THE LIGHT HIGGS IN SUPERSYMMETRIC MODELS WITH HIGGS TRIPLETS

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ABSTRACT

In supersymmetric models the presence of Higgs triplets introduce new quartic interactions for the doublets that may raise the mass of the lightest $CP$-even field up to 205 GeV. We show that the complete effect of the triplets can be understood by decoupling them from the minimal sector and then analyzing the vacuum and the spectrum of the effective two-Higgs doublet model that results. We find that the maximum value of $m_h$ is only achieved in a very definite region of the parameter space. In this region, however, radiative corrections decrease the bound to $\approx 190$ GeV.
1. Introduction. In the minimal supersymmetric standard model (MSSM) the quartic couplings of the Higgs fields are not free parameters, they are related to the $SU(2)_L \times U(1)_Y$ gauge couplings $[1]$. This implies that the tree-level mass of the lightest $CP$-even scalar field is smaller than $M_Z$:

$$m^2_h \leq M^2_Z \cos^2 2\beta,$$

where $\tan \beta$ is the ratio $\langle H_2 \rangle / \langle H_1 \rangle$ of vacuum expectation values (VEVs) of the two Higgs doublets.

There are two different mechanisms that can raise this upper bound on $m^2_h$. The first one is due to the fact that supersymmetry (SUSY) is broken. The SUSY values of the quartic couplings suffer radiative corrections $[2]$ from top loops proportional to $h_t^4 \ln(m^2_t/m^2_\tilde{t})$, where $h_t$ is the top Yukawa coupling and $m_t$ and $m_\tilde{t}$ are the fermion and the scalar masses, respectively. This translates into corrections to the Higgs mass of order $\Delta m^2_h \approx (90 \text{ GeV})^2$.

Also the soft SUSY-breaking trilinear $V \ni A_t \tilde{t} H^0_2$ may have an impact on $m^2_h$. At one loop it gives nonlogarithmic corrections $[3]$ proportional to $h_t^2 A_t^2 m^2_\tilde{t} - \frac{1}{12} A_t^4 m^4_\tilde{t}$ that can be of the same size.

The other mechanism to increase the bound on $m^2_h$ requires the presence of extra Higgs fields and couplings, namely, gauge singlets $[4]$ or triplets $[5]$ and couplings of these fields with the doublets. Trilinears in the superpotential $W$ involving two doublets and the extra field introduce new quartic interactions for the scalar doublets and then new contributions to $m^2_h$.

When introducing extra Higgs fields, however, there is another effect competing with the positive effect on $m^2_h$ of the quartic couplings. It is due to the mixing between the lightest and heavier states. Such a mixing tends to decrease the smallest eigenvalue in the Higgs mass matrix. In a particular scenario one would expect that the new (arbitrary) parameters present, soft masses and trilinears, may be adjusted in order to cancel the mixing. In that case, for real VEVs the lightest neutral state would be obtained diagonalizing the $2 \times 2$ submatrix defined by the $CP$-even scalars in the doublets. Then the maximum value of $m_h^2$ would be obtained for heavy $CP$-odd states.

In particular, let us consider the presence of a pair of $SU(2)_L$ triplets $(T, \overline{T})$ of $(-1, +1)$ hypercharge:

$$T = \begin{pmatrix} T^0 \\ T^- \\ T^{--} \end{pmatrix}; \quad \overline{T} = \begin{pmatrix} T^{++} \\ T^+ \\ T^0 \end{pmatrix}.$$  \(2\)

These fields admit a term in the superpotential of type

$$W \ni \frac{1}{2} \chi \left( T^0 H^0_2 H^0_2 - \sqrt{2} T^- H^0_2 H^+_2 + T^{--} H^+_2 H^+_2 \right).$$  \(3\)
An analogous term can be obtained exchanging $T \leftrightarrow \overline{T}$, $H_2 \leftrightarrow H_1$. With the quartic coupling $V \supset \lambda_2 H_2^0 H_2^0$ in the scalar potential the bound on $m_h^2$ becomes

$$m_h^2 \leq M_Z^2 \cos^2 2\beta + \chi^2 v^2 \sin^4 \beta ,$$  

(4)

where $v^2 = \langle H_1 \rangle^2 + \langle H_2 \rangle^2 \approx (174 \text{ GeV})^2$. Although $\chi$ is in principle a free parameter, it is constrained by the following argument. Its evolution with the energy scale is given by

$$\frac{d\chi}{dt} = \frac{\chi}{16\pi^2} \left( \frac{7}{2} \chi^2 + 6 h_t^2 - \frac{9}{5} g_1^2 - 7 g_2^2 \right) .$$  

(5)

The possibility to integrate the electroweak and the grand unification (GUT) scales is here a main motivation, as in any SUSY model. However, large initial values of $\chi$ would become nonperturbative before the GUT scale $M_X$. Evolving the model down from $M_X$ one finds that at low energies $\chi$ is always smaller than $\approx 0.9$. This value could be enough to put the bound in Eq. (4) around $205 \text{ GeV}$.

Our objective in this letter is to find out whether the bound in Eq. (4) can be saturated or not. We discuss a simple way to understand the effect of adding extra triplets (or any other vectorlike Higgs field) of mass $m = M + O(m_{SUSY})$. We decouple these scalar fields from the minimal sector keeping the terms of first order in $1/m$. The effective model that results is a particular two-Higgs doublet model that depends on the physics of the triplets. The decoupling effects are equivalent to the mixing in the complete mass matrix, but much simpler to analyze. The method is justified when $m^2 \gg (v^2, \frac{\chi^2}{4\pi} m_{SUSY}^2)$. In the triplet model this seems to be a necessary requirement because the VEVs grow as the inverse of their mass, and if $\langle T, \overline{T} \rangle \geq 10 \text{ GeV}$ it predicts an unacceptable value of the $\rho$ parameter (a different Weinberg angle measured from gauge boson masses and charged currents). We compare the value for $m_h$ obtained in the effective model with the exact numerical solution in the complete triplet model and find that the agreement is good even for $m \approx m_{SUSY}$. The method allows us to identify the only region in the parameter space that saturates the bound in Eq. (4). This region has a very definite pattern of SUSY breaking terms, with large radiative corrections that partially cancel top quark effects.

2. The triplet model. Let us consider the neutral Higgs sector of the model (from now on we drop the 0 superscript to indicate neutral fields). We include in $W$ the terms

$$W \supset -\mu H_1 H_2 - M T \overline{T} + \frac{\chi}{2} T H_1 H_2 .$$  

(6)

Adding soft SUSY-breaking terms the relevant part of the scalar potential is

$$V = m_1^2 H_1^1 H_1 + m_2^2 H_2^1 H_2 + m_3^2 T^1 T + m_4^2 \overline{T}^1 \overline{T}$$
\[-(m_{i_2}^2 H_1 H_2 + \text{h.c.}) + (A_M^2 T T + \text{h.c.}) - (A_\chi T H_2 H_2 + \text{h.c.}) \]
\[-\left(\frac{1}{2} \chi M \mathcal{T} H_2 H_2 + \text{h.c.} \right) - (\chi \mu H_1^\dagger T H_2 + \text{h.c.}) + \chi^2 H_2^\dagger T H_2 T \]
\[+ \frac{1}{4} \chi^2 (H_2^\dagger H_2)^2 + \tilde{g} (H_1^\dagger H_1 - H_2^\dagger H_2 + 2 T^\dagger T - 2 \mathcal{T} \mathcal{T})^2 , \]

(7)

where \( \tilde{g} = (g_1^2 + g_2^2)/8 \) and all the fields are neutral. Field redefinitions can be used to set \( m_{i_2}, A_\chi, \) and \( \chi M \) real and positive. We assume for simplicity that \( A_M^2 \) and \( \chi \mu \) are real, but this does not guarantee that all the VEVs are real and positive (which would be the case if \( \chi \mu \geq 0 \) and \( A_M^2 \leq 0 \)). The size of the mass parameters above is \( (m_1^2, m_2^2, m_{i_2}^2) = O(m_{\text{SUSY}}^2); (\mu, A_\chi) = O(m_{\text{SUSY}}); M \geq O(m_{\text{SUSY}}); (m_3^2, m_4^2) = M^2 + O(m_{\text{SUSY}}^2); \) and \( A_M^2 = O(m_{\text{SUSY}} M) \).

It is convenient to express the fields in terms of moduli and phases:
\[ H_1 = \frac{1}{\sqrt{2}} v_1 e^{\theta_1} ; \quad H_2 = \frac{1}{\sqrt{2}} v_2 e^{\theta_2} ; \]
\[ T = \frac{1}{\sqrt{2}} v_3 e^{\theta_3} ; \quad \mathcal{T} = \frac{1}{\sqrt{2}} v_4 e^{\theta_4} . \]

(8)

The determination of the minimum of the potential in Eq. (7) and of \( m_h^2 \) requires much algebra (in the general case with complex VEVs, the diagonalization of a \( 8 \times 8 \) matrix). We propose, instead, to integrate the fields \( (T, \mathcal{T}) \) out, analyze the effective model that results, and check numerically that for any particular choice of parameters the complete and the approximate models give the same spectrum for the four lightest fields.

3. The approximate model. We first rewrite \( (T, \mathcal{T}) \) in terms of mass eigenstates:
\[ T_1 = c_\alpha T - s_\alpha \mathcal{T}^\dagger ; \quad T_2 = s_\alpha T + c_\alpha \mathcal{T}^\dagger , \]

(9)

where \( s_\alpha = \sin \alpha, c_\alpha = \cos \alpha \) and \( \tan 2 \alpha = (2 A_M^2)/(m_4^2 - m_3^2) \). Their masses are
\[ M_1^2 = c_\alpha^2 m_3^2 + s_\alpha^2 m_4^2 - 2 s_\alpha c_\alpha A_M^2 ; \quad M_2^2 = s_\alpha^2 m_3^2 + c_\alpha^2 m_4^2 + 2 s_\alpha c_\alpha A_M^2 . \]

(10)

If \( M \gg m_{\text{SUSY}} \) then \( \alpha \approx \frac{\pi}{4}, M_1^2 \approx M^2 - A_M^2 \) and \( M_2^2 \approx M^2 + A_M^2 \).

Integrating \( T_1 \) and \( T_2 \) out it results the two Higgs doublet model
\[ V = m_1^2 H_1^\dagger H_1 + m_2^2 H_2^\dagger H_2 - (m_{i_2}^2 H_1 H_2 + \text{h.c.}) \]
\[ + \lambda_3 H_1^\dagger H_1 H_2^\dagger H_2 + \left( \frac{1}{2} \lambda_7 H_2^\dagger H_2 H_1 H_2 + \text{h.c.} \right) + \frac{1}{4} \chi^2 (H_2^\dagger H_2)^2 \]
\[ + \tilde{g} (H_1^\dagger H_1 - H_2^\dagger H_2)^2 . \]

(11)
where the contributions of the triplets to $\tilde{\chi}^2$, $\lambda_3$ and $\lambda_7$ come from the diagrams (a), (b) and (c) in Fig. 1, respectively. We obtain

$$
\tilde{\chi}^2 = \chi^2 - \frac{(\chi M s_\alpha - 2A \chi c_\alpha)^2}{M_1^2} - \frac{(\chi M c_\alpha + 2A \chi s_\alpha)^2}{M_2^2},
$$

$$
\lambda_3 = -\frac{(\chi \mu c_\alpha)^2}{M_1^2} - \frac{(\chi \mu s_\alpha)^2}{M_2^2},
$$

$$
\lambda_7 = -\frac{(\chi \mu c_\alpha)(2A \chi c_\alpha - \chi M s_\alpha)}{M_1^2} - \frac{(\chi \mu s_\alpha)(2A \chi s_\alpha + \chi M c_\alpha)}{M_2^2}.
$$

(12)

Higher dimensional operators would introduce corrections of order $v^2/M_{1,2}^2$, whereas one-loop corrections will be small if $\frac{v^2}{M_{SUSY}} < M_{1,2}^2$.

To find the minimum and the spectrum of this model we express the fields in terms of moduli and phases:

$$
H_1 = \frac{1}{\sqrt{2}} v_1 e^{\theta_1}; \quad H_2 = \frac{1}{\sqrt{2}} v_2 e^{\theta_2}.
$$

(13)

Then

$$
V = \frac{1}{2} m_{12}^2 v_1^2 + \frac{1}{2} m_{22}^2 v_2^2 - m_{12}^2 v_1 v_2 \cos(\theta_1 + \theta_2)
+ \frac{1}{4} \lambda_3 v_1^2 v_2^2 + \frac{1}{4} \lambda_7 v_1 v_2^3 \cos(\theta_1 + \theta_2) + \frac{1}{16} \tilde{\chi}^2 v_2^4
+ \frac{\tilde{g}}{4} (v_1^2 - v_2^2)^2.
$$

(14)

We can use an hypercharge transformation to set $\langle \theta_1 \rangle = 0$. The minimum conditions give then $\langle \theta_2 \rangle = 0$, $\langle v_1 \rangle$ and $\langle v_2 \rangle$. In the $4 \times 4$ mass matrix the $CP$-odd sector, $M_{ij} = \frac{1}{v_i v_j} \frac{\partial^2 V}{\partial \theta_i \partial \theta_j}$, does not mix with the the $CP$-even sector, $M_{2ij} = \frac{\partial^2 V}{\partial \theta_i \partial \theta_j}$. We find

$$
M_{11} = \tilde{m}_{12}^2 \tan \beta,
M_{12} = \tilde{m}_{12}^2,
M_{22} = \tilde{m}_{12}^2 \tan^{-1} \beta;
M_{33} = \tilde{m}_{12}^2 \tan \beta + M_Z^2 \cos^2 \beta,
M_{34} = -\tilde{m}_{12}^2 - M_Z^2 \sin \beta \cos \beta + 2\lambda_3 \sin \beta \cos \beta + \lambda_7 v^2 \sin^2 \beta,
M_{44} = \tilde{m}_{12}^2 \tan^{-1} \beta + M_Z^2 \sin^2 \beta + \tilde{\chi}^2 v^2 \sin^2 \beta + \frac{3}{2} \lambda_7 v^2 \sin \beta \cos \beta;
$$

(15)

where $\tilde{m}_{12}^2 = m_{12}^2 - \frac{\lambda_7}{4} (v_2)^2$, $v^2 = \frac{\langle m \rangle^2 + (\langle m \rangle)^2}{2}$, $\tan \beta = \frac{\langle v_2 \rangle}{\langle v_1 \rangle}$ and $M_Z^2 = 4\tilde{g} v^2$ (we neglect the contribution of triplet VEVs to $M_Z$).
It is now straightforward to find the mass eigenvalues. In the $CP$-odd sector there is, in addition to the massless Goldstone, a field of mass $m_A^2 = \tilde{m}_{12}^2 / (\sin \beta \cos \beta)$. The lightest Higgs is in the $CP$-even sector, together with a field of mass $m_H^2 = O(\tilde{m}_{12}^2)$. $m_h^2$ is bounded to be smaller than $(M_{33} M_{44} - M_{34}^2) / (M_{33} + M_{44})$, value that is saturated in the limit $\tilde{m}_{12}^2 \gg v^2$.

In this limit we can obtain an approximate expression for $m_h^2$:

$$m_h^2 \approx M_Z^2 \cos^2 2\beta + \tilde{\chi}^2 v^2 \sin^4 \beta + 4\lambda_\chi v^2 \sin^2 \beta \cos^2 \beta + \frac{7}{2} \lambda_\gamma v^2 \sin^3 \beta \cos \beta ,$$

where the couplings have the value specified in Eq. (12).

Before discussing how to tune the parameters in order to approach the bound in Eq. (4), let us check the efficiency of our approach. We have computed numerically the spectrum of light fields in the effective and in the complete triplet models for many different values of the parameters, changing the signs of $A_3^M$ and $\chi \mu$. We obtain that the results for $m_h$ in both models always agree within a 2% margin. For example, let us take a triplet model with $M = 3$ TeV, $\chi = 0.7$, $A_3^M = (0.5$ TeV$)M$, $A_\chi = 0.5\chi$, $\mu = 0.4$ TeV, $m_{12}^2 = 0.25$ TeV$^2$, $m_3^2 = 9.4$ TeV$^2$, $m_4^2 = 9.5$ TeV$^2$, and the mass parameters $m_1^2 = 0.50$ TeV$^2$ and $m_2^2 = 0.12$ TeV$^2$ (chosen to have $\tan \beta = 2$ and the right value of $M_Z$). The spectrum of this model is (3310, 2819, 790.7) GeV in the $CP$-odd sector and (3310, 2819, 794.3, 51.4) GeV in the $CP$-even sector. The value in the MSSM that corresponds to this value of $\tan \beta$ is $m_h = 55.2$ GeV, versus $m_h = 51.4$ GeV obtained here (both values would coincide if the triplets were completely decoupled). Defining an analogous model with opposite sign for $\chi \mu$ (i.e., $\mu = -0.4$ TeV) we obtain $m_h = 59.3$ GeV.

The corresponding two-Higgs doublet model is built with the couplings in Eq. (12): $\tilde{\chi} = 0.10$, $\lambda_3 = -0.0086$, $\lambda_7 = -0.0113$, and the masses $m_1^2 = 0.50$ TeV$^2$, $m_2^2 = 0.12$ TeV$^2$.

In order to keep the same value of $M_Z$ and $\tan \beta = 2$, these masses are not identical to $m_{1,2}$ in the complete triplet model (however, the difference is in next digits). Here we obtain a field of (790.8) GeV in the $CP$-odd sector and fields of (794.4, 51.2) GeV in the $CP$-even sector. The value obtained for the mass of the lightest Higgs, $m_h = 51.2$ GeV, is very close to the value $m_h = 51.4$ GeV of the complete triplet model. In the approximate model that corresponds to $\mu = -0.4$ TeV we obtain $m_h = 59.2$ GeV, also in agreement with the value $m_h = 59.3$ GeV of the complete model. For this choice of $\chi$ and $\tan \beta$ the bound in Eq. (4) is $m_h \leq 111.9$ GeV, a number that gives no information.

We also obtain an excellent approximation when the triplet fields to integrate are not heavier than the doublets. For example, taking $m_1^2 = 1.12$ TeV$^2$, $m_2^2 = 0.05$ TeV$^2$, $m_3^2 = 0.95$ TeV$^2$, $m_4^2 = 1.00$ TeV$^2$, $M = 0.7$ TeV, $\chi = 0.7$, $A_3^M = (0.5$ TeV$)M$, $A_\chi = 0.5\chi$, $\mu = -0.4$ TeV, and $m_{12}^2 = 0.25$ TeV$^2$ we have $\tan \beta = 5$, $CP$-odd scalars of (1175, 1119, 791).
GeV and CP-even scalars of (1175, 1119, 791, 91.9) GeV. Actually, those values correspond to a local minimum in an unbounded potential; this is the tendency (due to the large number of complex phases) in most of the parameter space when \( M \leq m_{\text{SUSY}} \). In this case the triplets to integrate out are lighter than one of the doublets. In the approximate model we have \( \tilde{\chi} = -0.296, \lambda_3 = -0.095 \) and \( \lambda_7 = 0.18 \). We obtain \( m_h = 91.2 \) GeV, in agreement with the result \( m_h = 91.9 \) GeV in the complete model (versus \( m_h = 84.9 \) GeV in the MSSM or \( m_h \leq 144.7 \) GeV in Eq. (4)). The numerical analysis shows that the effective two-Higgs doublet model describes very efficiently the effect of the triplets on \( m_h \).

4. Maximum value of \( m_h \).

Now, from the expressions in Eqs. (12,16) it is clear that for a SUSY triplet mass \( M \geq m_{\text{SUSY}} \) the bound in Eq. (13) is never approached. For \( M \) much larger than the SUSY-breaking masses \( \tilde{\chi}^2 \) goes to zero, and the triplets decouple (as expected). If \( 2A_\chi c_\alpha \) is large and tends to cancel \( \chi M s_\alpha \) in one of the terms defining \( \tilde{\chi}^2 \), then in the other term \( \chi M c_\alpha + 2A_\chi s_\alpha \) will be large and \( \tilde{\chi}^2 \) goes also to zero. The contribution to \( m_h^2 \) proportional to \( \lambda_3 \) is always negative, whereas the one proportional to \( \lambda_7 \) can be positive if \( \chi \mu < 0 \). However, in this case a sizeable contribution would require that all the couplings (\( \chi M, 2A_\chi \) and \( \chi \mu \)) are of the same order, implying complex VEVs and mixing of the light Higgs with the (heavy) CP-odd sector. Such a mixing also would lower \( m_h^2 \). In any case, the \( \lambda_3,7 \) terms are not relevant in the region of large \( \tan \beta \), where the bound in Eq. (14), if saturated, allows a light Higgs of up to 205 GeV [3].

The maximum value of \( m_h \) would be obtained for a SUSY mass \( M \) and a SUSY-breaking trilinear \( A_\chi \) both much smaller than the SUSY-breaking masses of the triplets: \( M^2 \ll (M_1^2, M_2^2), 2A_\chi^2 \ll (\chi^2 M_1^2, \chi^2 M_2^2) \). In this limit the scalar triplets decouple but the quartic coupling that they introduce in the Higgs doublet sector remains. In consequence, the tree-level bound in Eq. (11) would be approached.

This very definite pattern of SUSY breaking parameters, however, has obvious implications at the quantum level. We have here a large splitting between the fermion and the scalar components of the triplet superfields and also a large Yukawa coupling (see Eq. (16)), very much like in the top quark sector. To estimate the radiative effects let us focus on the region of large \( \tan \beta \), where the light neutral Higgs \( \phi \) is basically \( H_2 \). We simplify and assume that \( M \approx m_t \) and all the SUSY-breaking masses coincide, \( m_{\text{SUSY}} \approx 1 \) TeV, with a generic suppression of scalar trilinears that makes negligible nonlogarithmic corrections. Below \( m_{\text{SUSY}} \) we have the standard model, with \( V = m_\phi^2 \phi^\dagger \phi + \frac{1}{2} (\phi^\dagger \phi)^2 \), plus the fermion components of the triplets. At \( m_{\text{SUSY}} \) the quartic scalar coupling is

\[
\lambda_0 = \frac{g_Y^2 + g_L^2}{2} + \chi^2 ,
\]  

(17)
and its running down to the electroweak scale is given by

$$\frac{d\lambda}{dt} = \frac{3}{4\pi^2} \left[ \lambda^2 + (2h_t^2 + \chi^2)\lambda - 4h_t^4 - \frac{5}{3}\chi^4 \right].$$  (18)

In the MSSM one obtains $\lambda(m_t^2) \approx \lambda_0 + \frac{3}{16\pi^2} \ln \frac{m_{\text{SUSY}}^2}{m_t^2}$ and $m_h^2 \approx (92^2 + 88^2)$ GeV$^2$. Here there is a partial cancellation in the $\beta$ function, with

$$\lambda(m_t^2) \approx \lambda_0 + \frac{3}{16\pi^2} (4h_t^4 - 2h_t^2\chi^2 - \frac{1}{3}\chi^4) \ln \frac{m_{\text{SUSY}}^2}{m_t^2}$$  (19)

and

$$m_h^2 \approx M_Z^2 + \chi^2v^2 + \frac{3}{16\pi^2}(4h_t^4 - 2h_t^2\chi^2 - \frac{1}{3}\chi^4)v^2 \ln \frac{m_{\text{SUSY}}^2}{m_t^2} \approx (92^2 + 156^2 + 62^2) \text{ GeV}^2 = (190 \text{ GeV})^2,$$  (20)

where we have neglected the electroweak gauge couplings and the evolution of $h_t$ and $\lambda$ with the scale (both with decreasing effect on the estimated size of radiative corrections).

4. Conclusions. The mass of the Higgs in the MSSM is constrained to be smaller than $M_Z$ at the tree level and smaller than around 130 GeV once radiative (SUSY-breaking) corrections are included. In more general models there are new fields and new quartic Higgs interactions raising $m_h$. However, the new fields also introduce mixing with the Higgs doublets, which decreases $m_h$. In models with gauge singlets it is easy to see that this mixing can be fine tuned to zero. However, in models with triplets the larger number of fields and parameters makes the analysis too complicated.

We have presented a method that allows to understand the complete effect of the triplets on $m_h$. It is based on an effective model that results integrating the triplets out and keeping only their effect on the quartic couplings of the doublets. In this effective model the bound on $m_h$ is very simple to estimate (see Eq. (18)). The procedure gives an excellent approximation for triplet masses larger than the electroweak scale $v \approx 174$ GeV, even if the triplets and the doublets have similar masses.

Using this method we show that, in order to modify substantially the MSSM values of $m_h$, the scalar triplets must have a small SUSY mass (their mass must be basically a SUSY-breaking term) and their trilinears must be suppressed respect the masses. The large splitting in the triplet supermultiplet, with heavy scalars and light fermions, together with the large Yukawa coupling $\chi$ and the absence of scalar trilinears, define a very clear pattern of SUSY-breaking parameters with implications on the size of radiative corrections. We have estimated these corrections and obtained that they are significantly smaller than in the MSSM.
We conclude that the triplet model is able to provide values of $m_h$ of up to 190 GeV. The value 205 GeV seems an overestimate, since it would require $M \approx m_{SUSY}$ and large SUSY-breaking trilinears (it is based on a cancellation between trilinears of order $m_{SUSY}$ that we show cannot take place). The value $m_h = 190$ GeV is still larger than $m_h = 155$ GeV of the singlet model with intermediate vectorlike matter. Probably, the triplet model has ingredients that make it a less appealing framework from a model building point of view: the need to avoid triplet VEVs, the need for four pairs of colour triplets at low energy to obtain gauge unification, the need to avoid extra matter with electroweak charges at intermediate scales ($g_Y$ and $g_L$ are near their perturbative fixed point values), or the need to incorporate a generalized R-parity to be realistic (the usual $Z_2$ matter parity of the MSSM cannot avoid here, for example, unacceptable neutrino masses). However, the pattern of SUSY-breaking terms (large scalar masses versus trilinears) and the presence of double charged leptons at $\approx 200$ GeV required to saturate the bound define an interesting region of its parameter space. For example, there the triplet VEVs are naturally small:

$$\langle v_4 \rangle \approx \frac{\sqrt{2} \chi M v^2 \sin^2 \beta}{m^2_4},$$

(21)
easily within the experimental limit of 10 GeV.

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Figure 1: Diagrams contributing to $\tilde{\chi}^2$ (a), $\lambda_3$ (b) and $\lambda_7$ (c).