Raising the Higgs Mass in Supersymmetric Models

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ABSTRACT

The minimal supersymmetric standard model, and extensions, predicts a relatively light higgs particle if one supposes perturbativity until high scales. That fact is in conflict with nowadays data coming from LEPII fruitless searches for the Higgs boson. This bound is sufficient to rule out a big portion of the parameter space of the MSSM. In this letter we propose two mechanisms of raising the higgs mass at tree level, thus not needing very heavy sparticles that would reintroduce some fine-tuning in the problem.

1. Introduction

The minimal supersymmetric standard model (MSSM) has been the most promising candidate for physics beyond the standard model (SM) for the last two decades. It addresses one of the biggest problems in the SM, the hierarchy problem, the quadratic sensitivity of the higgs mass to high scales, while giving small contributions to the EW observables. At the same time it gives a natural scenario for unification and has a candidate for dark matter.

But the MSSM parameter space starts to be squeezed from the not-yet-found higgs boson. The higgs mass in the MSSM is predicted at tree-level to be smaller than $m_Z$, and this would have completely ruled it out had not been for radiative corrections. The upper bound on the higgs mass in the MSSM is around 130 GeV [1], and this suppose sparticles at 1 TeV and maximal mixing in the stop sector. These numbers start to reintroduce fine tuning in the problem, since the cancellation in order to get the correct value for the EW scale are on the order of few percent.

One way to alleviate that problem is achieved in the next to minimal supersymmetric standard model (NMSSM) where a new singlet coupled to the higgs is added. This singlet can have two roles, one is to solve the $\mu$ problem, the coincidence of the supersymmetric mass of the higgses with the EW scale, the other one is to raise the tree level value of the higgs mass. With nowadays bounds the first case tends to produce the same fine tuning as the MSSM, in the second case the bound on the higgs mass is lifted to 150 GeV [2] with the same assumptions as before. The gain is not too much because perturbativity up to $M_{GUT}$ impose the actual coupling of the singlet with the higgs to be small in the IR since it is not asymptotically free.

In this letter we will present two alternatives to increase the tree level value, both based on the same model, making use of asymptotically free interactions thus letting higher values of coupling in the IR. This letter is based on these two papers [3] and [4].
the reader should refer to those for further details and references.

2. D-Term (gauge extension)

The model we study is a gauge extension of the MSSM with gauge group $SU(3)_c \times SU(2)_1 \times SU(2)_2 \times U(1)_y$, the $SU(2)$ structure as follows, the third family plus the usual higgses are charged under the first $SU(2)$ whereas the first two families are charged under the second one, with this assignments the first $SU(2)$ is asymptotically free. Additional higgslike field exists on the second $SU(2)$ to generate appropriate yukawa interactions. To break the $SU(2)_1 \times SU(2)_2$ to the diagonal subgroup we use an extra bi-doublet $\Sigma$ which transforms as a $(2, 2)$. Above the scale of diagonal symmetry breaking, the $SU(2)_1 \times SU(2)_2$ D-term is

$$
\frac{g_2^2}{8} \left( \right. \left( \Sigma^\dagger \sigma^a \Sigma \right) + H^\dagger_2 \sigma^a H + \ldots \left) + \frac{g^2_2}{8} \left( \right. \left. \Sigma \sigma^a \Sigma^\dagger \right) \right)^2.
$$

(1)

The superpotential $W = \lambda S \left( \frac{1}{2} \Sigma \Sigma + w^2 \right)$ with an additional soft-mass $m^2$ for $\Sigma$ leads to the scalar potential

$$
V_\Sigma = \frac{1}{2} B \Sigma \Sigma + h.c. + m^2 |\Sigma|^2 + \frac{\lambda^2}{4} |\Sigma \Sigma|^2.
$$

(2)

Here, $\Sigma \Sigma$ is contracted with two epsilon tensors and $B = \lambda w^2$. For large $B$, $\Sigma$ acquires a VEV, $\langle \Sigma \rangle = u \mathbf{1}$, with $u^2 = (B - m^2)/\lambda^2$, which breaks $SU(2)_1 \times SU(2)_2$ to the diagonal subgroup. The minimum lies in a $D$-flat direction, leaving both Higgs fields massless.

Under the remaining $SU(2)$, $\Sigma$ contains a complex triplet, $T$, along with a complex singlet. Integrating out the real part of the heavy triplet at tree-level gives the effective Higgs potential below the triplet mass,

$$
\frac{g^2}{8} \Delta \left( H^\dagger_2 \bar{\sigma} H - \bar{\Pi} \sigma^a \Pi^a \right)^2 + \frac{g^2}{8} \left( \right. \left. |\Pi|^2 - |H|^2 \right)^2,
$$

with $\Delta = \frac{1 + \frac{m^2}{\bar{u}^2} \frac{1}{g_2^2}}{1 + \frac{m^2}{\bar{u}^2} \frac{1}{g_1^2 + g_2^2}}$ and $\frac{1}{g^2} = \frac{1}{g_1^2} + \frac{1}{g_2^2}$.

(3)

The MSSM D-term is recovered in the limit $u^2 \gg m^2$ (no SUSY breaking), for which SUSY protects the D-term below the gauge-breaking scale.

Electroweak symmetry breaking occurs under the same conditions as in the MSSM. We find the tree-level $W$ and $Z$ masses are corrected by the same relative amount, $(1 - g^4 v^2/2g^2_2 u^2 + \ldots)$ while the tree-level Higgs mass satisfies

$$
m^2_{h^0} < \frac{1}{2} \left( g^2 \Delta + g^2_1 \right) v^2 \cos^2 2\beta.
$$

(4)

To maximize the upper bound, $\Delta$ should be made as large as possible by sending $g_1 \to \infty$, $g_2 \to g$ and $m^2 \gg u^2$ by as much as possible without introducing fine-tuning. Also we have to be consistent with the EW fit.

We take the following example parameters:
• $g_1(u) = 1.80$, $g_2(u) = .70$, inspired by a GUT with $g_1(\Lambda_{GUT}) = .97$. Additional spectator fields are included in the running to aid in unification.

• $u = 2.4$ TeV, above the lower limit from electroweak constraints, giving $M_{W'}$, $M_{Z'} \sim 4.5$ TeV.

• $m = 10$ TeV. One-loop finite corrections to the Higgs mass parameter from supersymmetry breaking are $< 300$ GeV whereas two-loop RGE contributions can be somewhat larger if one assumes high-scale supersymmetry breaking.

We find $\Delta = 6.97$ and $m_h = 214$ GeV at tree-level in the large $\tan \beta$ and decoupling limits. Loop corrections to the effective potential from the top sector and the additional physics will make a relatively small shift in the tree-level result. This higgs mass is much greater that the MSSM or the NMSSM bound thus allowing for much larger parameter space and less fine-tuning.

3. F-term (NMSSM)

The second proposal to increase the higgs mass is similar to the usual NMSSM, we add an extra singlet with the following coupling in the superpotential to the higgs fields, $W = \lambda S H \bar{H}$, but with the same gauge structure of the previous section. In the original NMSSM, i.e. without gauge extension, the maximum value of $\lambda$ is relatively small if one supposes perturbativity to the GUT scale. The RGE of $\lambda$ can be written as:

$$\frac{d\lambda}{dt} = \frac{\lambda}{16\pi^2}(3|y_t|^2 + 4|\lambda|^2 - 3g^2)$$

where $y_t$ is the top yukawa and $g$ is the $SU(2)$ coupling. Since $g$ has a small value, the RGE for $\lambda$ is dominated by the first two terms and these make $\lambda$ non asymptotically free and hence pretty small in the IR. However in our framework $g$ should be replaced by the strong coupling $g_1$ and in this case the RGE evolution for $\lambda$ allows for much larger values in the IR.

The one-loop Higgs CP-even mass matrix can be written as a function of the CP-odd mass $(m_A)$, the $\mu$ parameter, the stop masses $m_{\tilde{t}}$, and the stop mixing parameter, $A_t$:

$$m_{11}^2 = m_Z^2 \cos^2 \beta + m_A^2 \sin^2 \beta - \frac{3 y_t^2}{16\pi^2} \mu^2 \frac{Z^2}{3}$$

$$m_{22}^2 = m_Z^2 \sin^2 \beta + m_A^2 \cos^2 \beta + \frac{3 y_t^2}{16\pi^2} \left(4m_t^2 \log \frac{m_{\tilde{t}} m_{\tilde{\tilde{t}}}}{m_t^2} + A_t(2m_t Z - A_t \frac{Z^2}{3})\right)$$

$$m_{12}^2 = -\frac{1}{2}(m_Z^2 + m_A^2 - 2\nu^2 \lambda^2) \sin 2\beta + \frac{3 y_t^2}{16\pi^2} \mu \left(m_t Z - A_t \frac{Z^2}{3}\right)$$

where:

$$Z = \frac{m_t(A_t + \mu \cot \beta)}{m_t^2 + \frac{1}{2}(m_Q^2 + m_U^2)}$$
and the stop masses are defined with respect to the soft-masses for left- and right-handed stops \((m_Q, m_U)\) as:

\[
m_{\tilde{t}_1,2}^2 = m_t^2 + \frac{1}{2}(m_Q^2 + m_U^2) \pm W
\]

\[
W^2 = \frac{1}{4}(m_Q^2 - m_U^2)^2 + y_t^2 v^2 |A_t \sin \beta - \mu \cos \beta|^2
\]  

(8)

The charged Higgs mass is (at one-loop):

\[
m_H^2 = m_A^2 + m_W^2 - v^2 |\lambda|^2.
\]  

(9)

From the above formulae there are several consequences that can be drawn. In the decoupling limit \((m_A \to \infty)\) and \(\tan \beta \sim 1\) the lightest higgs mass is \(\lambda v\) which can be as large as 214 GeV. In contrast to the usual MSSM and NMSSM where \(\tan \beta\) has to be greater than 1.8 due to the perturbativity of \(y_t\), in our case much lower values of \(\tan \beta\) are allowed since \(g_1\) also affects the RGE for \(y_t\). There are regions of the parameter space where the lightest of the higgs bosons is the charged higgs giving rise to interesting new signatures in colliders since the branching ratios are very different to those of the MSSM. Finally the LEPII bound of 114 GeV can be avoided for small values of \(m_A\) because in that case the scalar that couples to \(W\) and \(Z\) bosons is the heaviest hence that bound should apply to that higgs and not to the lightest.

4. Conclusions

The LEPII bound on the higgs mass makes the MSSM a fine-tuned theory. In this letter the bound is avoided thanks to a new strong interaction in the gauge sector broken at a similar scale as SUSY. In one case a non-decoupling D-term and in the second case a new superpotential interaction (NMSSM like) with large IR value give a new contribution to the higgs tree-level mass and large sparticle masses are not needed to accomplish the experimental bound. Both scenarios have interesting collider signatures.

5. References

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