



## Higgs Boson Mass in Models with Gauge-Mediated Supersymmetry Breaking

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### Abstract

We present the predictions for the mass  $M_h$  of the lightest Higgs boson in models with gauge-mediated supersymmetry breaking as a function of the SUSY-breaking scale. We include all radiative corrections up to two loops and point out that if the CDF  $e^+e^-\gamma\gamma$  event is interpreted in terms of these models, then the lightest Higgs boson should be lighter than 110 GeV.

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At present, supersymmetry (SUSY) is widely regarded as a leading candidate for physics beyond the Standard Model [1]. Although this is largely due to the fact that SUSY provides the only known perturbative solution to the problem of quadratic divergence in the standard-model Higgs mass, its additional virtues, such as providing a radiative mechanism to explain the origin of the electroweak symmetry breaking and opening up possible ways to unify gravity with other forces (via supergravity and superstrings), have made it especially appealing. Supersymmetry must of course be a broken symmetry in order to agree with observations, and an important unsolved problem of supersymmetric models is the nature and the scale of SUSY-breaking. The most convenient approach is to implement supersymmetry breaking in a hidden sector and then transmit it to the standard-model sector in one of the following two ways: either SUSY-breaking in the hidden sector is conveyed to the observable sector by gravitational interactions; this is the so-called  $N = 1$  supergravity scenario [1], or it is transmitted via the gauge interactions of a distinct messenger sector [2] which contains fields that transform nontrivially under the standard-model gauge group. In this paper we will be concerned with the latter class of models, those with gauge-mediated supersymmetry breaking (GMSB) [3].

The effective low-energy theory that emerges from either of these models contains soft SUSY-breaking mass terms for the scalar superpartners which carry information about the scale and nature of the hidden-sector theory. For instance, typical soft breaking terms for sfermions resulting from the  $N = 1$  supergravity mechanism have magnitude  $\widetilde{m}^2 \sim |F|^2/M_{\text{Pl}}^2$ , where  $|F|$  is the vacuum expectation value (VEV) of the  $F$ -term that breaks supersymmetry in the hidden sector. In order to generate soft masses of the order of  $M_W$  in the matter sector,  $\sqrt{|F|}$  should be around  $10^{11}$  GeV. On the other hand, in the GMSB models where SUSY is broken at the scale  $\Lambda$ , the magnitude of these terms is given by  $\widetilde{m}^2 \sim \frac{\alpha}{4\pi}\Lambda$ ; therefore, the same arguments imply a scale  $\Lambda \lesssim 10^2$  TeV which is much lower. This has the interesting consequence that flavour-changing-neutral-current

(FCNC) processes are naturally suppressed in agreement with experimental bounds. The reason for this suppression is that the gauge interactions induce flavour-symmetric SUSY-breaking terms in the observable sector at  $\Lambda$  and, because this scale is small, only a slight asymmetry is introduced by renormalization group extrapolation to low energies. This is in contrast to the supergravity scenarios where one generically needs to invoke additional flavor symmetries to achieve the same goal.

Another prediction of the GMSB models that distinguishes them from  $N = 1$  supergravity models is the existence of an ultra-light gravitino,  $\tilde{G}$  (which is the Nambu-Goldstone fermion corresponding to spontaneous SUSY breaking), with mass given by  $M_{\tilde{G}} \sim \Lambda^2/M_{\text{Pl}} = 10^{-2}(\Lambda/10 \text{ TeV})^2 \text{ eV}$ . It is therefore the lightest super-particle (LSP). The recent observation by the CDF collaboration of a single event with a final state containing hard  $e^+e^-\gamma\gamma$  and missing transverse energy [4] can indeed be given a straightforward interpretation in the context of GMSB models as selectron pair production in a  $p\bar{p}$  collision with  $\tilde{e} \rightarrow e + \tilde{\gamma}$  followed by  $\tilde{\gamma} \rightarrow \gamma + \tilde{G}$ , and similarly for  $\bar{e}$  [5].

An attractive feature of the GMSB models is that they are highly predictive. Indeed, at energies well below the scale  $\Lambda$ , the theory looks like the usual minimal supersymmetric standard model (MSSM) with the remarkable difference that all the free parameters (about 100) of the low-energy supersymmetric standard model are predicted in terms of three parameters: the SUSY-breaking scale  $\Lambda$ , the  $\mu$ -parameter of the  $H_d H_u$  term in the superpotential, and the soft bilinear mass term,  $B$ . Soft scalar masses  $\tilde{m}$  and gaugino masses  $M$  are induced at the two-loop and one-loop level, respectively, when the messenger sector is integrated out, and their values at the SUSY-breaking scale  $\Lambda$  depend only on  $\Lambda$ . Moreover, the trilinear soft breaking term  $A$  vanishes at  $\Lambda$ . This predictive power has recently been exploited to make a number of testable predictions for the model [6].

In this brief note we make use of the predictive power of the GMSB models to compute

the mass  $M_h$  of the lightest  $CP$ -even state  $h$  present in the Higgs sector as a function of the few parameters of the GMSB models. We use the two-loop corrected Higgs-boson mass spectrum to find accurate upper bounds on the mass of the lightest Higgs boson and, in particular, to extract any possible piece of information on  $M_h$  obtainable by combining these results with indications gathered from sources such as the CDF  $ee\gamma\gamma$  event. In view of planned Higgs searches at LEP2 and LHC [7], we believe that this prediction for  $M_h$  can provide an additional test of the important idea of gauge-mediated supersymmetry breaking.

The minimal GMSB models are defined by three sectors: (i) a secluded sector that breaks supersymmetry; (ii) a messenger sector that serves to communicate the SUSY breaking to the standard model and (iii) the SUSY standard model. The minimal messenger sector consists of a single  $\mathbf{5} + \bar{\mathbf{5}}$  of  $SU(5)$  (to preserve gauge coupling constant unification), *i.e.* color triplets,  $q$  and  $\bar{q}$ , and weak doublets  $\ell$  and  $\bar{\ell}$  with their interactions determined by the following superpotential:

$$W = \lambda_1 S \bar{q} q + \lambda_2 S \bar{\ell} \ell. \quad (1)$$

When the field  $S$  acquires a VEV for both its scalar and auxiliary components,  $\langle S \rangle$  and  $\langle F_S \rangle$  respectively, the spectrum for  $(q, \ell)$  is rendered non-supersymmetric. Integrating out the messenger sector gives rise to gaugino masses at one loop and scalar masses at two loops. For gauginos, we have

$$M_j(\Lambda) = k_j \frac{\alpha_j(\Lambda)}{4\pi} \Lambda, \quad j = 1, 2, 3, \quad (2)$$

where  $\Lambda = \langle F_S \rangle / \langle S \rangle$ ,  $k_1 = 5/3$ ,  $k_2 = k_3 = 1$  and  $\alpha_1 = \alpha / \cos^2 \theta_W$ . For the scalar masses one has

$$\tilde{m}^2(\Lambda) = 2 \sum_{j=1}^3 C_j k_j \left[ \frac{\alpha_j(\Lambda)}{4\pi} \right]^2 \Lambda^2, \quad (3)$$

where  $C_3 = 4/3$  for color triplets,  $C_2 = 3/4$  for weak doublets (and equal to zero otherwise) and  $C_1 = Y^2$  with  $Y = Q - T_3$ . Because the scalar masses are functions of only the gauge quantum numbers, these models automatically solve the supersymmetric flavor problem. Notice the structure of the theory at this level. Squarks are the most massive fields, their masses being roughly a factor of three higher than the slepton masses.

These relations receive significant corrections from the renormalization group evolution (RGE) from the scale  $\Lambda$  down to the weak scale. We have numerically solved the system of one-loop renormalization group equations. Radiative corrections drive the soft breaking mass squared  $m_{H_u}^2$  of the  $H_u$ -doublet, which couples to the top-quark, to negative values near  $M_Z$  leading to electroweak symmetry breaking. They also raise slightly the soft breaking mass squared for the sleptons. After including the effects of the RGE and  $D$ -terms, the experimental limits on the right-handed selectron mass requires

$$\Lambda \gtrsim 10 \text{ TeV}. \quad (4)$$

We notice here that, if the  $e^+e^-\gamma\gamma$  plus missing-transverse-energy event originates from slepton pair-production (*e.g.*  $\tilde{e}_L\bar{\tilde{e}}_L$  or  $\tilde{e}_R\bar{\tilde{e}}_R$ ), this restricts the values of slepton masses to  $(130 \gtrsim m_{\tilde{e}_{L,R}} \gtrsim 80) \text{ GeV}$  [5]. The  $\tilde{e}_R$ -mode in turn implies that

$$(30 \lesssim \Lambda \lesssim 50) \text{ TeV}, \quad (5)$$

whereas the  $\tilde{e}_L$ -mode implies  $20 \lesssim \Lambda \lesssim 35 \text{ TeV}$ . These upper bounds on  $\Lambda$  will be used in the following to constrain the mass of the lightest Higgs boson from above.

It is important to point out that the magnitude of the  $\mu$ - and  $B$ -parameters at the scale  $\Lambda$  depends crucially on the structure of the Higgs sector. In the *minimal* messenger model, which contains only the usual two Higgs doublets, one expects the  $B$ -parameter to be small at the scale  $\Lambda$  and to evolve to significant values at the scale  $M_Z$  in the process of running. In general, in order to generate the parameters  $\mu$  and  $B$  at the scale

$\Lambda$ , the Higgs sector should be enlarged [8],[9]. However, this is not expected to affect the results of this paper since, in general, the extra Higgs fields are so heavy that they decouple from the matter fields at low energy.

Let us now consider the low-energy spectrum of the GMSB models as far as the Higgs sector is concerned. As just mentioned, we assume that its particle content at low energies is exactly that of the MSSM. However, there are additional restrictions coming from the structure of the GMSB theories. The one-loop effective Higgs potential may be expressed as the sum of the tree-level potential plus a correction coming from the sum of one-loop diagrams with external lines having zero momenta,

$$V_{1\text{-loop}} = V_{\text{tree}} + \Delta V_1. \quad (6)$$

The right-hand side is independent of the running scale  $Q$  to one-loop order. The one-loop correction is given by (in the  $\overline{DR}$ -scheme)

$$\Delta V = \frac{1}{64\pi^2} \sum_j (-1)^{2s_j} (2s_j + 1) m_j^4 \left( \ln \frac{m_j^2}{Q^2} - \frac{3}{2} \right), \quad (7)$$

where  $m_j$  is the eigenvalue mass of the  $j^{\text{th}}$  particle with spin  $s_j$  in the  $(v_d, v_u)$  background, with  $v_d = \langle H_d^0 \rangle$  and  $v_u = \langle H_u^0 \rangle$ . The tree-level part of the potential of the MSSM Higgs sector reads

$$\begin{aligned} V_{\text{tree}} = & m_d^2 |H_d|^2 + m_u^2 |H_u|^2 - (m_3^2 H_d H_u + \text{h.c.}) \\ & + \lambda_1 |H_d|^4 + \lambda_2 |H_u|^4 + \lambda_3 |H_d|^2 |H_u|^2 + \lambda_4 |H_d H_u|^2. \end{aligned} \quad (8)$$

Here

$$\begin{aligned} \lambda_1 = \lambda_2 = & \frac{g_1^2 + g_2^2}{8}, \\ \lambda_3 = & \frac{g_2^2 - g_1^2}{4}, \\ \lambda_4 = & -\frac{g_2^2}{2}, \end{aligned} \quad (9)$$

where  $g_1$  and  $g_2$  are the gauge couplings of the  $U(1)_Y$  and  $SU(2)_L$  gauge groups respectively, and

$$m_d^2 = m_{H_d}^2 + |\mu|^2, \quad m_u^2 = m_{H_u}^2 + |\mu|^2, \quad m_3^2 = B\mu. \quad (10)$$

The parameters of the potential are allowed to run; that is, they vary with scale according to the RGE. We must use the RGE to evolve the parameters of the potential to a convenient scale such as  $M_Z$  (where the experimental values of the gauge couplings are determined). After the following redefinition

$$\bar{m}_i^2 = m_i^2 + \frac{\partial \Delta V}{\partial (v_i^2)}, \quad i = d, u, \quad (11)$$

minimization of the potential yields the following conditions among the parameters:

$$\frac{1}{2}M_Z^2 = \frac{\bar{m}_d^2 - \bar{m}_u^2 \tan^2 \beta}{\tan^2 \beta - 1}, \quad (12)$$

$$B\mu = -\frac{1}{2}(\bar{m}_d^2 + \bar{m}_u^2) \sin 2\beta, \quad (13)$$

where  $\tan \beta = v_u/v_d$ .

After  $M_Z^2$  has been fixed to its physical value, all masses may be expressed in terms of only two parameters and we have chosen them to be the SUSY-breaking scale  $\Lambda$  and  $\tan \beta$ . The  $\mu$ -parameter at the scale  $M_Z$  is then fixed by Eq. (12).

The minimization conditions lead to the determination of the the tree-level mass  $M_h^{\text{tree}} = M_Z |\cos 2\beta|$  of the lightest  $CP$ -even state  $h$  of the Higgs spectrum. However, it is well-known that radiative corrections contribute significantly to the physical mass  $M_h$ . The Higgs-boson mass was first determined by the renormalization-group resummation of all-loop leading log (LL) corrections in [10]. Some next-to-leading log (NTLL) corrections were further introduced in [11] and [12], and finally a complete NTLL analysis was performed in [13] and [14]. One of the main issues in [14] was the comparison between

the LL and the NTLL approximations. As expected, the LL approximation shows a strong scale-dependence, while the NTLL is almost scale-independent. This implies not only that, working in the NTLL approximation, the choice of scale is almost irrelevant, but also that the LL approximation may yield accurate results if a correct choice of the renormalization scale is made. The scale where both results coincide turns out to be close to the pole top-quark mass  $M_t$  [14].

Very useful analytical approximations to the numerical all-loop renormalization-group improved LL result, including two-loop leading-log effects, may be found in [15] where the reader is referred to for more details. We report here the expression for  $M_h$  only in the case in which the mass  $M_A$  of the  $CP$ -odd state in the Higgs spectrum is much larger than  $M_Z$ <sup>†</sup>:

$$M_h^2 = M_Z^2 \cos^2 2\beta \left( 1 - \frac{3}{8\pi^2} \frac{m_t^2}{v^2} t \right) + \frac{3}{4\pi^2} \frac{m_t^4}{v^2} \left[ \frac{1}{2} \widetilde{X}_t + t + \frac{1}{16\pi^2} \left( \frac{3}{2} \frac{m_t^2}{v^2} - 32\pi\alpha_3 \right) (\widetilde{X}_t t + t^2) \right], \quad (14)$$

where  $v^2 = v_d^2 + v_u^2$ ,

$$t = \ln \left( \frac{M_S^2}{M_t^2} \right), \quad (15)$$

$$m_t = \frac{M_t}{1 + \frac{4}{3\pi} \alpha_3(M_t)} \quad (16)$$

is the on-shell running mass and  $\alpha_3$  indicates

$$\alpha_3(M_t) = \frac{\alpha_3(M_Z)}{1 + \frac{b_3}{4\pi} \alpha_3(M_Z) \ln(M_t^2/M_Z^2)}, \quad (17)$$

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<sup>†</sup>In this case all degrees of freedom except the lightest  $CP$ -even state decouple, leaving an effective theory which is similar to the standard model with different boundary conditions for the Higgs quartic coupling. In the opposite case  $M_A \lesssim M_Z$ ,  $M_h$  depends on  $M_A$ , see [15].

where  $b_3$  is the one-loop QCD beta function. Moreover,  $\widetilde{X}_t$  is the stop mixing parameter

$$\begin{aligned}\widetilde{X}_t &= \frac{2\widetilde{A}_t^2}{M_S^2} \left(1 - \frac{\widetilde{A}_t^2}{12 M_S^2}\right), \\ \bar{A}_t &= A_t - \mu \cot\beta.\end{aligned}\tag{18}$$

The expressions above assume that only the squarks of the third generation contribute to the radiative corrections (this translates into the bound  $\tan\beta \lesssim 35$ ).

The scale  $M_S$  is to be associated with the characteristic stop mass scale and we have computed it in the following way. We have solved the RGE for the soft SUSY-breaking parameters which enter the stop mass matrix. They are  $\widetilde{A}_t$ ,  $\widetilde{m}_Q$  and  $\widetilde{m}_U$ , where the latter are the soft SUSY-breaking mass terms of the left-handed and right-handed stop, respectively. The initial conditions at the scale  $\Lambda$  are given by Eq. (3) and by

$$A_t(\Lambda) = 0.\tag{19}$$

Defining the stop squared-mass eigenvalues by  $M_{\widetilde{t}_1}^2$  and  $M_{\widetilde{t}_2}^2$ , the scale  $M_S$  has been defined as the scale at which<sup>†</sup>

$$M_{\widetilde{t}_1}^2(M_S)M_{\widetilde{t}_2}^2(M_S) = M_S^2.\tag{20}$$

Other operative definitions are possible, for example  $M_S^2 = (M_{\widetilde{t}_1}^2(M_S) + M_{\widetilde{t}_2}^2(M_S))/2$ , but these different distinctions have no significant impact on the final result for  $M_h$ . We have generally found the  $\mu$ -parameter to be so large that the pseudoscalar mass  $M_A$  is driven to values much larger than  $M_Z$ , rendering the expression for  $M_h$  in Eq. (14) very reliable.

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<sup>†</sup>Since, strictly speaking, the operator expansion leading to the expression (14) is performed in the symmetric phase, one should have used the product of the SUSY-breaking squared masses  $\widetilde{m}_Q(M_S)\widetilde{m}_U(M_S)$  to define the scale  $M_S$ . We have checked that the numerical shift in the final result for  $M_h$  is negligible when adopting this definition instead of the one in Eq. (20).

Notice that the expression (14), which we made use of in the case  $M_A \gg M_Z$ , is only valid under the assumption (see Refs. [15] for a thorough discussion)

$$\frac{M_{\tilde{t}_1}^2(M_S) - M_{\tilde{t}_2}^2(M_S)}{M_{\tilde{t}_1}^2(M_S) + M_{\tilde{t}_2}^2(M_S)} \lesssim 0.5. \quad (21)$$

We have checked numerically that this condition was satisfied.

In Fig. 1 and Fig. 2, we present our predictions for the mass of the lightest Higgs boson in the GMSB models as a function of the scale  $\Lambda$  for different values of  $M_t$  and  $\tan\beta$ . From Fig. 1, we see that the values of  $M_h$  for a top quark mass of 175 GeV range from 85 to 110 GeV for  $\Lambda = 50$  TeV. Fig. 2 shows the  $\Lambda$  dependence of  $M_h$  for  $M_t = 175$  GeV. The requirement that the CDF  $ee\gamma\gamma$  event is explained by the GMSB scenario constrains the values of  $M_h$  to lie on the left-hand side of the vertical lines which show the upper bounds on  $\Lambda$  coming from the constraint  $m_{\tilde{e}_R} \lesssim 130$  GeV (long-dashed line) and  $m_{\tilde{e}_L} \lesssim 130$  GeV (dashed line). We infer that  $M_h \lesssim 110$  GeV if  $\tilde{e}_R$  leads to the CDF event and  $M_h \lesssim 105$  GeV for the  $\tilde{e}_L$  case. Interestingly enough, this mass range is accessible at LEP2 with a center-of-mass energy  $\sqrt{s} = 205$  GeV. This opens the exciting possibility that one can obtain useful information about the GMSB models once these ranges of Higgs masses are explored at LEP2 and LHC.

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## Figure Captions

**Fig. 1:** The physical mass of the lightest  $CP$ -even state  $M_h$  as a function of  $M_t$  for  $\tan\beta = 2, 15$  and  $\Lambda = 50$  TeV.

**Fig. 2:** The physical mass of the lightest  $CP$ -even state  $M_h$  as a function of the SUSY-breaking scale  $\Lambda$  for  $\tan\beta = 2, 15$  and  $M_t = 175$  GeV. The dashed and the long-dashed vertical lines indicate the kinematical upper bounds on  $\Lambda$  from the interpretation of the CDF  $ee\gamma\gamma$  event as originating from  $\tilde{e}_L\tilde{e}_L$  and  $\tilde{e}_R\tilde{e}_R$  production, respectively.



