{2}} \\
-\lambda_S v \frac{(M_S - A_S)}{\sqrt{2}} & m_{aS}^2 & \lambda_S v \frac{v_s}{2} \\
-\lambda_T v \frac{(A_T - M_T)}{\sqrt{2}} & \lambda_S v \frac{v_s}{2} & m_{T}^2 + \lambda_T v^2 \frac{v_s}{2}
\end{pmatrix}
\]

(3.23)

where
\[ m_{aS}^2 = m_{SI}^2 + \lambda_S v^2 + \kappa \lambda_S v^2 s_{2\beta} + \kappa^2 v_s^2 + \sqrt{2} k v_s (M_S - A_S) \]

(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:
\[ G_T^\pm = \frac{1}{\sqrt{\rho}} (G^\pm + \sqrt{2} T^T (T_\pm + (T_\mp)^*)) \]

(3.25)

where the erstwhile Goldstone boson (before the triplet coupling is switched on) is
\[ G^+ = c_\beta \overline{H}_d^--s_\beta H_u^+ \]
\[ G^- = c_\beta H_d^- - s_\beta \overline{H}_u^+ \]

(3.26)

the tree-level \( \rho \)-parameter given by:
\[ \rho = 1 + 4 \frac{v_T^2}{v^2} \]

(3.27)

and the remaining orthogonal combinations of charged states can be written as
\[ T_T^+ = \frac{1}{\sqrt{2}} (T_+ + (T_-)^*) \quad T_T^- = (T_T^+)^* \]
\[ T_{RG}^+ = \frac{1}{\sqrt{\rho}} (\frac{1}{\sqrt{2}} (T_+ + (T_\mp)^*) - 2 \frac{v_T}{v} G^+) \quad T_{RG}^- = (T_{RG}^+)^* . \]

(3.28)

(3.29)
The mass matrix in the Lagrangian \( \mathcal{L}_{\text{gaugino}}^{\text{Dirac}} = - (\Sigma^+)^T M_C \Sigma^- \) with \( \Sigma^+ = \{ H^+, T_I^+, T_{RG}^+ \} \) and \( \Sigma^- = \{ H^-, T_I^-, T_{RG}^+ \} \) reads then:

\[
\begin{pmatrix}
\tilde{m}_{H^+} \\
\left( \frac{g^2 - 2 \lambda^2}{2} \nu T s_{2\beta} - \frac{\lambda_T}{\sqrt{2}} (M_T - A_T) \right) \nu \\
-\sqrt{\rho} \Delta t
\end{pmatrix}
\begin{pmatrix}
\nu \\
-\sqrt{\rho} \Delta t
\end{pmatrix}
\begin{pmatrix}
\tilde{m}_{H^+}^2 \\
\tilde{m}_{H^+}^2 + \lambda_T^2 \nu^2 + g^2 \nu_T^2 \\
-\sqrt{\rho} \frac{g^2 - 2 \lambda^2}{4} \nu^2 c_{2\beta}
\end{pmatrix}
\]

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

\[
\tilde{m}_{H^\pm}^2 = M_T^2 c_W^2 + M_A^2 + \left( \lambda_T^2 - \lambda_S^2 \right) \nu^2 + 2 g m_2 D c_{2\beta} v_T
+ 2 \lambda_T^2 \nu_T^2 - 2 \sqrt{2} \lambda_T v_T \bar{\mu} - \sqrt{2} \lambda_T (M_T + A_T) v_T.
\]

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 \( \kappa \).

4 The SU(2) \(_W^W\) 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:

\[
v_T^2 = \frac{v^2}{2(M_T^2 + m_T^2 + 4 m_2^2 D + B_T)} \left[ - g m_2 D c_{2\beta} - \sqrt{2} \bar{\mu} \lambda_T + \frac{\lambda_T}{\sqrt{2}} (M_T + A_T) s_{2\beta} \right]
\]

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

\[
\Delta \rho \approx 4 \frac{v_T^2}{v^2} \lesssim 8 \cdot 10^{-4}
\]

which is satisfied for \( v_T \lesssim 3 \text{ GeV} \).

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 \( M_T \rightarrow \infty \), 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 \( m_2 D \) large, of the order of \( \gtrsim 2 \text{ TeV} \). 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 \( m_2 D \rightarrow \infty, m_T \rightarrow \infty \).
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:

$$V_{EW} = \left[ \frac{g^2}{4} + \frac{\lambda^2_S + \lambda^2_T}{2} \right] \frac{v^4}{8}$$

$$+ \left[ m_{H_u}^2 s^2_\beta + m_{H_d}^2 c^2_\beta + \mu^2 - \tilde{B}_\mu s_{2\beta} - g' m_{1D} v S c_{2\beta} \right] \frac{v^2}{2}$$

$$+ \frac{\kappa^2}{4} v_S^4 + \frac{\kappa}{\sqrt{2}} \frac{(3 M_S + A_S)}{3} v_S^3 + \frac{1}{2} \tilde{m}^2_{sR} v_S^2$$  \hspace{1cm} (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:

$$V_{EW} = \left[ \frac{g^2 + g'^2}{4} c^2_\beta + \frac{\lambda^2_S + \lambda^2_T}{2} \right] \frac{v^4}{8}$$

$$+ \left[ m_{H_u}^2 s^2_\beta + m_{H_d}^2 c^2_\beta + \mu^2 - \tilde{B}_\mu s_{2\beta} - g' m_{1D} v S c_{2\beta} \right] \frac{v^2}{2}$$

$$+ \frac{\kappa^2}{4} v_S^4 + \frac{\kappa}{\sqrt{2}} \frac{(3 M_S + A_S)}{3} v_S^3 + \frac{1}{2} \tilde{m}^2_{sR} v_S^2$$  \hspace{1cm} (4.4)$$

where we see that both contributions from the $SU(2)_W$ $D$-term and $\lambda_T$ 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:

$$m_h^2 \leq M_Z^2 c^2_\beta + \frac{v^2}{2} (\lambda^2_S + \lambda^2_T) s^2_{2\beta}. \hspace{1cm} (5.1)$$

In order to obtain a spectrum with $m_h > M_Z$, we need both a large $\lambda^2_S + \lambda^2_T$ and $s^2_{2\beta} \to 1$. While the second condition can be satisfied quite easily, the first one requires the presence of new physics at energies $\Lambda_{NP}$ of a few tens of TeV, and may be subject to constraints from electroweak precision tests. In fact, as the couplings $\lambda_T, \lambda_S$ grow with energy, and now start with a large value at the electroweak scale, they become non-perturbative in the ultraviolet very rapidly at the scale $\Lambda_{NP}$. 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:

\[
16\pi^2 \frac{d}{dt}\lambda_S = \lambda_S \left[ 4\lambda_S^2 + 6\lambda_T^2 + 2|\kappa_S|^2 + \cdots \right]
\]

\[
16\pi^2 \frac{d}{dt}\lambda_T = \lambda_T \left[ 2\lambda_S^2 + 8\lambda_T^2 + \cdots \right]
\]

\[
16\pi^2 \frac{d}{dt}\kappa_S = \kappa_S \left[ 6\lambda_S^2 + 6|\kappa_S|^2 \right].
\]

(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 \( \lambda_S \) is chosen to be large enough while \( \lambda_T \ll \lambda_S \).
   A bound on the size of the \( \lambda_S^2 \) 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 \( \lambda_S \) at low scale ranges from approximately \( \lambda_S \sim 1.4 \) for \( \Lambda_{NP} \sim 50 \text{ TeV} \), to \( \lambda_S \sim 2 \) for \( \Lambda_{NP} \sim 10 \text{ TeV} \).

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 \( \Lambda_{NP} \sim 10–50 \text{ TeV} \).

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 \( \Lambda_{NP} \sim 10–50 \text{ TeV} \).

Note that in order to allow the largest values of \( \lambda_S \), one needs to keep \( \kappa \) small as it contribute to the corresponding RGEs.

First, let us assume the presence of both \( \mu \) and \( B_{\mu} \) terms not generated by the singlet vev \( v_S \). The alternative possibility shall be investigated in the next section. The \( \mu \) term by itself does not break \( R \)-symmetry, but then \( B_{\mu} \) does. Expecting the same source of \( R \)-symmetry breaking to provide contributions to other soft terms, requires then to assume \( B_{\mu} \) 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:

\[
\begin{pmatrix}
M_Z^2 s_W^2 + \Delta_h s_{2\beta}^2 & \Delta_h s_{2\beta} c_{2\beta} \\
\Delta_h s_{2\beta} c_{2\beta} & M_A^2 - \Delta_h s_{2\beta}^2
\end{pmatrix}
\]

(5.3)

which has as eigenvalues:

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

(5.4)

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

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

(5.5)
6 The singlet vacuum expectation value

A scalar component $S$ 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 \cite{31}). For simplicity (and in line with our philosophy of $\kappa$ 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 \cite{14}, one gets:

$$v_S \simeq \frac{\tilde{m}_{S0}^2}{m_S^2} \lambda_3 v^2 \left[ \sqrt{2} \mu \lambda_S - g' m_{1D} c_2 \beta \right]. \quad (6.1)$$

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

We would like to switch on the trilinear coupling $\kappa$ 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 $A_\kappa$ smaller than the other relevant mass parameters.

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

$$|v_S| \simeq \frac{\tilde{m}_{S0}^2}{\kappa} + \frac{v_0^2}{2\tilde{m}_{S0}^2} + \cdots \quad (6.2)$$

$$|v_S| \simeq \frac{v}{\sqrt{2\kappa}} \left( \frac{\lambda_3}{v^2} - \frac{\lambda_3^2 + \kappa \lambda_S s_2 \beta}{2m_{SR}^2} \right)^{1/2} \quad (6.3)$$

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

$$-B_S = |B_S| \gg m_{1D} \quad (6.5)$$

This hierarchy is not unnatural, it is even quite generic in models of gauge mediation with Dirac gauginos \cite{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 \cite{13}. Therefore:

$$|v_S| \simeq \frac{\tilde{m}_{S0}/\kappa}{v_0^2}. \quad (6.6)$$

This singlet vev (6.4) induces both a $\mu$-term

$$\mu^2 = \frac{\lambda_3^2}{2} \frac{v_S^2}{2\kappa^2} \simeq \frac{\lambda_3^2}{2\kappa^2} \left( \frac{\lambda_3 - \kappa \lambda_S s_2 \beta}{v_S^2} \right) v^2 \simeq \frac{\lambda_3^2}{\kappa} \left( |B_S| - m_{SR}^2 - 4m_{1D}^2 \right). \quad (6.7)$$
and a $B\mu$-term

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

(6.8)

with sizes controlled by $\lambda_S$, $\kappa$ 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 $\mu^2$ 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 $\frac{2}{3}S^3$ induces through the vev of $S$ a Majorana mass for the singlino $\chi_S$.

The neutralino mass matrix, in the basis $\tilde{S}, \tilde{B}, \tilde{\tilde{T}}, \tilde{W}^0, \tilde{H}_d^0, \tilde{H}_u^0$ reads:

$$
\begin{pmatrix}
M_S + \sqrt{2}k\nu_S & m_{1D} & 0 & 0 & -\frac{\sqrt{2}S}{g} M_{ZSW\beta} & -\frac{\sqrt{2}S}{g} M_{ZSW\beta} \\
m_{1D} & M_1 & 0 & 0 & -\frac{\sqrt{2}S}{g} M_{ZSW\beta} & -\frac{\sqrt{2}S}{g} M_{ZSW\beta} \\
0 & 0 & M_T & m_{2D} & -\frac{\sqrt{2}S}{g} M_{ZSW\beta} & -\frac{\sqrt{2}S}{g} M_{ZSW\beta} \\
0 & 0 & m_{2D} & M_2 & M_{ZCW\beta} & 0 \\
-\frac{\sqrt{2}S}{g} M_{ZSW\beta} & -\frac{\sqrt{2}S}{g} M_{ZSW\beta} & -\frac{\sqrt{2}S}{g} M_{ZSW\beta} & M_{ZCW\beta} & 0 & -\tilde{\mu} \\
-\frac{\sqrt{2}S}{g} M_{ZSW\beta} & M_{ZSW\beta} & M_{ZSW\beta} & M_{ZCW\beta} & -\tilde{\mu} & 0
\end{pmatrix}
$$

(6.9)

Therefore a Majorana mass of order:

$$M'_1 = \sqrt{2}k\nu_S \simeq \frac{2\kappa}{\lambda_S} \tilde{\mu}$$

(6.10)

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

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

(6.11)

Noting (6.7) we see that this requires

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

(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 $v^+ = (\tilde{T}^+, \tilde{W}^+, \tilde{H}_d^+)$, $v^- = (\tilde{T}^-, \tilde{W}^-, \tilde{H}_u^0)$, are given by

$$M_{Ch} = \begin{pmatrix}
M_T & m_{2D} + gv_T & \frac{2\kappa}{g} M_{ZCW\beta} \\
m_{2D} - gv_T & M_2 & \sqrt{2}M_{ZCW\beta} \\
-\frac{2\kappa}{g} M_{ZCW\beta} & \sqrt{2}M_{ZCW\beta} & \tilde{\mu} - \sqrt{2}\lambda_T v_T
\end{pmatrix}$$

(6.13)

and do not depend on $\kappa$. 

13
6.2 The scalar spectrum

The scalar spectrum

The scalar spectrum of a CP even scalar singlet $S_R$ is given by the effective mass:

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

while the one for CP-odd $S_I$ reads:

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

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

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

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

The CP odd neutral scalar mass is given by

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

The charged Higgs mass is

$$m_{H^\pm}^2 = M_W^2 s_{2\beta}^2 + M_A^2 + (\lambda_T^2 - \lambda_S^2) \frac{\nu^2}{2}. \quad (6.19)$$

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

$$\Delta_{hs} \simeq -\sqrt{2}(\lambda_S + \kappa s_{2\beta})\tilde{\mu}v, \quad \Delta_{ht} \simeq \sqrt{2}\lambda_T \tilde{\mu}v + gm_{2D}\nu s_{2\beta} \quad (6.20)$$

$$\Delta_s \simeq g'm_{1D}\nu s_{2\beta}, \quad \Delta_t \simeq -gm_{2D}\nu s_{2\beta} \quad (6.21)$$

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

With both $\mu$ and $B_\mu$ 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 $\lambda_S$. 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 $\lambda_S = 0.8, \kappa = 0.6, \lambda_T = 0.1, \tan \beta = 1.38$ 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 $\lambda_T$ and small $\lambda_S$. 

14
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 approximately 72 percent of bino, 18 percent of singlino, while they become approximately 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 $\lambda_S, \lambda_T$ 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 $\mu$ and $B\mu$ parameters. The $S$ vev, easily computed, was found in all previous studies to be typically too small. As the combination of $\mu$ 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 $\frac{1}{9}S^3$, and shown that it is sufficient. Both $\mu$ 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 $\lambda_S, \lambda_T$ 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 $\lambda_S, \lambda_T$, $m_{2D}$ 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 $\kappa S^3$, and investigation of the resulting pattern of the soft-terms are needed. We will return to these important issues in future work.
### Table 1: Examples of model parameters.

| Input parameter | Model I | Model II | Model III |
|-----------------|---------|----------|-----------|
| $\lambda_S$     | 1.2     | 0.8      | 0.1       |
| $\lambda_T$     | 0.1     | 0.1      | 0.7       |
| $\kappa$        | 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$^2$ |
| $m_T^2$         | $10^6$ GeV$^2$ | $10^6$ GeV$^2$ | $10^6$ GeV$^2$ |
| $B_T$           | 0       | 0        | 0         |
| $A_\kappa$      | 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, 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

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| Input parameter | Model IV |
|-----------------|----------|
| $\lambda_S$     | 0.13     |
| $\lambda_T$     | 0.13     |
| $\kappa$        | 0.13     |
| $\tan \beta$    | 31.8     |
| $m^2_S$         | 400000 GeV$^2$ |
| $B_S$           | -1500000 GeV$^2$ |
| $m^2_T$         | 6000000 GeV$^2$ |
| $B_T$           | 0        |
| $A_{\kappa}$    | 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.

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