

SUSY WITH R -SYMMETRY: CONFRONTING EW PRECISION OBSERVABLES AND LHC CONSTRAINTS*

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After motivation and short presentation of the minimal supersymmetric model with R -symmetry (MRSSM), we address the question of accommodating the measured Higgs boson mass in accordance with electroweak precision observables and LHC constraints.

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1. Introduction

This article is based on two recent papers by Diessner, Kalinowski, Kotlarski and Stöckinger [1, 2].

The discovery of a Higgs-boson candidate with a mass near 125 GeV by both the ATLAS and CMS collaborations at the LHC in 2012 [3] seemingly completes the Standard Model (SM). So far, the properties of the new state are within experimental errors consistent with the predictions of the SM. Nevertheless, the true nature of the discovered state has to be thoroughly explored. The Run 2 of the LHC (and later the high-lumi run) should clarify if the couplings of the new state are exactly as predicted by the SM, whether it unitarizes the WW scattering amplitude or not, and eventually discover new particles. Although the SM is able to describe a vast number of experimental measurements, there are many questions which cannot be address: *e.g.* Dark Matter, baryogenesis, *etc.* In particular, the question of stabilization of the Higgs boson mass with respect to the Planck scale has fueled theoretical speculations on beyond the SM physics. Among these, the TeV-scale supersymmetry is one of the most theoretically and experimentally studied options. So far, no direct signal of supersymmetry has been observed by the LHC experiments, and only limits on superpartner masses have been

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derived. However, those limits should be taken with a grain of salt as the experimental analyses were performed for simplified models with many additional assumptions. The current limits may be not valid in more general supersymmetric scenarios and dedicated phenomenological studies are required for non-minimal models. In fact, the absence of any direct signal of supersymmetric particle production at the LHC, and the measured Higgs boson mass ~ 125 GeV close to the upper value of ~ 135 GeV achievable in the MSSM are a strong motivation to consider non-minimal SUSY scenarios. R -symmetric supersymmetric models, invariant under a global $U(1)_R$ transformation $\theta \rightarrow e^{i\alpha}\theta$, are particularly well motivated. R -invariance is indeed a symmetry [4] of all basic building blocks of the supersymmetric extension of the SM. This symmetry is stronger than R -parity because forbids not only baryon- and lepton-number changing terms in the superpotential, as well as dimension-five operators mediating proton decay, but also removes left–right sfermion mixing, the higgsino μ -term and Majorana gaugino masses. As a result, several of the most important experimental constraints on supersymmetry are alleviated: contributions to CP- and flavor-violating observables can be suppressed even in the presence of flavor violation in the sfermion sector, and the production cross section for squarks reduced, making squarks below the TeV scale generically compatible with LHC data.

Since this is the write-up of a lecture given at a school, the next section is devoted to a brief exposition of the MRSSM structure. Then, we address the problem of confronting the MRSSM with the observed Higgs boson mass in accordance with electroweak precision measurements. This is not an obvious task since, as to be seen in the next section, in the MRSSM, the lightest Higgs boson tree-level mass is typically reduced compared to the MSSM due to the mixing with additional scalars. Moreover, an important MSSM mechanism of generating large radiative corrections due to the stop mixing is absent, and R -symmetry necessarily introduces an $SU(2)$ scalar triplet, which can increase m_W already at the tree level. Nevertheless, in Refs. [1, 2], a number of benchmark points illustrating different viable parameter regions have been identified and verified that they are not excluded by further experimental constraints from Higgs observables, collider and low-energy physics. In this write-up, we will show the results for only one of the benchmarks BMP1 corresponding to $\tan\beta = 3$; for other BMPs, we refer to our original publications, where a comprehensive analysis of the parameter space is discussed and a complete list of references can also be found. We conclude with summary and outlook.

2. MRSSM

Under the global $U(1)_R$ R -symmetry, the Grassmann coordinate in the superspace $\{x^\mu, \theta, \bar{\theta}\}$ transforms as $\theta \rightarrow e^{i\alpha}\theta$. By convention, we can assign the R -charge 1 (-1) to the coordinate θ ($\bar{\theta}$). If a super field $\hat{X} = \phi_x + \theta\chi + \bar{\theta}\bar{\chi} + \dots$ has a well defined R -charge R_X , so that it transforms as

$$\hat{X}(x^\mu, \theta, \bar{\theta}) \rightarrow e^{i\alpha R_X} \hat{X}(x^\mu, e^{i\alpha}\theta, e^{-i\alpha}\bar{\theta}),$$

the component fields must transform differently. Obviously, the scalar component ϕ_x has R -charge R_X , the fermionic component χ ($\bar{\chi}$) has R -charge $R_X - 1$ ($R_X + 1$), *etc.* For the gauge invariant kinetic term of a chiral super field $\hat{\Phi}$ (irrespectively of its R -charge)

$$\mathcal{L} \ni \int d^2\theta d^2\bar{\theta} \hat{\Phi}^\dagger e^{-2g\hat{V}} \hat{\Phi} \quad (1)$$

to be R -invariant, the gauge vector super field $\hat{V}$ must be uncharged under R -symmetry. Since $\hat{V} = \dots - \bar{\theta}\sigma_\mu\theta V^\mu - i\bar{\theta}\bar{\theta}\theta\lambda + \dots$, the gauge vector field V^μ is uncharged, while the gaugino λ must carry R -charge $+1$. Then, the kinetic term for the gauge superfields

$$\mathcal{L} \ni \int d^2\theta \hat{W}\hat{W}, \quad (2)$$

where $\hat{W}$ stands for the gauge superfield stress tensor $\hat{W} \ni -i\lambda + \sigma_\mu\bar{\sigma}_\nu\theta F^{\mu\nu} + \dots$, is also automatically R -invariant. On the other hand, the Majorana gaugino mass terms $\frac{1}{2}M^M\lambda\lambda$ are forbidden in the soft-SUSY breaking Lagrangian. However, Dirac mass terms $M^D\lambda\lambda'$ are perfectly allowed if additional fermions λ' with opposite R -charge in the adjoint representations of each gauge group factor are introduced. This can be achieved by introducing gauge chiral-superfield adjoints $\hat{O}, \hat{T}, \hat{S}$ corresponding to $SU(3)_c$, $SU(2)_L$ and $U(1)$, respectively, each with R -charge 0. Such a construction amounts to promoting the gauge/gaugino sector to the $N = 2$ supersymmetric structure, which necessarily brings in new scalars, *i.e.* for each group factor, apart from gauge vector fields and gauge Dirac fermions, there are scalars in the adjoint representations. The MRSSM, therefore, contains sgluons — color-octet scalars, O , a scalar $SU(2)$ triplet T , and a scalar singlet S .

The assignment of R -charges to the matter chiral superfields is model dependent. In the Minimal R -symmetric Supersymmetric Standard Model (MRSSM) [5], it is done in such a way that all SM particles have R -charge 0 (in analogy to discrete R -parity). Thus, the left-chiral quark and lepton superfields have R -charge 1 and left-chiral Higgs superfields have R -charge 0.

With this assignment, the standard Yukawa terms in the superpotential are perfectly allowed, while all baryon- and lepton-number violating terms, as well as dimension-five operators mediating proton decay are forbidden. For the same reason, the standard Higgs/higgsino μ term is also forbidden. Therefore, to generate R -symmetric μ terms (and, consequently, higgsino mass terms), the Higgs sector of the MRSSM is extended by adding two iso-doublet superfields $\hat{R}_u$ and $\hat{R}_d$ with R -charge 2 (to be called R -Higgs). In all, the spectrum of fields in the R -symmetric supersymmetry theory consists of the standard MSSM matter, Higgs and gauge superfields augmented by chiral adjoints $\hat{O}, \hat{T}, \hat{S}$ and two R -Higgs iso-doublets. The R -charges of the superfields and their component fields are listed in Table I.

TABLE I

The R -charges of the superfields and the corresponding bosonic and fermionic components.

Field	Superfield		Boson		Fermion	
Gauge vector	$\hat{g}, \hat{W}, \hat{B}$	0	g, W, B	0	$\tilde{g}, \tilde{W}, \tilde{B}$	+1
Matter	$\hat{l}, \hat{e}^c$	+1	$\tilde{l}, \tilde{e}_R^*$	+1	l, e_R^*	0
	$\hat{q}, \hat{d}^c, \hat{u}^c$	+1	$\tilde{q}, \tilde{d}_R^*, \tilde{u}_R^*$	+1	q, d_R^*, u_R^*	0
H -Higgs	$\hat{H}_{d,u}$	0	$H_{d,u}$	0	$\tilde{H}_{d,u}$	-1
R -Higgs	$\hat{R}_{d,u}$	+2	$R_{d,u}$	+2	$\tilde{R}_{d,u}$	+1
Adjoint chiral	$\hat{O}, \hat{T}, \hat{S}$	0	O, T, S	0	$\tilde{O}, \tilde{T}, \tilde{S}$	-1

The MRSSM superpotential takes the following form

$$\begin{aligned}
 W = & \mu_d \hat{R}_d \cdot \hat{H}_d + \mu_u \hat{R}_u \cdot \hat{H}_u - Y_d \hat{d} \hat{q} \cdot \hat{H}_d - Y_e \hat{e} \hat{l} \cdot \hat{H}_d + Y_u \hat{u} \hat{q} \cdot \hat{H}_u \\
 & + \lambda_d \hat{S} \hat{R}_d \cdot \hat{H}_d + \lambda_u \hat{S} \hat{R}_u \cdot \hat{H}_u + \Lambda_d \hat{R}_d \cdot \hat{T} \hat{H}_d + \Lambda_u \hat{R}_u \cdot \hat{T} \hat{H}_u. \quad (3)
 \end{aligned}$$

Note that the Λ, λ -terms are similar to the usual Yukawa terms, where the $\hat{R}$ -Higgs and $\hat{S}$ or $\hat{T}$ play the role of the quark/lepton doublets and singlets. They will turn to be instrumental in achieving the required Higgs boson mass.

Turning to soft-SUSY breaking, the usual soft mass terms of the MSSM scalar fields are allowed just like in the MSSM. Similarly, the soft SUSY breaking B_μ , the Higgs, the adjoint scalar and R -Higgs scalar masses are consistent with R -symmetry. Although the holomorphic soft mass terms for the adjoint scalars, like $(m^2 SS + \text{h.c.})$, are also allowed, for simplicity, we will neglect them, as well as their trilinear couplings among themselves and to the

Higgs bosons since their presence does not influence our results significantly. On the other hand, all trilinear scalar couplings involving Higgs bosons to squarks and sleptons, which in the MSSM are notoriously unwanted sources of flavor violation, are removed since they carry non-vanishing R -charge. Likewise, the bilinear coupling of the R -Higgs has R -charge 4 and, therefore, is forbidden as well. The B_μ term is thus the only one which destroys the exchange symmetry between the H and R -Higgs fields. The soft-SUSY breaking scalar mass terms that we take, read

$$\begin{aligned}
 V_{\text{SB}} = & B_\mu (H_d^- H_u^+ - H_d^0 H_u^0) + \text{h.c.} \\
 & + m_{H_d}^2 \left(|H_d^0|^2 + |H_d^-|^2 \right) + m_{H_u}^2 \left(|H_u^0|^2 + |H_u^+|^2 \right) \\
 & + m_{R_d}^2 \left(|R_d^0|^2 + |R_d^+|^2 \right) + m_{R_u}^2 |R_u^0|^2 + m_{R_u}^2 |R_d^-|^2 \\
 & + m_S^2 |S|^2 + m_T^2 |T^0|^2 + m_T^2 |T^-|^2 + m_T^2 |T^+|^2 + m_O^2 |O|^2 \\
 & + \tilde{d}_{L,i}^* m_{q,ij}^2 \tilde{d}_{L,j} + \tilde{d}_{R,i}^* m_{d,ij}^2 \tilde{d}_{R,j} + \tilde{u}_{L,i}^* m_{q,ij}^2 \tilde{u}_{L,j} + \tilde{u}_{R,i}^* m_{u,ij}^2 \tilde{u}_{R,j} \\
 & + \tilde{e}_{L,i}^* m_{l,ij}^2 \tilde{e}_{L,j} + \tilde{e}_{R,i}^* m_{e,ij}^2 \tilde{e}_{R,j} + \tilde{\nu}_{L,i}^* m_{l,ij}^2 \tilde{\nu}_{L,j}. \quad (4)
 \end{aligned}$$

Although the familiar MSSM Weyl fermions $\tilde{B}, \tilde{W}, \tilde{g}$ cannot receive the soft Majorana masses, they can be paired with the corresponding fermionic component of the chiral adjoints $\tilde{S}, \tilde{T}, \tilde{O}$. When the Dirac mass terms are generated by D -type spurions, additional terms with auxiliary $\mathcal{D}$ -fields appear in the Lagrangian

$$\begin{aligned}
 V_D = & M_B^D \left(\tilde{B} \tilde{S} - \sqrt{2} \mathcal{D}_B S \right) + M_W^D \left(\tilde{W}^a \tilde{T}^a - \sqrt{2} \mathcal{D}_W^a T^a \right) \\
 & + M_O^D \left(\tilde{g}^a \tilde{O}^a - \sqrt{2} \mathcal{D}_g^a O^a \right) + \text{h.c.} \quad (5)
 \end{aligned}$$

When the auxiliary fields are eliminated via equations of motion, the Dirac mass parameters enter the scalar sector as well.

Since the R -Higgs fields carry non-vanishing R -charge, they do not develop vacuum expectation values (vev). The electroweak gauge symmetry breaking is triggered only by the vevs of neutral EW scalar fields, parameterized as

$$\begin{aligned}
 H_d^0 &= \frac{1}{\sqrt{2}}(v_d + \phi_d + i\sigma_d), & H_u^0 &= \frac{1}{\sqrt{2}}(v_u + \phi_u + i\sigma_u), \\
 T^0 &= \frac{1}{\sqrt{2}}(v_T + \phi_T + i\sigma_T), & S &= \frac{1}{\sqrt{2}}(v_S + \phi_S + i\sigma_S). \quad (6)
 \end{aligned}$$

The non-vanishing v_T contributes to the W -boson mass and shifts the ρ parameter away from one already at tree level. Therefore, experimental constraints put an upper limit $|v_T| \leq 4$ GeV.

Solving the tadpole equations for H_d, H_u , the soft masses $m_{H_d}^2$ and $m_{H_u}^2$ can be eliminated using v_d and v_u , and $v^2 = v_u^2 + v_d^2$ and $\tan \beta = v_u/v_d$ are defined as in the MSSM. The other two tadpole equations are solved for v_T and v_S , allowing us to use m_S^2 and m_T^2 as input parameters, which we assume to be positive to avoid tachyons.

The neutral scalar fields $(\phi_d, \phi_u, \phi_S, \phi_T)$ mix giving rise to the 4×4 Higgs boson mass matrix. The 2×2 sub-matrix for the (ϕ_d, ϕ_u) fields takes the same form as in the MSSM, and for large m_S^2, m_T^2 (which we take in the TeV range), the 2×2 sub-matrix for the ϕ_S, ϕ_T is approximately diagonal. The 2×2 off-diagonal sub-matrix that mixes the two sectors reads as

$$\mathcal{M}_{21} = \begin{pmatrix} v_d (\sqrt{2}\lambda_d \mu_d^+ - g_1 M_B^D) & v_u (\sqrt{2}\lambda_u \mu_u^- + g_1 M_B^D) \\ v_d (\Lambda_d \mu_d^+ + g_2 M_W^D) & -v_u (\Lambda_u \mu_u^- + g_2 M_W^D) \end{pmatrix}, \quad (7)$$

where the effective μ -parameter is given by $\mu_i^\pm = \mu_i + \frac{\lambda_i v_S}{\sqrt{2}} \pm \frac{\Lambda_i v_T}{2}$, $i = u, d$.

In general, the mixing between ϕ_d, ϕ_u and ϕ_S, ϕ_T leads to a reduction of the lightest tree-level Higgs boson mass compared to the MSSM. In [1], an approximate formula has been derived

$$m_{H_1}^2 \approx \left(m_Z^2 - v^2 \left(\frac{(g_1 M_B^D + \sqrt{2}\lambda\mu)^2}{4(M_B^D)^2 + m_S^2} + \frac{(g_2 M_W^D + \Lambda\mu)^2}{4(M_W^D)^2 + m_T^2} \right) \right) \cos^2 2\beta \quad (8)$$

in the limiting case of large m_A^2 and assuming $\lambda = \lambda_u = -\lambda_d$, $\Lambda = \Lambda_u = \Lambda_d$, $\mu_u = \mu_d = \mu$ and $v_S = v_T = 0$. It is clear that the MSSM upper limit of m_Z at tree level can be substantially reduced by terms depending on the new model parameters. Therefore, in the MRSSM loop corrections must play even more significant role, which we discuss in the next section.

3. Loop-corrected Higgs boson masses

The parameters of the model are renormalized in the $\overline{\text{DR}}$ scheme and v_d, v_u, v_S and v_T are given by the minimum of the loop-corrected effective potential. The pole mass m_{pole}^2 of a field is given by the pole of the full propagator

$$0 = \det \left[p^2 \delta_{ij} - \hat{m}_{ij}^2 + \Re \left(\hat{\Sigma}_{ij}(p^2) \right) \right]_{p^2 = m_{\text{pole}}^2}, \quad (9)$$

where p is the momentum, $\hat{m}^2$ the tree-level mass matrix and $\hat{\Sigma}(p^2)$ the finite part of the self-energy corrections. The one-loop self energies have been computed exactly using FeynArts [6], FormCalc [7] and Feynman rules generated by SARAH [8] properly modified to match our model. Since the

above equation cannot be solved analytically for m_{pole}^2 , the solution has to be found numerically. With SARAH, an MRSSM version of the SPheno spectrum generator [9] has been created to calculate the mass spectrum at full one-loop level. The results have been checked with a recent framework FlexibleSUSY [10].

Before presenting numerical results of full one-loop calculations, it is instructive to discuss the self energies in the effective potential approach. In the MSSM, the dominant one-loop contribution to the Higgs mass matrix comes from the top/stop sector. In the MRSSM, they are also important but because of the absence of stop mixing they are simpler

$$\Delta m_{H_1}^2 \sim \frac{6v^2}{16\pi^2} Y_t^4 \log \frac{m_{\tilde{t}_1} m_{\tilde{t}_2}}{m_t^2}$$

and, as a result, for the same stop mass, cannot reach the value as high as in the MSSM. Since the MRSSM superpotential contains new $\lambda_{u,d}$ and $\Lambda_{u,d}$ terms with a Yukawa-like structure, one can expect additional corrections proportional to λ^4, Λ^4 and logarithms of soft masses. Using the same approximation as in Eq. (8), the lightest Higgs state is given mainly by ϕ_u and only (ϕ_u, ϕ_u) component of the mass matrix needs to be computed and simple analytical expressions can be derived. For example, the λ^4 term gives the following contribution

$$\Delta m_{H_1}^2 \sim \frac{2v^2}{16\pi^2} \lambda^4 \log \frac{M_{R_u} m_S}{(M_B^D)^2} \quad (10)$$

with a similar structure of the Λ^4 and somewhat more complicated for the $\lambda^2 \Lambda^2$ terms.

For numerical analyses of one-loop corrections, no approximation is used and full dependencies on the parameters are taken into account. In [1], three representative benchmarks BMP1, BMP2, BMP3 have been identified with $\tan \beta = 3, 10, 40$, respectively, for which the Higgs boson mass can be met in accordance with LHC constraints, and with EW precision observables, to be discussed in the next section. The dependence of the lightest Higgs boson mass calculated at tree, one-loop (and two-loop) levels for one of the benchmarks, the BMP1, is shown in Fig. 1 as a function of one of the parameters, with all others set to the benchmark values. As already mentioned, the tree-level mass is significantly reduced below $m_Z \cos \beta$ for large values of λ_u, Λ_u , as the mixing between doublets and the singlet and the triplet gets enhanced. The dependence on λ_d, Λ_d is significantly weaker since the lightest Higgs gets dominant contribution from ϕ_u even for a low $\tan \beta = 3$. The size of one-loop top/stop Yukawa contribution alone can be judged from the value read at $\Lambda_u = \lambda_u = 0$, since then only stop/top

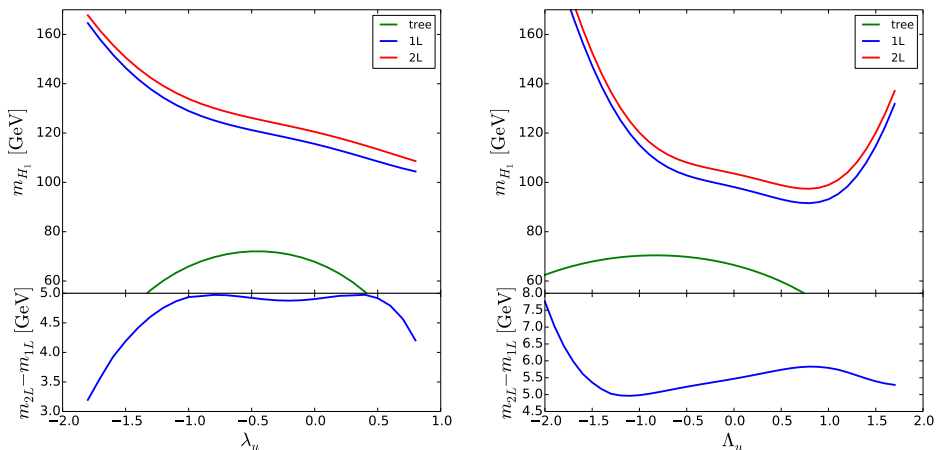


Fig. 1. The lightest MRSSM Higgs boson mass m_{H_1} , and the difference $m_{2L} - m_{1L}$ between masses calculated at the two-loop and one-loop level, as a function of λ_u , A_u , respectively. In the upper parts of the figure, lines from top to bottom correspond to two-loop, one-loop and tree level calculations. Other parameters are set to the values of benchmark point BMP1 with $\tan\beta = 3$ (from Ref. [2]).

contributions are significant. Evidently, a stop mass of 1 TeV, as set in the benchmark points, is not enough in the MRSSM to achieve the correct value of the Higgs boson mass, due to the absence of left–right mixing.

The full one-loop result shows large positive contributions from λ, A terms. Although the tree-level result falls quadratically with λ, A , as expected from Eq. (8), the one-loop result shows quartic dependence, as seen *e.g.* from Eq. (10), which explains the behavior of the sum. Thus, the λ, A one-loop contributions can push the Higgs boson mass to the measured value for values of λ_u, A_u close to unity.

Since the one-loop corrections are large, the question arises about the size of higher-order corrections. In Ref. [1], an estimate of higher-order corrections has been given with a conclusion that an expected two-loop contribution for the lightest Higgs boson mass should not exceed 6 GeV. This estimate has been verified in Ref. [2] using the recently updated **SARAH** code [11] that provides **SPheno** routines to calculate two-loop corrections in the effective potential approach and the gauge-less limit $g_{1,2} = 0$.

At two-loops, the λ, A corrections should behave in a manner similar to the pure top/stop two-loop contributions in the MSSM without stop mixing. And in fact, their numerical impact turns to be rather small, typically below 1 GeV, unless the couplings λ, A become very large, $|\lambda, A| \gg 1$. But at two-loops, also the strongly interacting sector and strong coupling α_s enter

directly to the Higgs boson mass predictions and these corrections can be expected to be sizable. Apart from the gluon, they involve the Dirac gluino and the sgluon, the scalar components of the octet superfield $\hat{O}$, which depend on the sgluon soft mass parameter m_O and the Dirac mass M_O^D . Note that M_O^D appears not only directly as the gluino mass but, via Eq. (5), also in couplings and mass terms of sgluons. In particular, it causes the splitting between the real and imaginary parts of the complex sgluon field $O = \frac{1}{\sqrt{2}}(O_S + iO_A)$, with tree-level masses $m_{O_S}^2 = 4(M_O^D)^2 + m_O^2$ and $m_{O_A}^2 = m_O^2$.

Compared to the MSSM, there are important differences due to the Dirac nature of the gluino, the lack of left–right stop mixing and the vanishing μ -parameter. For example, the diagrams with fermion mass insertions, corresponding to $\overline{FF}S$ -type contributions in the notation of Ref. [12], are not present in the MRSSM due to the absence of L–R mixing between squarks (for a comprehensive discussion of similarities and differences of two-loop results in the MSSM and MRSSM, see Ref. [1]).

Figure 2 shows the gluino mass dependence of the complete two-loop correction to the lightest Higgs boson mass. Curves are drawn for two different values of the sgluon mass parameter $m_O = 2$ and 10 GeV with all other

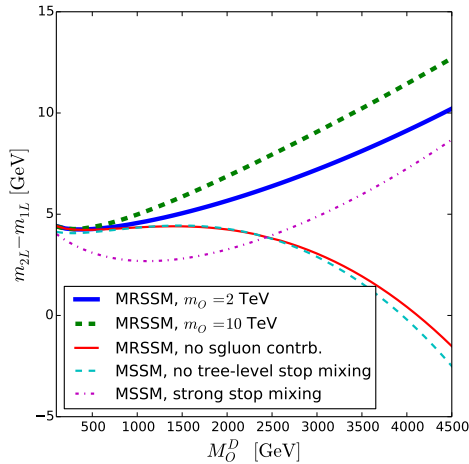


Fig. 2. (Color on line) Gluino mass dependence of two-loop contributions to the lightest Higgs boson mass in the MRSSM for BMP1. The results are shown for two different values of the soft sgluon mass parameter $m_O = 2$ TeV (thick solid/blue line) and 10 TeV (thick dashed/green line) with all contributions, respectively, and without the sgluon contributions (thin solid/red line). For comparison, also the MSSM contributions for no (thin dashed/light blue line) and maximal (purple/dotted line) stop mixing are plotted (from Ref. [2]).

parameters corresponding to BMP1. For comparison, the two-loop result without the sgluon contribution is shown as well, and the MSSM result with strong stop mixing and without any sfermion mixing at the tree-level.

In the MSSM without sfermion mixing, the gluino contribution is precisely the same as in the MRSSM since the Dirac or Majorana nature of the gluino does not matter as the Dirac partner, the octet superfield $\hat{O}$, has no direct couplings to quark superfields. This explains why the lines for the MRSSM without sgluon and the MSSM without stop mixing have very similar M_O^D dependence.

Including the sgluon diagram in the MRSSM strongly changes the behavior. Surprisingly, the full MRSSM two-loop contributions resemble the MSSM contributions with large stop mixing (corresponding to $X_t = 2000$ GeV), however for different reasons. In the MSSM, the increase is due to the additional $\overline{F}FS$ -type diagram which is directly proportional to gluino mass, while in the MRSSM, the sgluon diagram grows with M_O^D , both due to the sgluon–stop–stop coupling, which scales like M_O^D , and to an increase of the sgluon mass m_{O_S} . With the sgluon contributions, the total two-loop contributions to the Higgs boson mass in the MRSSM are larger than the ones in the MSSM. They are further increased by heavy sgluons. With the positive two-loop correction, a somewhat smaller value of the Λ_u is needed to meet the experimentally measured Higgs boson mass.

Overall, the two-loop contribution amounts to approximately +5 GeV, in agreement with previous estimate, and confirms the validity of the perturbative expansion in spite of the large one-loop result.

4. Electroweak observables

Since the non-vanishing v_T of the scalar triplet contributes to the tree-level W mass and shifts the ρ parameter from 1, it is constrained to be small: for all our benchmarks, it is below 1 GeV. Small v_T implies, through tadpole equations, a large value of the triplet soft mass and, consequently, somewhat split spectrum of Higgs bosons. In the SM and the MSSM, the top Yukawa coupling dominates loop corrections to m_W . Therefore, it should be expected that due to their Yukawa-like character, the λ, Λ couplings will also contribute at loop-level to electroweak observables (EWO), in particular to the W -boson mass.

Beyond tree-level, the W -boson mass can be obtained from the precisely measured muon decay constant using (hats denote $\overline{DR}$ -renormalized quantities in the MRSSM)

$$m_W^2 = \frac{1}{2} m_Z^2 \hat{\rho} \left[1 + \sqrt{1 - \frac{4\pi\hat{\alpha}}{\sqrt{2}G_\mu m_Z^2 \hat{\rho} (1 - \Delta\hat{r}_W)}} \right], \quad (11)$$

where $\hat{\rho}$ contains only oblique and $\Delta\hat{r}_W$ both oblique and non-oblique corrections which depend on the entire particle content of the model. The above formula also properly resums leading two-loop SM corrections [13]. The numerical calculation of $\hat{\rho}$ and $\Delta\hat{r}_W$ has been performed with the help of SARAH appropriately modified to account for the triplet scalar contribution.

It is convenient to rewrite the one-loop approximation to the W -boson mass in terms of the electroweak precision parameters S , T and U as

$$m_W = m_W^{\text{ref}} + \frac{\hat{\alpha} m_Z \hat{c}_W}{2(\hat{c}_W^2 - \hat{s}_W^2)} \left(-\frac{S}{2} + \hat{c}_W^2 T + \frac{\hat{c}_W^2 - \hat{s}_W^2}{4\hat{s}_W^2} U \right), \quad (12)$$

where $\hat{c}_W^2 = 1 - \hat{s}^2 = m_W^2/m_Z^2\hat{\rho}$, and m_W^{ref} is the W mass calculated in the SM. The advantage of computing S, T, U parameters is that they can be used in the calculation of several EWO. The main contribution to m_W from the MRSSM sector can be described in terms of the T parameter. It receives input from three sectors: charginos/neutralinos, Higgs and R -Higgs bosons. The contribution from R -Higgses has a similar structure to the stop/sbottom. Since in our benchmarks soft masses $m_{R_u}^2$ and $m_{R_d}^2$ are large, the mixing between R -Higgses and mass-splitting is small, leading to negligible contribution to T . The contribution from the triplet scalar is suppressed by the large soft triplet mass m_T^2 . The dominant contribution thus comes from the chargino/neutralino sector. For example, in the simplifying case of $\lambda_u = g_1$ and $\mu_u = M_W^D$, it can be written as

$$T = \frac{1}{16\hat{s}_W^2\hat{m}_W^2} \frac{v_u^4}{(M_W^D)^2} \times (4^{\text{th}} \text{ order polynomial in } g_2, A_u). \quad (13)$$

Figure 3 shows the A_u dependence of the full calculation of the W mass in the MRSSM (solid black line), as well as in various approximations. Other lines in the figure contain the full SM contribution, but the MRSSM contributions are taken into account either completely, or only via the T -parameter in various approximations, or from the tree-level triplet vev contribution. The figure shows that the chargino/neutralino approximation already gives an excellent approximation to the full T -parameter. The T -parameter, together with the tree-level triplet vev contribution, provides a good approximation to the full result. The remaining difference from non- T -parameter oblique corrections, vertex and box contributions, and leading higher loop contributions, is within ± 20 MeV, except for $|A_u| \gg 1.5$.

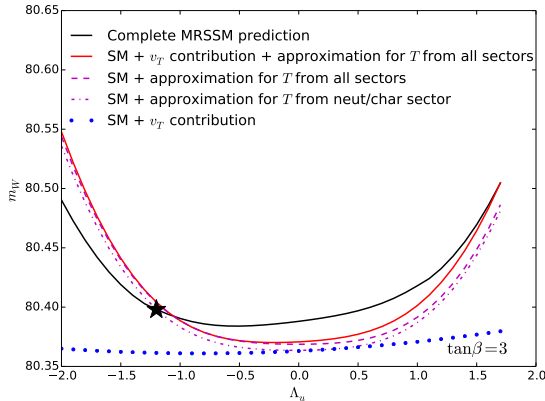


Fig. 3. The W -boson mass as a function of A_u , calculated using full MRSSM contributions and different approximations for the T -parameter for BMP1 (marked by the black star).

5. Conclusions and outlook

In my paper, I have discussed the structure of the R -symmetric supersymmetric extension of the Standard Model and recent progress in the precision calculation of the Higgs and W -boson masses. Compared to the MSSM, the model contains new states: R -Higgs $SU(2)$ -doublets and singlet, $SU(2)$ -triplet and $SU(3)$ -octet superfields, whose fermionic components allow us to write down the Dirac mass terms for gauginos and higgsinos.

We have seen that one can accommodate the observed Higgs boson mass in accordance with precision observables. The experimental values of m_{H_1} and W impose stringent and non-trivial constraints on the parameter space of the model. Nevertheless, it is easy to identify regions in the parameter space which accommodate the measured values and are in accordance with experimental data, as checked explicitly with **HiggsBounds** [14] and **HiggsSignals** [15], as well as selected low-energy flavor constraints. We have computed the full one-loop corrections to both m_{H_1} and W , and the two-loop correction to the Higgs mass in the effective potential approach. Numerical calculations have been cross-checked with analytic calculations of the most important new corrections. We have found that large scalar masses are favorable, of the order of 1–3 TeV. The resulting large mass ratios enhance loop-corrections to the lightest Higgs boson mass and suppress contribution from new states to the W -boson mass. The most instrumental are the new superpotential couplings λ, A , which play a role similar to the top/bottom Yukawa couplings with R -Higgses and singlet/triplet replacing quark doublets and singlets. With λ, A of the order of 1, like the top Yukawa coupling, the Higgs boson mass of ~ 125 GeV can easily be obtained even for top squarks below 1 TeV in spite of lack of L–R sfermion mixing.

The proposed benchmark points have many of the new states within the reach of the Run 2 of the LHC, in particular the supersymmetric fermions. It would be extremely exciting to see some of them in the current run of the LHC.

So far, we have exploited scenarios in which the lightest Higgs boson is the SM-like. In such cases, the mixing with new states lowers the tree-level m_{H_1} compared to the MSSM value, calling for even larger loop corrections to meet the measured value. However, one can contemplate an alternative scenario in which the lightest Higgs boson is mostly singlet, and the next one is the SM-like. In such a case, the second-lightest Higgs state gets pushed up via mixing already at the tree-level, thereby reducing the required loop corrections [16]. Similar scenarios have been considered in the next-to-minimal MSSM.

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