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Dirac gauginos, R symmetry and the 125 GeV Higgs

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Abstract: We study a supersymmetric scenario with a quasi exact R-symmetry in light of the discovery of a Higgs resonance with a mass of 125 GeV. In such a framework, the additional adjoint superfields, needed to give Dirac masses to the gauginos, contribute both to the Higgs mass and to electroweak precision observables. We analyze the interplay between the two aspects, finding regions in parameter space in which the contributions to the precision observables are under control and a 125 GeV Higgs boson can be accommodated. We estimate the fine-tuning of the model finding regions of the parameter space still unexplored by the LHC with a fine-tuning considerably improved with respect to the minimal supersymmetric scenario. In particular, sizable non-holomorphic (non-supersoft) adjoints masses are required to reduce the fine-tuning.

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1 Introduction

The discovery of a 125 GeV particle closely resembling the Standard Model (SM) Higgs [1, 2] may represent a challenge for Supersymmetry (SUSY). Indeed, at least in its minimal version, large loop contributions are needed to raise the mass of the lightest Higgs boson to the observed value, the most relevant ones coming from the stop system. This points toward very heavy stops, and/or large left-right stop mixing.

While this is perfectly consistent with the non observation of any superpartner at the LHC, it is widely believed to be at odds with the concept of naturalness, which requires stop masses as well as the corresponding A-terms to be below 1 TeV. Needless to say, after the first LHC run and the Higgs discovery, understanding whether the concept of naturalness as it stands is or not a principle followed by nature has become of the utmost importance.

If we insist on naturalness, we need to consider alternatives to the Minimal Supersymmetric Standard Model (MSSM). An interesting possibility is given by models with Dirac gauginos, which have relaxed naturalness bounds on the gluino mass. This is most welcome since a relaxed naturalness bound on this mass — being the most constrained after the first LHC run — would result in less tension with data. The mechanism behind the improved naturalness is the generation of Dirac gaugino masses through supersoft operators, which give only finite contributions to scalar masses [3]. Models with Dirac gauginos are also interesting from a purely phenomenological point of view: first of all, squark pair production is suppressed at the LHC due to the absence of Majorana mass insertions [4]. Moreover, Dirac gaugino masses are compatible with the presence of a global $\mathrm{U}(1)_R$ symmetry, which would be otherwise broken by the Majorana mass term. The R-symmetry can be used as an alternative to R-parity to forbid operators leading to proton decay [5, 6], but has far

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richer consequences. Indeed, the absence of A terms, the μ term and Majorana gaugino masses has a drastic beneficial effect on the SUSY flavor problem [7].

A peculiar aspect of R-symmetric models is the Higgs sector particle content. Models have been proposed in the literature with four Higgs doublets [7], two Higgs doublets in which the role of the down type Higgs is played by one of the lepton doublets [8], one up type Higgs doublet [9] or even with no Higgs doublets at all, with the role of the Higgs being played by one of the slepton doublets [10].

As already pointed out, naturalness is among the reasons motivating the study of models with Dirac gauginos. However, a solid and complete statement about the fine-tuning cannot be done without a full analysis of how a 125 GeV Higgs mass is obtained within this framework. The situation has been studied in [11], where however the R-symmetric case was not considered.[1] This case is going to be the focus of this paper. As we will explain, respecting the R-symmetry in the Higgs sector changes dramatically how the lightest Higgs mass is raised up to 125 GeV (see [13]). Indeed, while in [11] this is achieved through an NMSSM-like tree level enhancement of the Higgs mass, here this possibility is forbidden by the R-symmetry.[2] However, it turns out that the extra matter necessary to respect the R-symmetry, i.e. the adjoint scalars and the inert doublets, can provide radiative corrections comparable to the stop one, giving a 125 GeV Higgs with a few percent level fine-tuning.

2 Electroweak symmetry breaking in R-symmetric models

As already explained, preserving the R-symmetry typically requires an enlarged Higgs sector. For definiteness, we will present the Lagrangian for the four Higgs doublet model of ref. [7], in which the two doublets with R-charge 0, H_u and H_d, acquire a vev while the two with R-charge 2, R_u and R_d, are inert doublets. Another, more economical, possibility is to have the sneutrino as the down type Higgs so that just two doublets, H_u and R_d are needed.[3] We will focus on the large $\tan\beta$ limit (where $\tan\beta \gtrsim 10$) in which most of electroweak symmetry breaking is through H_u, with the extra Higgs states decoupled from the electroweak symmetry breaking sector. In this limit, we expect the various models to give similar results.

The superpotential of the model is given by:

$$\begin{aligned} W &= W_{\text{Yukawa}} + W_{\text{Higgs}} \\ W_{\text{Yukawa}} &= H_u Q Y_u u^c + H_d Q Y_d d^c + H_d L Y_e e^c, \\ W_{higgs} &= \sqrt{2}\lambda_T^u H_u T R_d + \sqrt{2}\lambda_T^d R_u T H_d + \lambda_S^u H_u S R_d + \lambda_S^d R_u S H_d \\ &\quad + \mu_u H_u R_d + \mu_d R_u H_d \,. \end{aligned} \tag{2.1}$$

[1]In ref. [12] the question of the Higgs mass and fine-tuning is investigated in a scenario with additional right handed neutrinos.

[2]Another possible extension is through new U(1) D-terms, as explored in [14]. This proposal, however, also involves R-symmetry breaking effects.

[3]It is also possible to have an even more economical Higgs sector [10] where the sneutrino gives mass to the up type fermions via SUSY breaking Yukawa couplings. In this case the Higgs quartic is generated by SUSY breaking as well.

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We write the triplet superfield normalized as

$$T = \frac{1}{\sqrt{2}} \begin{pmatrix} T^0 & \sqrt{2}T^+ \\ \sqrt{2}T^- & -T^0 \end{pmatrix} , \tag{2.2}$$

so that the kinetic terms for the (complex) triplet components are automatically canonically normalized; the factor $\sqrt{2}$ in front of λ_T^i is chosen such that $W \supset \lambda_T H_u^0 T^0 R_d^0$.

The R-symmetry allows the gaugino fields λ_i to pair up with the fermionic components of the adjoint superfields, ψ_i, through soft SUSY breaking Dirac masses

$$\mathcal{L}_D = M_{\tilde{B}} \lambda_{\tilde{B}} \psi_{\tilde{B}} + M_{\tilde{W}} \lambda^a_{\tilde{W}} \psi^a_{\tilde{W}} + M_{\tilde{g}} \lambda_{\tilde{g}} \psi_{\tilde{g}} + h.c. \tag{2.3}$$

Moreover, the soft SUSY breaking scalar terms read

$$\begin{aligned} V_{\text{soft}}^{EW} &= \tilde{Q}^\dagger m^2_{\tilde{Q}} \tilde{Q} + \tilde{u}^\dagger m^2_{\tilde{u}} \tilde{u} + \tilde{d}^\dagger m^2_{\tilde{d}} \tilde{d} + \tilde{L}^\dagger m^2_{\tilde{L}} \tilde{L} + \tilde{e}^\dagger m^2_{\tilde{e}} \tilde{e} + B_\mu H_u H_d \\ &+ m^2_{H_u} |H_u|^2 + m^2_{H_d} |H_d|^2 + m^2_{R_u} |R_u|^2 + m^2_{R_d} |R_d|^2 + m^2_S |S|^2 + m^2_T T^{a\dagger} T^a \\ &+ t_S\, S + B_S S^2 + \frac{1}{3} A_S\, S^3 + B_T T^a T^a \\ &+ A_{ST}\, S T^2 + A_{SH}\, S H_u H_d + A_{TH}\, H_u T H_d + h.c. \end{aligned} \tag{2.4}$$

We notice that the R-symmetry forbids all the A-terms except for those written above, which together with t_S we will assume to be negligible for simplicity.[4]

Let us now comment on the soft breaking terms in the adjoint sector. As already explained in the Introduction, Dirac gaugino masses are generated by supersoft operators and give finite contributions to the scalar masses. This property has the beneficial effect of relaxing the gluino naturalness bound, reducing the tension with the direct searches [4]. However not all possible R-invariant terms that can be constructed out of the adjoint superfields are supersoft: indeed the non-holomorphic adjoints masses for the singlet m_S^2, the triplet m_T^2, and the octect m_O^2 contribute at the two-loop level to the β functions for the scalar masses, pushing down their values at low energy. In particular, a too large octect scalar mass can eventually induce tachyonic squark masses, causing charge and color breaking at the weak scale [15, 16]. Furthermore, it is also important for the Dirac gaugino masses and the holomorphic and non-holomorphic adjoints scalar masses to be of the same order, to avoid tachyons already at tree level. It turns out, however, that realizing such a requirement in a UV complete model is quite challenging. This resembles the $\mu - B_\mu$ problem in gauge mediation, and leads to a source of fine-tuning estimated in [16] to be of order of 0.1%. In what follows we will discuss a generic case where also non-holomorphic masses are present, assuming that the mass hierarchy among the adjoint soft terms is such as to ensure color and charge conservation at the weak scale. For concreteness, we will take the gluino and its scalar octect partner to have masses around $4-5$ TeV, i.e. large enough to be safe from any direct search bound. Moreover, we will assume their ratio to be such that the induced tuning on the stop masses is not larger than 20% [4]. We also emphasize that we will focus in this work on the sectors most closely connected to electroweak symmetry breaking (EWSB),

[4]We can in any case invoke a $\mathbb{Z}_2$ parity under which S, T and $R_{u,d}$ are odd to forbid these terms.

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which can affect the Higgs mass in a significant way. Since, for instance, the slepton or squark sectors (other than the stops) depend on parameters that, generally, are not strongly constrained from the considerations in this paper, we will remain agnostic as to their detailed properties. As a result, we do not need to commit to a particular scenario in this regard, which can be rather model-dependent. For instance, our Higgs sector can be embedded in a baryonic RPV scenario [17], which would affect significantly the interpretation of the current LHC bounds. After this clarifications, we turn to the Higgs sector of the model.

The total scalar potential is

$$\begin{aligned} V^{EW} &= V_F^{EW} + V_D^{EW} + V_{\text{soft}}^{EW}\,, \\ V_F^{EW} &= \sum_i \left|\frac{\partial W}{\partial \phi_i}\right|^2\,, \qquad V_D^{EW} = \frac{1}{2}\sum_{a=1}^{3}(D_2^a)^2 + \frac{1}{2}D_Y^2\,, \end{aligned} \tag{2.5}$$

with W defined in eq. (2.1) and V_{soft}^{EW} given in eq. (2.4). The presence of additional chiral superfields charged under $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ modifies the expression for the D-terms:

$$\begin{aligned} D_2^a &= g\left(H_u^\dagger\tau^a H_u + H_d^\dagger\tau^a H_d + R_u^\dagger\tau^a R_u + R_d^\dagger\tau^a R_d + \vec{T}^\dagger\lambda^a\vec{T}\right) + \sqrt{2}M_{\tilde{W}}\left(\vec{T}^a + \vec{T}^{\dagger a}\right), \\ D_Y &= \frac{g'}{2}\left(H_u^\dagger H_u + R_u^\dagger R_u - H_d^\dagger H_d - R_d^\dagger R_d\right) + \sqrt{2}M_{\tilde{B}}\left(S + S^\dagger\right), \end{aligned} \tag{2.6}$$

where $M_{\tilde{B}}$ and $M_{\tilde{W}}$ are the Dirac Bino and Wino masses, τ^a and λ^a are the two and three dimensional SU(2) generators respectively, while $\vec{T}^a = \sqrt{2}\,\mathrm{tr}(\tau^a T) = \left\{\frac{T^+ + T^-}{\sqrt{2}}, \frac{T^- - T^+}{\sqrt{2}i}, T^0\right\}$.

Writing the neutral fields as

$$H_{u,d}^0 = \frac{h_{u,d} + i a_{u,d}}{\sqrt{2}}, \quad R_{u,d}^0 = \frac{r_{u,d} + i a_{r_{u,d}}}{\sqrt{2}}, \quad T^0 = \frac{t + i a_t}{\sqrt{2}}, \quad S = \frac{s + i a_s}{\sqrt{2}}\,, \tag{2.7}$$

the scalar potential for the CP even components reads:

$$\begin{aligned} V = \frac{1}{2}\Big[& \left(m_{H_u}^2 + \mu^2\right)h_u^2 + \left(m_{R_u}^2 + \mu^2\right)r_u^2 + \left(m_{H_d}^2 + \mu^2\right)h_d^2 + \left(m_{R_d}^2 + \mu^2\right)r_d^2 \\ & - 2B_\mu h_u h_d + \left(4M_{\tilde{B}}^2 + m_S^2 + 2B_S\right)s^2 + \left(4M_{\tilde{W}}^2 + m_T^2 + 2B_T\right)t^2\Big] \\ & + \frac{1}{2}\Big[\sqrt{2}\mu\left(\lambda_S s + \lambda_T t\right)\left(h_u^2 + h_d^2 + r_u^2 + r_d^2\right) \\ & + \left(g M_{\tilde{W}} t - g' M_{\tilde{B}} s\right)\left(h_d^2 + r_d^2 - h_u^2 - r_u^2\right)\Big] + \frac{1}{32}\left(g^2 + g'^2\right)\left[\left(h_u^2 - h_d^2\right)^2 + \left(r_u^2 - r_d^2\right)^2\right] \\ & + \frac{g^2 + g'^2}{16}\left(h_u^2 r_u^2 + h_d^2 r_d^2\right) + \left(\frac{\lambda_S^2 + \lambda_T^2}{4} - \frac{g^2 + g'^2}{16}\right)\left(h_u^2 r_d^2 + h_d^2 r_u^2\right) \\ & + \frac{\lambda_S\lambda_T}{2}st\left(h_u^2 + h_d^2 + r_u^2 + r_d^2\right) + \frac{\lambda_T^2 t^2 + \lambda_S^2 s^2}{4}\left(h_u^2 + h_d^2 + r_u^2 + r_d^2\right), \end{aligned} \tag{2.8}$$

where we have assumed for simplicity $\lambda_S^u = \lambda_S^d = \lambda_S$, $\lambda_T^u = \lambda_T^d = \lambda_T$, $\mu_u = \mu_d = \mu$ and set $t_S = A_{ST} = A_{SH} = A_{TH} = 0$. The minimization conditions for this potential are written in the appendix. The triplet acquires a vev which is constrained by EWPM to be $|v_T| \lesssim 3\,\mathrm{GeV}$. We will discuss more precisely the bounds from EWPM in section 3.

Inspecting the various contributions, we notice that the D-terms produce the usual MSSM quartic. However, the Dirac gaugino masses contribute to reduce the tree level Higgs mass with respect to the MSSM. Indeed, V_D contains trilinear interactions between the active Higgs fields (those participating in EWSB) and the scalar adjoints,

$$V_D \supset \frac{1}{2}\left(-gM_{\tilde{W}}t + g'M_{\tilde{B}}s\right)h_u^2 - \frac{1}{2}\left(-gM_{\tilde{W}}t + g'M_{\tilde{B}}s\right)h_d^2\,, \tag{2.9}$$

which after EWSB push down the lightest eigenvalue due to mixing.

In addition, the R-symmetry forces the active Higgs fields to couple only with the inert doublets (those that do not get vevs) and not among themselves, so that any NMSSM-like quartic term $\lambda_{S,T}^2 h_u^2 h_d^2$ is forbidden. As a consequence, the MSSM tree level upper bound $(m_h)^2_{\text{tree}} \leq m_Z^2 \cos^2 2\beta$ applies, and the lightest scalar mass is maximized in the large $\tan\beta$ regime. The situation is different when the R-symmetry is broken in the Higgs sector. In this case $W \supset \lambda_T H_u T H_d + \lambda_S H_u S H_d$ and in the low $\tan\beta$ regime the usual NMSSM-like tree level enhancement is recovered [11].

A more complete discussion of the tree level scalar masses will be presented in section 2.1, where in order to maximize the lightest eigenvalue we will focus on the large $\tan\beta$ regime.[5] In section 2.2 we will instead study the loop corrections to the Higgs boson mass.

2.1 Tree level Higgs mass

We have already pointed out that Dirac gaugino masses constitute an irreducible source of mixing between active Higgs fields and adjoint scalars, which tends to lower the smallest mass eigenvalue in the CP-even Higgs sector at tree level. Appealing to loop-level corrections to reproduce the observed $m_h \approx 125$ GeV could then require a heavier SUSY spectrum, with the associated increase in fine-tuning. However, we note that the above push-down effect may be minimized by additional contributions from the supersymmetric couplings λ_T, λ_S and the μ term. This can be seen from the off-diagonal elements of the mass matrix for CP even scalars (see appendix A for the complete expressions):

$$\begin{aligned} m^2_{h_u,t} &= v(-\sqrt{2}gM_{\tilde{W}} + 2\lambda_T(\lambda_S v_S + \lambda_T v_T + \mu))\,, \\ m^2_{h_u,s} &= v(+\sqrt{2}g'M_{\tilde{B}} + 2\lambda_S(\lambda_S v_S + \lambda_T v_T + \mu))\,. \end{aligned} \tag{2.10}$$

Anticipating that λ couplings of order one are helpful to increase the Higgs boson mass at loop level when the stops are not too heavy (as will be studied in section 2.2), and insisting on relatively small μ values as suggested by naturalness (see section 4), we see that the terms in eq. (2.10) can be smaller than naively expected when the singlet and triplet vevs are small. This follows from a possible partial cancellation between the first and last terms in eqs. (2.10) (those involving the Dirac gaugino masses and the μ-term, respectively). Since by field redefinitions we can always choose $g > 0$ and $M_{\tilde{B}}, M_{\tilde{W}}, \mu > 0$, we conclude that $\lambda_T > 0$ and $\lambda_S < 0$ are preferred to obtain smaller $m^2_{h_u,t}$ and $m^2_{h_u,s}$. This is confirmed in figure 1, where we show the tree level Higgs boson mass, together with the singlet vev

[5]However, we have checked that it is easy to deform our benchmark points to obtain examples with moderate $\tan\beta$ (~ 10) without affecting our conclusions.

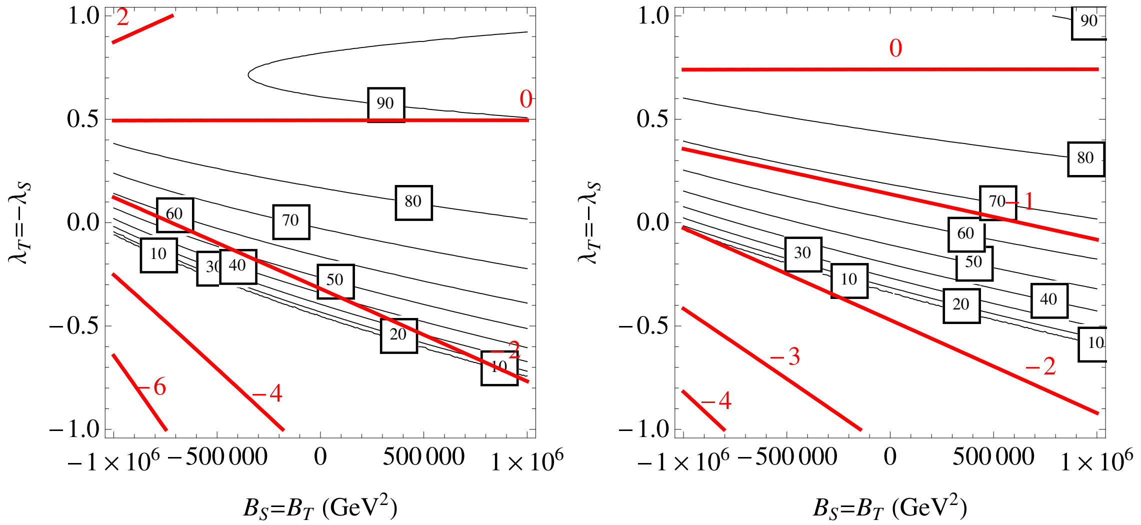

Figure 1. Tree level Higgs boson mass in GeV (black lines) and singlet vev v_S (red lines) as a function of $B_T = B_S$ and $\lambda_T = -\lambda_S$. Left: $M_{\tilde{W}} = M_{\tilde{B}} = 600\,\text{GeV}$, $m_T = m_S = 1500\,\text{GeV}$ and $\mu = 300\,\text{GeV}$. Right: $M_{\tilde{W}} = M_{\tilde{B}} = 900\,\text{GeV}$, $m_T = m_S = 1500\,\text{GeV}$ and $\mu = 300\,\text{GeV}$.

v_S, as a function of $B_T = B_S$ and $\lambda_T = -\lambda_S$. Apart from showing that the region $\lambda_T > 0$ is preferred, we also see that one gets $m_h \simeq m_Z$ when v_S is small and positive. Thus, it is possible to start from a tree-level Higgs mass not too far from the MSSM limit, even in the presence of Dirac gaugino masses, provided the couplings λ_T and λ_S are sizable.

To illuminate the above discussion, it is useful to derive a simple formula for the smallest CP-even Higgs mass eigenvalue in the limit of small v_T, v_S as well as a large hierarchy between Dirac gaugino masses and non-holomorphic adjoint masses, $M_D \ll m_{\text{adj}}$. For $\tan\beta \gg 1$, the lightest tree-level mass is:

$$(m_h^2)_{\text{tree}} \simeq m_Z^2 - v^2 \frac{(-\sqrt{2} g M_{\tilde{W}} + 2\lambda_T \mu)^2}{m_{T_R}^2} - v^2 \frac{(\sqrt{2} g' M_{\tilde{B}} + 2\lambda_S \mu)^2}{m_{S_R}^2}\,, \tag{2.11}$$

where $m_{T_R}^2 = 4M_{\tilde{W}}^2 + m_T^2 + 2B_T$ and $m_{S_R}^2 = 4M_{\tilde{B}}^2 + m_S^2 + 2B_S$ are the masses of the real parts of the adjoint scalars before EWSB. Eq. (2.11) can be used as a first estimate for the tree-level Higgs mass in our region of interest. Let us also stress that the presence of supersymmetric couplings, as well as holomorphic and non-holomorphic masses for the adjoint scalars, improves the situation with respect to [3], where the quartic coupling vanishes for decoupled adjoint scalars (see also [11]).

2.2 Radiative corrections to the Higgs mass

The 1-loop level corrected Higgs mass can be computed by adding the Coleman-Weinberg potential:

$$V_{\text{Higgs}}^{CW} = \frac{1}{64\pi^2}\left[\sum_i (-1)^{2J_i+1}\,(2J_i+1)\,m_i^4 \left(\log\frac{m_i^2}{Q^2} - \frac{3}{2}\right)\right]\,, \tag{2.12}$$

to the tree-level potential in eq. (2.8). Here the sum is to be taken over all the states coupled to the Higgs, with m_i^2's the field-dependent masses. In section 4 we will perform a numerical

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analysis of the 1-loop corrected Higgs mass using eq. (2.12). However, it will be illuminating to obtain also analytic expressions for these loop corrections, at least in certain limits. For instance, we will be interested in a region of parameter space where the triplet vev is small (due to EW constraints) and, as motivated in the previous subsection, also where the singlet vev is small (in order to maximize the Higgs mass and reduce the tuning). To obtain simple analytic expression, we therefore set to zero both the singlet and triplet vevs. We also diagonalize perturbatively the h_u-dependent mass matrices for the scalar and fermionic sectors (recall that we are focusing on the large $\tan\beta$ limit where h_d plays a negligible role), and expand the resulting 1-loop potential in powers of h_u. The resulting analytical expressions are still fairly cumbersome, and not particularly transparent. However, simple expressions can be obtained for $M_D \ll m_{\rm adj}$ or $m_{\rm adj} \ll M_D$, where M_D and $m_{\rm adj}$ are common mass scales for Dirac gauginos and adjoint scalars, respectively. These limits turn out to be useful to understand the numerical results to be presented later (and which do not assume extreme hierarchies between these parameters). In addition, they will serve to further motivate the preferred choices for parameters such as λ_T and λ_S. For this reason we will focus in the following discussion on the contributions associated with these two couplings. In the numerical analysis to be presented in section 4 we will include also subdominant effects such as those coming from the gauge interactions.

Region 1. When the scalar CP even and CP odd masses are significantly larger than the gaugino masses, $\mu \ll M_D \ll m_{\rm adj}$, we have the following contribution to the Higgs quartic coupling:

$$\begin{aligned} V^{CW}_{\rm Higgs} \supset \frac{1}{4}\Bigg[& \frac{5\lambda_T^4+2\lambda_T^2\lambda_S^2+\lambda_S^4}{32\pi^2}\log\frac{m^2_{R_d}}{Q^2}+\frac{\lambda_T^2}{32\pi^2}\left(5\lambda_T^2+2\lambda_S^2\frac{m_T^2}{m_T^2-m_S^2}\right)\log\frac{m_T^2}{Q^2} \\ & +\frac{\lambda_S^2}{32\pi^2}\left(\lambda_S^2-2\lambda_T^2\frac{m_S^2}{m_T^2-m_S^2}\right)\log\frac{m_S^2}{Q^2}-\frac{\lambda_T^2\lambda_S^2}{16\pi^2}\Bigg]h_u^4 \\ & -\frac{1}{4}\Bigg[\frac{\lambda_T^2}{16\pi^2}\left(5\lambda_T^2+2\lambda_S^2\frac{M^2_{\tilde W}}{M^2_{\tilde W}-M^2_{\tilde B}}\right)\log\frac{M^2_{\tilde W}}{Q^2}-\frac{\lambda_S^2\lambda_T^2}{8\pi^2} \\ & +\frac{\lambda_S^2}{16\pi^2}\left(\lambda_S^2-2\lambda_T^2\frac{M^2_{\tilde B}}{M^2_{\tilde W}-M^2_{\tilde B}}\right)\log\frac{M^2_{\tilde B}}{Q^2}\Bigg]h_u^4\,, \end{aligned} \tag{2.13}$$

where Q is the renormalization scale and the first two lines show the scalar contribution while the third one shows the fermionic one. A particularly simple expression can be obtained in the limit $m^2_{R_d} \simeq m_T^2 \simeq m_S^2 = m^2_{\rm adj}$ and $M_{\tilde W} \simeq M_{\tilde B} = M_D$. The Higgs quartic is then

$$V^{CW}_{\rm Higgs} \supset \frac{1}{4}\left[\frac{5\lambda_T^4+2\lambda_T^2\lambda_S^2+\lambda_S^4}{16\pi^2}\log\frac{m^2_{\rm adj}}{M_D^2}+\frac{\lambda_S^2\lambda_T^2}{16\pi^2}\right]h_u^4\,, \tag{2.14}$$

so that a relevant positive contribution to the quartic can be obtained for a large enough ratio $m_{\rm adj}/M_D$.

Region 2. In the opposite limit, $m_{\rm adj} \ll M_D$, the one-loop contribution to the Higgs quartic is:

$$V^{CW}_{\rm Higgs} \supset \frac{1}{4}\left[\frac{5\lambda_T^4}{32\pi^2}\log\frac{M^2_{\tilde{W}}}{Q^2} + \frac{\lambda_S^4 + 2\lambda_T^2\lambda_S^2}{32\pi^2}\log\frac{M^2_{\tilde{B}}}{Q^2}\right] h_u^4$$
$$- \frac{1}{4}\left[\frac{5\lambda_T^4}{16\pi^2}\log\frac{M^2_{\tilde{W}}}{Q^2} + \frac{\lambda_S^4}{16\pi^2}\log\frac{M^2_{\tilde{B}}}{Q^2} - \frac{\lambda_T^2\lambda_S^2}{8\pi^2}\right] h_u^4\,, \qquad (2.15)$$

where we have also assumed $m^2_{R_d} \ll M_D^2$. The first line shows the scalar contribution, the second line the fermionic one.

Putting all together, we end up with

$$V^{CW}_{\rm Higgs} \supset \frac{1}{4}\left[-\frac{5\lambda_T^4}{32\pi^2}\log\frac{M^2_{\tilde{W}}}{Q^2} - \frac{\lambda_S^4 - 2\lambda_T^2\lambda_S^2}{32\pi^2}\log\frac{M^2_{\tilde{B}}}{Q^2} + \frac{\lambda_T^2\lambda_S^2}{8\pi^2}\right] h_u^4\,, \qquad (2.16)$$

so that we expect this contribution to be always negative. Let us notice that this region corresponds to the pure supersoft spectrum, where indeed the non-holomorphic scalar masses are negligible and $M_D^2 \gtrsim B$ in order to avoid problems with tachyonic masses. Furthermore, $m^2_{R_d}$ is given by the gaugino induced one-loop correction [3]:

$$m^2_{R_d} = \frac{\alpha_2}{\pi} M^2_{\tilde{W}} \log\frac{4M^2_{\tilde{W}} - 2B_T}{M^2_{\tilde{W}}}\,, \qquad (2.17)$$

with an inert doublet therefore too light to give any significant boost to the Higgs mass.

For comparison, the well known stop contribution is given by [18]

$$V^{CW}_{\rm Higgs} \supset \frac{1}{4}\left[\frac{3}{16\pi^2}y_t^2\left(y_t^2 - \frac{m_Z^2}{2v^2}\right)\log\frac{M^2}{m_t^2}\right.$$
$$\left. + \frac{3y_t^4}{(16\pi^2)^2}\left(\frac{3}{2}y_t^2 - 32\pi\alpha_3(m_t)\right)\log^2\frac{M^2}{m_t^2}\right] h_u^4\,, \qquad (2.18)$$

where y_t is the top Yukawa coupling, α_3 the strong coupling constant and M is a common soft SUSY breaking stop mass scale. We show also the two-loop contribution, since the term proportional to the strong gauge coupling may reduce in a significant way the Higgs quartic, and we have set to zero the stop A terms, which are forbidden by the R-symmetry.

The simplified expressions, eqs. (2.13), (2.14), suggest that for $|\lambda_T| \simeq |\lambda_S| \simeq y_t$ the new states may give a contribution comparable to the stop one, depending on the mass hierarchy. In this sense, they can be regarded as "additional stops", which in principle can allow for a collective loop enhancement of the Higgs quartic. Moreover, since both the triplet and the singlet are uncolored, we do not expect the two loop terms proportional to the gauge couplings to give a reduction analogous to the one proportional to α_3 in the stop sector, making more effective the loop boost achieved through these states. Whether or not this scenario will allow to obtain a 125 GeV Higgs with less fine-tuning than in the MSSM will be studied in detail in section 4.

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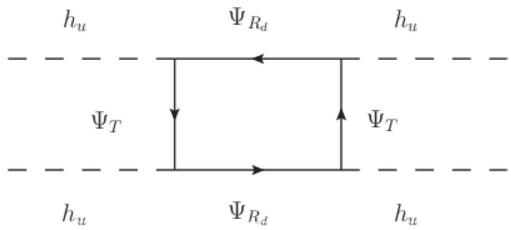

Figure 2. Box diagram that contributes to the T parameter.

As a last comment, let us notice that the condition $|\lambda_T| \simeq |\lambda_S| \simeq y_t$ (i.e. rather large values for the trilinear couplings at the weak scale) may imply a loss of perturbativity at relatively low scales. Solving the RGE's [19] requiring $\lambda_T = 1 = -\lambda_S$ at the weak scale, we find that the coupling that runs faster, λ_T, reaches $\sqrt{4\pi}$ for scales above 100 TeV.

3 Electroweak precision measurements

Getting a significant help from the triplet and the singlet to raise the Higgs mass through radiative corrections requires an appreciable value for the couplings λ_T and/or λ_S. However, the same couplings contribute to the T parameter at loop level. Therefore there exist a potential tension between generating a large Higgs mass and electroweak precision data. Besides the loop-level corrections to T there is already at tree level a dangerous effect due to the vev for the triplet T^0, which can lead to a large contribution to the $\hat{T}$ parameter (with the standard $T = \hat{T}/\alpha$):

$$\hat{T} = 4\frac{v_T^2}{v^2}, \tag{3.1}$$

which constrain the triplet vev to be $|v_T| \lesssim 3\,\text{GeV}$, where

$$v_T = \frac{\sqrt{2}gM_{\tilde{W}} - 2\lambda_S\lambda_T v_S - 2\lambda_T\mu}{4B_T + 8M_{\tilde{W}}^2 + 2m_T^2 + 2\lambda_T^2 v^2}\, v^2 \ . \tag{3.2}$$

It can be minimized by taking m_T large, or otherwise arranging for the numerator to be small. Besides the tree level contributions, there are contributions coming from loops of superpartners. A detailed study of all these contributions will be presented in [20], but the dominant effect comes from contributions to $\hat{T}$ from loops involving the fermionic part of the superfield H_u, T, R_d and S. Integrating them out at loop level, trough the diagram of figure 2 lead to the higher-dimension operator associated with $\hat{T}$:

$$\frac{\left|H_u^\dagger D_\mu H_u\right|^2}{\Lambda^2}, \tag{3.3}$$

with a coefficient proportional to λ_T^4. Thus, the same coupling which can help to make the Higgs heavier will also lead to sizeable contributions to T. To estimate the region of parameter space excluded by electroweak precision data we compute the $\hat{T}$ parameter due

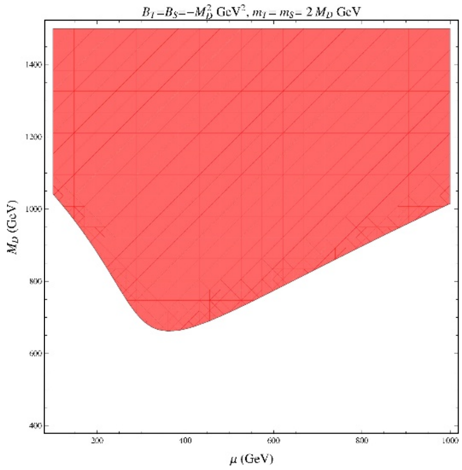

Figure 3. Shown in red, the region allowed at 95% C.L. by EWPM ($T < 0.2$) as a function of $M_D = M_{\tilde{W}} = M_{\tilde{B}}$ and μ. The adjoint holomorphic and non-holomorphic masses are fixed to $m_S = m_T = 2M_D$ and $B_S = B_T = -M_D^2$, respectively. The supersymmetric trilinear couplings are $\lambda_T = 1 = -\lambda_S$.

to v_T, and from loops of the scalar and fermionic sector of H_u, T, R_d and S. Imposing $T < 0.2$ [21], we find that the fermion loops will force $\lambda_T \lesssim 1$, whenever $M_D \lesssim 1\,\text{TeV}$. We notice however that, as can be anticipated from eq. (3.2), the μ parameter plays an essential role in keeping the contributions to $\hat{T}$ under control, at least at tree level. This is confirmed in figure 3, from which it can be clearly seen that the expected lower bound on M_D is dramatically modified in the $\mu = (300 - 400)\,\text{GeV}$ region, where gaugino masses as low as $M_D \simeq 650\,\text{GeV}$ are allowed. This can be essentially traced back to the smallness of the numerator of eq. (3.2), that makes the tree level contribution to $\hat{T}$ basically negligible.

4 125 GeV Higgs boson and fine-tuning

We are now in a position to analyze the region of parameter space in which not only a 125 GeV Higgs boson mass can be obtained, but also the contributions to the T parameter can be kept under control. Before doing so we comment on the fine-tuning in our set up.

Following [22], we consider two possible sources of tuning. The first one measures the sensitivity of the vev on the fundamental parameters $a_i = \{m_T^2, m_S^2, m_{R_d}^2, m_{Q_3}^2, m_{u_3}^2, \mu, M_{\tilde{W}}, M_{\tilde{B}}, B_T, B_S, \lambda_T, \lambda_S\}$,

$$\Delta = \max_{a_i} \left| \frac{\partial \log m_Z^2}{\partial \log a_i} \right| . \tag{4.1}$$

The second one measures the sensitivity of the physical Higgs mass m_h^2 on the same set of parameters (this time for fixed vev's):

$$\Delta_h = \max_{a_i} \left| \frac{\partial \log m_h^2}{\partial \log a_i} \right| . \tag{4.2}$$

Let us notice that eq. (4.2) effectively measures the tuning on the Higgs quartic coupling.

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As customary, the dependence of m_Z^2, eq. (A.1), on the high energy parameters (defined at the scale Λ at which they are generated) is taken into account solving the RGE's for $m_{H_u}^2$ [19]. In the leading-log approximation,

$$\delta m_{H_u}^2 \simeq \frac{1}{8\pi^2}\Big[\left(\lambda_S^2+3\lambda_T^2\right)m_{R_d}^2+\lambda_S^2 m_S^2+3\lambda_T^2 m_T^2+3y_t^2\left(m_{Q_3}^2+m_{u_3}^2\right)\Big]\log\frac{\Lambda}{\mathrm{TeV}}\,. \tag{4.3}$$

In contrast to what happens in models with Majorana gaugino masses, no dependence on the Dirac gaugino masses is present in the RGE's for $m_{H_u}^2$.[6]

Regarding the computation of the fine-tuning, the usual estimates are not valid in this case. Indeed, once v_T and v_S are inserted in the expression for m_Z^2, eq. (A.1), the dependence on the parameters is different from the MSSM one, which is recovered only in the $m_{\mathrm{adj}} \to \infty$ limit. Although the expressions are quite involved, we can obtain a good approximation solving perturbatively the minimum equations in powers of $\alpha_i = v^2/(2B_i + 4M_i^2 + m_i^2 + \lambda_i^2 v^2)$, which are small quantities in the region of parameter space we are going to consider (i.e. $M_D, m_{\mathrm{adj}} \gtrsim 500\,\mathrm{GeV}$). For instance, to order $\alpha_{s,t}$, the triplet and singlet vev's read

$$v_T = \alpha_T\left(\frac{gM_{\tilde{W}}}{\sqrt{2}} - \lambda_T\mu\right), \qquad v_S = -\alpha_S\left(\frac{g'M_{\tilde{B}}}{\sqrt{2}} + \lambda_S\mu\right), \tag{4.4}$$

from which it is clear that both quantities are small. As discussed in section 2.1, this implies that for $\lambda_T > 0$ and $\lambda_S < 0$, the tree level mixing between h_u and s, t can be small, and the tree-level Higgs mass may be close to the MSSM upper bound. Computing the variation of the first of eqs. (A.1) we obtain

$$\begin{aligned}
\Delta_\mu &= \left|\frac{4\mu^2}{m_Z^2} + \frac{8\mu}{m_Z^2}\left[\lambda_S v_S + \lambda_T v_T + \frac{2\mu}{m_Z^2}\left(\frac{v_S^2}{\alpha_S} + \frac{v_T^2}{\alpha_T}\right)\right]\right|\,,\\
\Delta_{M_{\tilde{W}}} &= \left|\frac{4\sqrt{2}gM_{\tilde{W}}v_T}{m_Z^2}\right|\,,\\
\Delta_{M_{\tilde{B}}} &= \left|\frac{4\sqrt{2}g'M_{\tilde{B}}v_S}{m_Z^2}\right|\,,\\
\Delta_{B_T} &= \left|\frac{8B_T}{m_Z^2}\frac{v_T^2}{v^2}\right|\,,\\
\Delta_{B_S} &= \left|\frac{8B_S}{m_Z^2}\frac{v_S^2}{v^2}\right|\,,\\
\Delta_{m_{R_d}^2} &= \left|\frac{\lambda_S^2+3\lambda_T^2}{4\pi^2}\frac{m_{R_d}^2\log\frac{\Lambda}{\mathrm{TeV}}}{4m_Z^2}\left[1+\frac{4}{m_Z^2}\left(\frac{v_S^2}{\alpha_S}+\frac{v_T^2}{\alpha_T}\right)\right]\right|\,,
\end{aligned} \tag{4.5}$$

[6] There is nonetheless a dependence through the finite one-loop contribution analogous to eq. (2.17); we checked however that the tuning due to this contribution is never the dominant one in the interesting regions.

$$\Delta_{m_T^2} = \left| \frac{3\lambda_T^2 m_T^2 \log \frac{\Lambda}{\text{TeV}}}{4\pi^2 m_Z^2} \left[1 + \frac{4}{m_Z^2}\left(\frac{v_S^2}{\alpha_S} + \frac{v_T^2}{\alpha_T}\right)\right]\right| ,$$

$$\Delta_{m_S^2} = \left| \frac{\lambda_S^2 m_S^2 \log \frac{\Lambda}{\text{TeV}}}{4\pi^2 m_Z^2} \left[1 + \frac{4}{m_Z^2}\left(\frac{v_S^2}{\alpha_S} + \frac{v_T^2}{\alpha_T}\right)\right]\right| ,$$

$$\Delta_{m_{Q_3}^2, m_{u_3}^2} = \left| \frac{3y_t^2 m_{Q_3,u_3}^2 \log \frac{\Lambda}{\text{TeV}}}{4\pi^2 m_Z^2} \left[1 + \frac{4}{m_Z^2}\left(\frac{v_S^2}{\alpha_S} + \frac{v_T^2}{\alpha_T}\right)\right]\right| .$$

We immediately see that in the limit $v_S = v_T = 0$ we get the same expressions we would have obtained from the MSSM's minimum condition, i.e. only Δ_μ, and $\Delta_{m_{Q_3}^2, m_{u_3}^2}$ do not vanish. Switching on the singlet and triplet contributions, we get corrections which start at $\mathcal{O}(\alpha_{T,S})$ (notice that, according to eq. (4.4), v_S and v_T are $\mathcal{O}(\alpha_{T,S})$, and so are v_S^2/α_S and v_T^2/α_T). In particular, since $\Delta_{M_{\tilde{W},\tilde{B}}} \sim \mathcal{O}(\alpha_{T,S})$ and $\Delta_{B_T,B_S} \sim \mathcal{O}(\alpha_{T,S}^2)$, the related tuning is never relevant.

Among the tunings due to the soft SUSY breaking masses, for $\lambda_T \sim \lambda_S \sim 1$ we get comparable results from $m_{R_d}^2$, m_T^2 and m_{Q_3,u_3}^2 (with a slightly worse sensitivity associated to the inert doublet mass, $m_{R_d}^2$). Indeed, taking all the soft masses to be of the same order, we get

$$\Delta_{m_{R_d}^2} \simeq 4\Delta_{m_S^2} \simeq \frac{4}{3}\Delta_{m_T^2} \simeq \frac{4}{3}\Delta_{m_{Q_3}^2} . \tag{4.6}$$

In particular, there is no worsening in the fine-tuning for $m_{\tilde{t}} \simeq m_{\text{adj}}$.

As usual, we need also to keep under control the tree level tuning due to μ, whose contribution may rapidly become the dominant one (especially in the region with relatively small m_{adj} in which the remaining sensitivities are not particularly severe). In particular, in the region $M_D, m_{\text{adj}} \lesssim 2\,\text{TeV}$, we have

$$\begin{aligned} \mu = 100\,\text{GeV} &\rightarrow \Delta_\mu \sim 5 - 10\,, \\ \mu = 200\,\text{GeV} &\rightarrow \Delta_\mu \sim 20 - 30\,, \\ \mu = 300\,\text{GeV} &\rightarrow \Delta_\mu \sim 45 - 55\,. \end{aligned} \tag{4.7}$$

Since μ gives the Higgsino mass, we need to be careful with the limits imposed by direct searches and EWPM. In principle, we would expect smaller values of μ to be preferred since they give smaller Δ_μ. However, as shown in figure 3, light Higgsinos require heavier gauginos to be compatible with EWPM (unless we take $\mu = (300 - 400)\,\text{GeV}$). As we are going to see, this in turn implies heavier scalars to accommodate $m_h = 125\,\text{GeV}$, with a general worsening of the fine-tuning.

Let us now discuss the sensitivity of the Higgs quartic couplings on the parameters, eq. (4.2). Integrating out the heavy fields in eq. (2.8) we obtain

$$V \supset \frac{1}{4v^2}\left[\frac{m_Z^2}{4} - \left(\frac{v_T^2}{\alpha_T} + \frac{v_S^2}{\alpha_S}\right) + \lambda_{\text{loop}}\right] h_u^4 , \tag{4.8}$$

from which we easily compute

$$
\begin{aligned}
\Delta_h^{(\mu)} &= \frac{8\mu}{m_h^2}\left(\lambda_T v_T + \lambda_S v_S\right) + \Delta_h^{(\mu)}\Big|_{\text{loop}} \,, \\
\Delta_h^{(M_{\tilde W})} &= \frac{4M_{\tilde W} v_T}{m_h^2}\left(\frac{8M_{\tilde W} v_T}{v^2} - \sqrt{2} g\right) + \Delta_h^{(M_{\tilde W})}\Big|_{\text{loop}} \,, \\
\Delta_h^{(M_{\tilde B})} &= \frac{4M_{\tilde B} v_S}{m_h^2}\left(\frac{8M_{\tilde B} v_S}{v^2} - \sqrt{2} g'\right) + \Delta_h^{(M_{\tilde B})}\Big|_{\text{loop}} \,, \\
\Delta_h^{(m_T^2)} &= \frac{4m_T^2 v_T^2}{m_h^2 v^2} + \Delta_h^{(m_T^2)}\Big|_{\text{loop}} \,, \\
\Delta_h^{(m_S^2)} &= \frac{4m_S^2 v_S^2}{m_h^2 v^2} + \Delta_h^{(m_S^2)}\Big|_{\text{loop}} \,.
\end{aligned}
\tag{4.9}
$$

The contributions dubbed $\Delta_h^{(i)}\Big|_{\text{loop}}$ are those coming from λ_{loop} in eq. (4.8). From eq. (4.9) it is clear that all the tree level contributions start either at $\mathcal{O}(\alpha_{T,S})$ or at $\mathcal{O}(\alpha_{T,S}^2)$ and are thus going to be irrelevant. It can nevertheless happen that some or all of the $\Delta_h^{(i)}\Big|_{\text{loop}}$ are large. We can get an idea of the typical tuning arising at loop level using the approximation of eqs. (2.14) together with the stop contribution, eq. (2.18): we always have $\Delta|_{\text{loop}} \lesssim 1$. We checked numerically that this is also the case when the complete loop contributions are taken into account, so that in the following we will discard Δ_h.

We now focus on the more precise numerical analysis. For this we turn again to the Coleman-Weinberg potential, eq. (2.12). The first step is to find the physically interesting minimum (after adding the tree-level potential). At tree-level, it is simple to use the minimization conditions to determine $m_{H_u}^2$ so that $\langle h_u \rangle \approx v = 246$ GeV is reproduced. This depends on the triplet and singlet vev's, which could themselves be exchanged for m_T^2 and m_S^2, respectively. The corresponding expressions are given in eq. (A.1) of the appendix. In order to minimize the one-loop corrected potential, we use the previous tree-level relations in the Coleman-Weinberg expression, eq. (2.12), and find again the tree-level $m_{H_u}^2$ by differentiating w.r.t. the vev's, but keeping the vev dependence given by eq. (A.1) fixed. This procedure is appropriate at 1-loop order. As remarked earlier, we need to ensure that both v_T and v_S be small,[7] so that the potential is essentially a function of $\langle h_u \rangle$, with $m_{H_u}^2$ fixed so that $\langle h_u \rangle$ gets the required value to reproduce m_Z. Indeed, we have checked that setting $v_T = v_S = 0$ from the start and focusing on the h_u dependence is a good approximation throughout the region of parameter space presented in our plots. Furthermore, we have also checked that expanding the potential and truncating

[7]We recall that for the singlet it is possible to write a tadpole term. Even though, for simplicity, we are setting this tadpole to zero at tree level, it will be generated at 1-loop. In order to keep v_S small, we then need to increase m_S^2. Once this is done, the effect on the minimization, which enters only through v_S is a small effect. The main effect is on the real part of the singlet fluctuations, whose mass is directly controlled by m_S^2. Since, in practice, this mass is relatively large (in the multi-TeV range), the mixing with the lighter states is also small.

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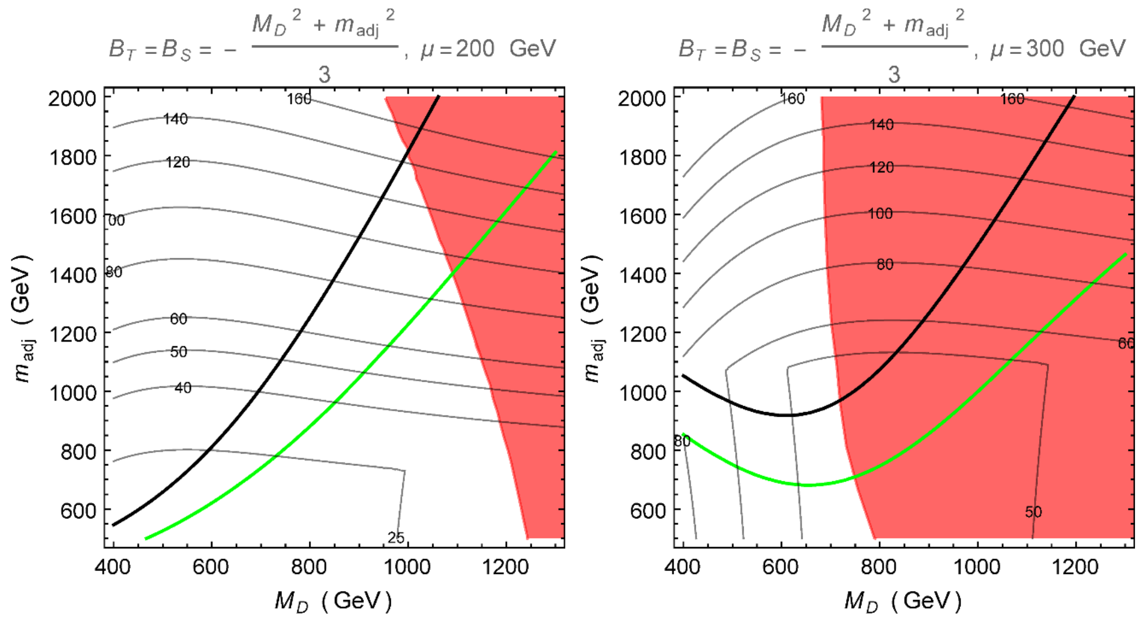

Figure 4. Higgs boson mass $m_h = 125\,\text{GeV}$ (black and green thick lines) and fine-tuning parameter Δ (thin lines), as a function of $M_D = M_{\tilde{W}} = M_{\tilde{B}}$ and $m_{\text{adj}} = m_T = m_S = m_{R_d}$, for $B_T = B_S = -\frac{1}{3}(m_{\text{adj}}^2 + M_D^2)$. We fix $\lambda_T = 1 = -\lambda_S$. The upper (black) curve refers to a common stop mass of $m_{stop} = 300\,\text{GeV}$, the lower (green) curve to $m_{stop} = m_{\text{adj}}$. Left panel: $\mu = 200\,\text{GeV}$; right panel: $\mu = 300\,\text{GeV}$. The red region is allowed at 95% C.L. by EWPM ($T < 0.2$.)

at order h_u^4 is an excellent approximation. Since making these approximations allows for a significantly faster evaluation without losing too much precision, we show the results obtained in this limit. Hence, in our plots the tree-level masses that enter the Coleman-Weinberg potential have been evaluated perturbatively. We have, however, checked by a full numerical evaluation of the potential (on a coarser grid of points) that the errors incurred by this procedure are small (more comments later on). In summary, compared to the analytical expressions given in section (2.2), other than the small field approximation, we keep the full dependence on all other parameters, in particular not assuming a large hierarchy between M_D, m_{adj} and μ. The Higgs boson mass is obtained by evaluating the second derivatives of the potential at its minimum, and diagonalizing. We have also checked that the difference in the Higgs mass between the approximate and fully numerical evaluations is at most a couple of GeV, which is comparable to other uncertainties that have not been included here, such as higher-loop orders. Even though the Higgs mass is currently known to a significantly better precision, we believe that a more sophisticated and precise analysis would not alter in an important way our conclusions.

Our main results are shown in figure 4. We present them as a function of $M_D = M_{\tilde{W}} = M_{\tilde{B}}$ and $m_{\text{adj}} = m_T = m_S = m_{R_d}$, with couplings fixed to $\lambda_T = 1 = -\lambda_S$. On the left panel $\mu = 200\,\text{GeV}$, while on the right panel $\mu = 300\,\text{GeV}$. The solid thick lines correspond to $m_h = 125\,\text{GeV}$, with the red region allowed at 95% C.L. by EWPM. We also show the largest among the fine-tuning parameters Δ_i (thin black lines), eq. (4.5), fixing $\Lambda = 20\,\text{TeV}$.[8] The $m_h = 125\,\text{GeV}$ thick lines refer to two different stop masses: $m_{\tilde{t}} = 300\,\text{GeV}$[9] for the upper curve and $m_{\tilde{t}} = m_{\text{adj}}$ for the lower curve.

Let us comment on two counterintuitive features of our results: the correct Higgs mass is achieved with less fine-tuning for *heavier* stops and for *heavier* Higgsinos. This can be understood as follows: for the upper (black) curves, the lightness of the stops is such that the main boost to the Higgs quartic comes from the adjoint and inert fields. On the contrary, for the lower (green) curves the stop boost to the Higgs quartic is relevant. However, as already pointed out, there is no worsening in the tuning for $m^2_{\tilde{s}top} = m^2_T = m^2_{R_d}$, eq. (4.6). Moreover, the "collective" quartic enhancement in the lower curves allows for smaller soft SUSY breaking masses, implying thus less tuning. Turning to the μ parameter, we already observed that compatibility with EWPM for lighter Higgsinos require heavier gauginos (i.e. larger M_D). In addition, it is clear from the shape of the Higgs mass curves in figure 4 that this in turn requires heavier scalars to get $m_h = 125\,\text{GeV}$, so that a worsening in the tuning is expected. This is indeed the case in figure 4: for $\mu = 300\,\text{GeV}$ compatibility between $m_h = 125\,\text{GeV}$ and EWPM is achieved for $m_{\text{adj}} \gtrsim 800 - 1100\,\text{GeV}$ (for $m_{\tilde{t}} = m_{\text{adj}}$ or $300\,\text{GeV}$, respectively), i.e. when the sensitivity is still dominated by μ. On the contrary, for $\mu = 200\,\text{GeV}$ the scalar masses are pushed up to $m_{\text{adj}} \gtrsim 1500 - 1900\,\text{GeV}$ (again for $m_{\tilde{t}} = m_{\text{adj}}$ or $m_{\tilde{t}} = 300\,\text{GeV}$, respectively), in a region in which the soft SUSY breaking masses dominate the tuning.

For comparison, in the MSSM with maximal stop mixing $m_h = 125\,\text{GeV}$ is achieved with $\Delta \gtrsim 100 - 200$ [23], while since for $A_t = 0$ stop masses around 10 TeV are needed [24], we can estimate $\Delta \gtrsim 2000$. We do not attempt here to follow [22] and find the minimum tuning achievable in the model; it is however clear that among the beneficial effects of our R-symmetric scenario we have a significant fine-tuning reduction (modulo potential additional sources from the UV that might be reduced with further model building).

We also show in table 1 the spectrum of the scalar and neutralino/chargino sectors for a representative benchmark point. The numbers have been obtained by evaluating numerically the 1-loop effective potential, without the approximations discussed earlier. In particular, we evaluate the spectrum of the real fluctuations by taking the second derivative of the 1-loop effective potential numerically. However, for the imaginary parts we diagonalize numerically the tree-level mass matrices.[10] For this particular benchmark, we find that the approximate evaluation would produce a Higgs about 2.5 GeV lighter, while the difference in the other masses is never larger than $6 - 7\%$. The benchmark example illustrates a typical spectrum, with most scalar states at around $1 - 2$ TeV, and with

[8]In this work we assume that the required soft parameters can be obtained naturally with the appropriate values at this scale. However, as mentioned previously, this is a somewhat non-trivial task and could be the source of additional fine-tuning [15].

[9]A detailed study of the LHC phenomenology of the model is outside the scope of the present work, therefore we assume $m_{\tilde{t}} \sim 300\,\text{GeV}$ to be still allowed by the LHC either because of a very compressed spectrum or because of baryonic R-Parity violating couplings in the superpotential.

[10]This is because we have not allowed for imaginary vev's in the Coleman-Weinberg potential and therefore cannot evaluate the second derivatives in those directions. We expect the 1-loop corrections in this sector to be small so that a tree level treatment will be sufficient to illustrate our main points. Also, since we are assuming no CP violation, the real and imaginary sectors are decoupled.

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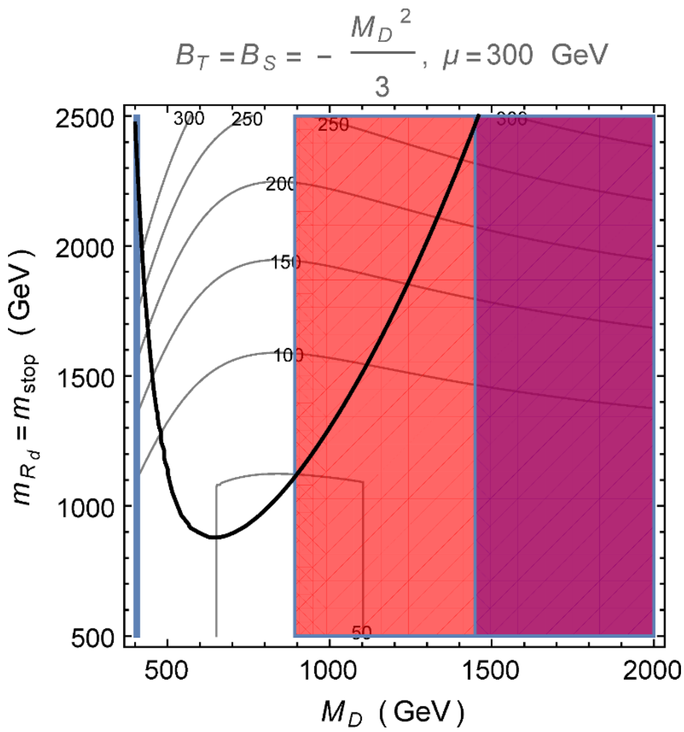

Figure 5. Higgs boson mass $m_h = 125\,\text{GeV}$ (solid line) and fine-tuning parameter (thin lines), as a function of $M_D = M_{\tilde{W}} = M_{\tilde{B}}$ and $m_{\tilde{t}} = m_{R_d}$, with $m_T = m_S = 0$ (supersoft limit), and $B_T = B_S = -\frac{M_D^2}{3}$. We fix $\lambda_T = 1 = -\lambda_S$ and $\mu = 300\,\text{GeV}$. Blue region: spontaneous charge and/or CP breaking. Red region: allowed at 95% C.L. by EWPM ($T < 0.2$.) . Purple region: $m_{\tilde{\ell}_R} > 100\,\text{GeV}$.

Input Parameters	μ	$M_{\tilde{B}}$	$M_{\tilde{W}}$	m_S^2	B_S	m_T^2	B_T	λ_S	λ_T
	300	1040	1040	$(1000)^2$	$-(833)^2$	$(1000)^2$	$-(833)^2$	-1	1

Scalar Spectrum	h	R_d	Re(S)	Im(S)	Re(T)	Im(T)
	126	1070	3660	1550	1860	1570

Fermion Spectrum	χ_1^0	χ_2^0	χ_3^0	$\chi_1^\pm$	$\chi_2^\pm$	$\chi_3^\pm$
	313	1040	1070	316	1040	1070

Table 1. We exhibit a benchmark point that illustrates a typical spectrum within our model. We show the input parameters in the upper table. We work in the large $\tan\beta$ limit (≈ 60), and we have fixed $m_{H_u}^2$ by requiring that $v = 246$ GeV. The triplet and singlet vev's are both around 1 GeV. We have fixed, in addition, the common stop masses in eq. (2.18) as $M = 1000$ GeV, and the renormalzation scale at $Q = 600$ GeV. In the lower tables, we show the physical scalar and fermion spectrum (showing only those states that play a role in EWSB). All masses are in GeV.

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the lightest neutralino/chargino at a few hundred GeV. For this point, we find that the T-parameter is $T \approx 0.16$, comfortably within the experimental limits, even though both λ_T and λ_S are order one. As already remarked earlier, the details in the slepton/squark sector are largely decoupled [with the exception of the stop mass, which we have taken here at a scale $M = 1000$ GeV, see eq. (2.18)]; hence, they are not shown in the table. Since we do not include two-loop effects, the gluino and octet scalar masses do not enter. However, as we remarked earlier these can be taken around $4 - 5$ TeV, without incurring in excessive tuning from this sector. The tuning, as measured by the sensitivity to the lagrangian parameters of this point is about 2%, and is dominated by μ.

Regarding the phenomenology at the LHC run II, the discovery potential for a 1 TeV stop strongly depends on its decay modes. For instance, such stops are within the LHC reach assuming either R-parity conservation or Leptonic R-parity violation [25, 26] (we assume the sensitivity in the LRPV case to be roughly the same as in the R-parity conserving case, as it is for current searches), while in general we can expect that they will be more difficult to discover assuming Baryonic R-parity violation [27, 28]. In addition, in generic models with Dirac gauginos degenerate squarks with 1 TeV masses are still an open possibility [4], independently of the details of the R-parity violating sector. Sleptons may be within the LHC reach, but as already pointed out, our conclusions are largely independent of the details of this sector. The other CP-even states in the Higgs sector are generically too heavy to be discovered, both in direct LHC searches and through their effect on Higgs coupling measurements. On the contrary, 300 GeV Higgsinos may be in the LHC13 reach, for sufficient luminosity [25, 26].

As a final comment, we consider the case of a pure supersoft spectrum, i.e. the case in which we set the non-holomorphic adjoint masses m_T^2 and m_S^2 to zero. With $m_{R_d}^2$ given by (2.17), we easily see from eq. (2.16) that the only relevant radiative correction comes from the stop sector, with the gaugino contributions decreasing the Higgs mass. However, it is easy to imagine that the full UV completion of the model may contain a sector which couples the Higgs multiplets to messenger fields. This can generate the required μ and B_μ terms, as well as extra contributions to the R_d mass which may then push up the Higgs mass even with moderately light stops. In any case, relying only on the loop corrections from R_d makes the boost to the Higgs quartic less efficient than in the non-supersoft case, making the model less natural. There is also another important constraint to take into account in a pure supersoft spectrum: all the sfermions acquire mass via finite gaugino one loop contribution and are therefore predicted (modulo the running from the adjoint mass scale down to the weak scake) in term of M_D. This tends to make the sleptons, especially the right-handed ones which only have hypercharge gauge couplings, quite light. This is illustrated in figure 5, where we show once again, in the $M_D - m_{\text{stop}}$ plane, the line corresponding to a 125 GeV Higgs together with the region allowed by EWPM (in red). We also show the region where the slepton has a mass greater than 100 GeV (purple region), which correspond to the LEP bound. We see that the slepton constraint pushes all the masses to be very heavy, into a region with very large fine-tuning. This leads to the conclusion that sizable non-holomorphic adjoints masses are required to reduce the fine-tuning.

5 Conclusions

We are finally in an era in which experiments are directly exploring the electroweak scale, and will (at least in part) shed light on whether or not the electroweak scale is natural. The first LHC run has already provided us with some indications. The general message seems to be that our simplest natural models are by now tuned at the percent level, or worse. While it can turn out that this is the level of tuning of the EW scale, it may also be taken to motivate the search for more natural (although less minimal) models. One such possibility is the supersymmetric model with a quasi exact $\mathrm{U}(1)_R$ symmetry considered in this work. This kind of models are more natural with respect to the MSSM, since there is no gluino induced one-loop contribution to the squark masses. Moreover, the usual supersymmetric flavor problem is greatly ameliorated. The point on which we focus here is that the adjoint superfields needed to write Dirac gaugino masses may give relevant loop corrections to the Higgs boson mass: they act effectively as "additional stops", at least in part of the parameter space. A possible drawback is that the very same couplings that help increasing the Higgs boson mass break custodial symmetry, potentially leading to large contributions to the electroweak precision measurements. Our main results are summarized in figure 4, in which we show the region in parameter space in which a 125 GeV Higgs mass can be obtained in a way compatible with EWPM. We also presented on the same plot the required fine-tuning. A first conclusion that can be drawn is that there are regions in which the fine-tuning is ameliorated with respect to the MSSM, roughly reduced to twice the tuning of the NMSSM, $\Delta \sim 20 - 30$ [22]. Let us stress however that the mechanism that allows for the increased naturalness is completely different: while in the NMSSM it is due to the enhanced tree level Higgs boson mass, here it is due to the collective loop enhancement which reduces the sensitivity to the single mass involved. Moreover, we do not attempt to scan the parameter space to find the minimum achievable tuning of the model compatible with experimental data. It may well be that in some region of parameter space the tuning can be better than the one here presented. A further point which is worth mentioning is that stop masses in the TeV range do not increase the fine-tuning, which is basically driven by $m^2_{R_d}$, eq. (4.5). Together with the already mentioned improved naturalness bound on the gluino mass, this makes the non observation so far of any superpartner at the LHC less worrisome. Note also that squarks could be of the same order as the stop and have escaped detection due to a reduced production cross section which weakens the LHC bounds [4]. Depending on the details of the model, such squarks could be discovered soon.

What can we expect to observe at LHC-13, given this framework? It is of course quite difficult to make a robust a solid statement. As we have seen, since the fine-tuning is driven by m_{R_d} and m_{adj} in the interesting part of the parameter space, naturalness does not require the stop to be as light as possible. On the contrary, a relatively heavy stop (with a mass around 1 TeV) is preferred since it can give a sizable contribution to the Higgs mass, allowing for the state which are driving the fine-tuning to be lighter. In any case, we still expect μ to be as low as possible, with the Higgsino possibly "right around the corner".

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A Potential minimization and mass matrices

We collect here useful formulas obtained from the minimization of the tree level scalar potential. For simplicity, we will take from the beginning the limit $\tan\beta \gg 1$.

Using for the vacuum expectation values the convention $\langle h_u\rangle = v$, $\langle t\rangle = v_T$, $\langle s\rangle = v_S$, the minimization of the scalar potential, eq. (2.8) gives

$$\begin{aligned} m_{H_u}^2 &= \sqrt{2}\left(gM_{\tilde{W}}v_T - g'M_{\tilde{B}}v_S\right) - \frac{m_Z^2}{2} - (\lambda_S v_S + \lambda_T v_T + \mu)^2 \,, \\ v_T &= \frac{\sqrt{2}gM_{\tilde{W}} - 2\lambda_S\lambda_T v_S - 2\lambda_T\mu}{2\left(2B_T + 4M_{\tilde{W}}^2 + m_T^2 + \lambda_T^2 v^2\right)}v^2 \,, \\ v_S &= -\frac{\sqrt{2}g'M_{\tilde{B}} + 2\lambda_S\lambda_T v_T + 2\lambda_S\mu}{2\left(2B_S + 4M_{\tilde{B}}^2 + m_S^2 + \lambda_S^2 v^2\right)}v^2 \,. \end{aligned} \tag{A.1}$$

The squared mass matrix for the CP-even scalars, in the (h_u, t, s, r_d) basis, reads

$$M^2_{\mathrm{CP-even}} = \begin{pmatrix} m_Z^2 & \cdot & \cdot & \cdot \\ v\left(2\lambda_T\left(\lambda_S v_S + \lambda_T v_T + \mu\right) - \sqrt{2}gM_{\tilde{W}}\right) & m_{T_R}^2 + \lambda_T^2 v^2 & \cdot & \cdot \\ v\left(2\lambda_S\left(\lambda_S v_S + \lambda_T v_T + \mu\right) + \sqrt{2}g'M_{\tilde{B}} +\right) & \lambda_S\lambda_T v^2 & m_{S_R}^2 + \lambda_S^2 v^2 & \cdot \\ 0 & 0 & 0 & m_H^2 \end{pmatrix}, \tag{A.2}$$

where $m_{T_R}^2$ and $m_{S_R}^2$ are defined below eq. (2.11), while m_H^2, the mass of the CP-even inert doublet, is given by

$$\begin{aligned} m_H^2 = \mu^2 + m_{R_d}^2 - \frac{m_Z^2}{2} + \sqrt{2}\left(gM_{\tilde{W}}v_T - g'M_{\tilde{B}}v_S\right) + 2\left(\lambda_S v_S + \lambda_T v_T\right)\mu + \\ + \left(\lambda_S v_S + \lambda_T v_T\right)^2 + \left(\lambda_S^2 + \lambda_T^2\right)v^2 \end{aligned} \tag{A.3}$$

Turning to the CP-odd squared mass matrix, in the (a_t, a_s, a_{r_d}) basis it is

$$M^2_{\mathrm{CP-odd}} = \begin{pmatrix} m_T^2 - 2B_T + \lambda_T^2 v^2 & \cdot & \cdot \\ \lambda_S\lambda_T v^2 & m_S^2 - 2B_S + \lambda_S^2 v^2 & \cdot \\ 0 & 0 & m_H^2 \end{pmatrix}, \tag{A.4}$$

with the CP-odd component of the inert doublet degenerate in mass with the CP-even part.

To conclude, the entries of the charged scalar squared mass matrix in the basis $(H_u^+, T^+, (T^-)^*, (R_d^-)^*)$ are

$$
\begin{aligned}
M_{11}^2 &= 2v_T\left[\sqrt{2}gM_{\tilde{W}} - 2\lambda_T\left(\lambda_S v_S + \mu\right)\right] \\
M_{12}^2 &= -\frac{v}{2}\left[\sqrt{2}g^2 v_T - 2gM_{\tilde{W}} + 2\sqrt{2}\lambda_T\left(\lambda_S v_S - \lambda_T v_T + \sqrt{2}\mu\right)\right] \\
M_{13}^2 &= \frac{v}{2}\left[\sqrt{2}g^2 v_T + 2gM_{\tilde{W}} - 2\sqrt{2}\lambda_T\left(\lambda_S v_S + \lambda_T v_T + \sqrt{2}\mu\right)\right] \\
M_{22}^2 &= m_T^2 + 2M_{\tilde{W}}^2 + 2\lambda_T^2 v^2 + \frac{g^2}{2}\left(2v_T^2 - v^2\right) \\
M_{23}^2 &= 2\left(M_{\tilde{W}}^2 + B_T\right) - g^2 v^2 \\
M_{33}^2 &= m_T^2 + 2M_{\tilde{W}}^2 + \frac{g^2}{2}\left(2v_T^2 + v^2\right) \\
M_{44}^2 &= m_H^2 + m_W^2 - 2\sqrt{2}gM_{\tilde{W}}v_T - 4\lambda_S\lambda_T v_S v_T - 4\sqrt{2}\lambda_T\mu v_T + \left(\lambda_T^2 - \lambda_S^2\right)v^2
\end{aligned}
\tag{A.5}
$$

with all the other entries vanishing. The 3×3 submatrix obtained by taking out the R_d^- entry has vanishing determinant as expected, since one combination of the charged scalars is the would-be Goldstone boson eaten up by the $W^\pm$.

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