Naturalness constraints on gauge-mediated supersymmetry breaking models

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Abstract

The question of naturalness is addressed in the context of gauge-mediated supersymmetry breaking models. Requiring that $M_Z$ arises naturally imposes upper limits on the right-handed selectron mass in these models that are stronger than in the Minimal Supersymmetric Standard Model and are interesting from the point of view of searches at the current and future colliders.

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How natural it is to conceive large superpartner scalar masses keeping the weak scale light? In this Letter we address this question of naturalness in the context of gauge-mediated supersymmetry breaking (GMSB) models \(^1\). There is a naturalness problem in the Standard Model itself, following from the assumption of the Higgs boson as a fundamental object, which is best cured by the elegant introduction of supersymmetry, providing a cut-off for the quadratic divergences associated with the Higgs. Since supersymmetry has the virtue of providing a radiative mechanism for spontaneous breakdown of electroweak symmetry, the superpartner scalar masses cannot be arbitrarily heavy, if we require that the weak scale arises naturally. In other words, as those scalar masses are pushed higher, the weak scale still remains light only at the price of an increasingly delicate cancellation among the soft supersymmetry breaking masses in the theory. Thus a quantitative requirement of naturalness puts an upper limit on superparticle masses. In the Minimal Supersymmetric Standard Model with universal boundary conditions, this requirement leads to a \(\sim 10^2\) TeV scale upper bound on the superpartner masses \(^3\). We find that in GMSB models, the naturalness upper limits are comparatively stronger.

In the GMSB scenario there exists a set of heavy chiral superfields, called the “messenger” fields. Supersymmetry is broken in the messenger sector due to the interaction of the messengers with a GUT-singlet spurion field which has a non-vanishing vacuum expectation value (vev) in its auxiliary component \((F)\). The information of supersymmetry breaking is then transmitted to the superpartners of the quarks, leptons, gauge and Higgs bosons via the usual SU(3)\(\times\)SU(2)\(\times\)U(1) gauge interactions. The gauge-mediation attributes these models the virtue of automatically suppressing the Flavor-Changing Neutral Currents. In the supergravity scenario, where supersymmetry breaking in the hidden sector is communicated to the visible sector by gravitational strength interactions, the breaking scale \(\sqrt{F}\) is \(\sim 10^{10}\) GeV in order that the scalar masses obtained from \(\tilde{m}^2 \sim F^2/M_{Pl}^2\) are in the TeV range. On the other hand, in the GMSB models, where the scalar masses are generated radiatively, the analogous expressions of the squared masses have a loop-suppression factor of \((\alpha/4\pi)^2\) and \(M_{Pl}\) is replaced by the messenger mass \((10\text{ TeV} \lesssim M \lesssim M_{Pl})\). Depending upon the messenger scale, this can lead to a smaller intrinsic supersymmetry breaking scale and consequently to a superlight gravitino \(\tilde{G}\) constituting the lightest supersymmetric particle. Of late, there has been a resurgence of interests \(^6\) in these models as these massless gravitinos help to provide a comfortable explanation to the recent CDF event with \(e^+e^-\gamma\gamma + \text{missing energy}\) in the final state \(^7\). Very recently, detailed phenomenological analyses of a general class of GMSB models have been presented in refs. \(^8\) and \(^9\).

The simplest version of the messenger models involves a set of vector-like superfields \((\tilde{M}_i + \bar{M}_i)\). These couple to a gauge singlet superfield \(X\) through the following superpotential,\(^{1}\)

\[
W = \lambda_i XM_i\bar{M}_i.
\]  

\(M_i\) and \(\bar{M}_i\) transform under a complete SU(5) representation and there could be \(n_5\) copies of \((5 + \bar{5})\) and \(n_{10}\) copies of \((10 + \bar{10})\) messengers. The quantity \(n = (n_5 + 3n_{10})\) could at most be four from the requirement of perturbative unification. The superfield \(X\) must acquire non-zero vevs both in its scalar and auxiliary components through its interaction

\(^1\)This issue has been reinvestigated with a somewhat modified criterion of naturalness in ref. \(^4\) and with nonuniversal boundary conditions in ref. \(^5\).
with the hidden sector fields. Assuming for simplicity $\lambda \equiv \lambda_i$ (for each $i$), it follows from the superpotential written above that the messenger fermions acquire a supersymmetric mass $M \equiv \lambda \langle X \rangle$ and the messenger scalars are split as having masses $M_{\pm}^2 = M^2 \pm \lambda \langle F_X \rangle$. The effective supersymmetry breaking scale is $\Lambda \equiv \langle F_X \rangle / \langle X \rangle$. Integrating out heavy messenger fields generate gaugino masses at one-loop at the $M$-scale, as

$$\tilde{M}_i(M) = n \, R \, \tilde{\alpha}_i(M) \, \Lambda \, g(\Lambda/M),$$

where $\tilde{\alpha}_i \equiv \alpha_i / 4\pi$ correspond to the gauge couplings and $g(x)$ is the loop function. $R \leq 1$ arises from the fact the $U(1)_R$-symmetry could break at a scale smaller than $M$ rendering the gauginos lighter. The scalar masses arise at the two-loop level and their expressions at the generating scale $M$ are given by

$$\tilde{m}(M) = 2 \, n \, \Lambda^2 \, f(\Lambda/M) \sum_{i=1}^{3} C_i \, \tilde{\alpha}_i^2(M),$$

where $C_i = 4/3, 3/4$ for the fundamental representations of $SU(3)$, $SU(2)$ respectively and zero for singlets, $C_i = 3Y^2/5$ for $U(1)$ ($Y \equiv Q - T_3$) and $f(x)$ is the associated loop function. The functions $g(x)$ and $f(x)$ [10] are very close to unity in the limit $x \to 0$ and, therefore, the expressions of the gaugino and scalar masses become particularly simple when $\Lambda \ll M$, which is a reasonable limit that we rely upon in the rest of the paper. Also, in this limit the gaugino and scalar masses are independent of the Yukawa couplings appearing in the superpotential. We do not treat $n$ as an independent parameter, since it can always be absorbed in a redefinition of $R$ and $\Lambda$ as $\sqrt{n} \, R \rightarrow R$ and $\sqrt{n} \, \Lambda \rightarrow \Lambda$. Actually, $g$ and $f$ can also be absorbed in a further redefinition of $R$ and $\Lambda$.

As has been pointed out in ref. [11], there is an intrinsic “$\mu$-problem” associated with the $\mu \tilde{H}_u \tilde{H}_d$ term in the superpotential and the $B\mu \tilde{H}_u \tilde{H}_d$ term in the scalar potential, since these Peccei-Quinn symmetry-violating terms cannot be generated by gauge interactions. Both are generated at the same loop level in generic models rendering either $\mu$ to be at the weak scale and $B$ too large beyond the naturalness bound, or $B$ at the weak scale and $\mu$ too small to be phenomenologically acceptable. The problem can be cured by arranging the model in such a way that $\mu$ is generated at one-loop but $B\mu$ at two-loop [11]. This keeps both $\mu$ and $B$ at the weak scale but indeed at the price of introducing additional interactions which could affect the Higgs boson soft mass terms. We parametrize this effect by adding two quantities $\Box_u^2$ and $\Box_d^2$ at the scale $M$ in the following way:

$$m_{\tilde{H}_u(\tilde{H}_d)}^2(M) = m_{\tilde{L}}^2(M) + \Box_u^2(\Box_d^2).$$

Soft trilinear $A$-terms are zero at $M$. Their radiative generation requires the simultaneous violation of the $U(1)_R$ symmetry and the chiral flavor symmetry in the observable sector in the broken supersymmetry phase. Since the messengers do not violate observable sector

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$^2$We denote henceforth the superfields and their scalar components by the same symbols.

$^3$Our $\alpha_1$ always corresponds to the one which unifies with $\alpha_2$ and $\alpha_3$, i.e. it is related to the Standard Model $\alpha'_1$ by $\alpha_1 = 3\alpha'_1/5$.

$^4$The limit $\Lambda = M$ leads to massless messenger scalars which are definitely unwanted. Also $\Lambda > M$ leads to unphysical messenger scalar masses.
chiral flavor, the $A$-terms are not generated at one-loop. They are generated only at the two-loop level requiring a gaugino mass insertion in the internal line [8].

The GMSB models that we consider are characterized by the following parameters: $M$, $\Lambda$, $R$, $\square_u^2$, $\square_d^2$, $\mu$ and $B_0$ (the value of $B$ at $M$). They determine the electroweak scale $M_Z$ and $\sin 2\beta$ through

\begin{align}
M_Z^2 &= 2 \frac{m_{H_u}^2 - m_{H_d}^2 \tan^2 \beta}{\tan^2 \beta - 1} - 2\mu^2, \quad (5) \\
\sin 2\beta &= \frac{2B\mu}{m_{H_u}^2 + m_{H_d}^2 + 2\mu^2}; \quad (6)
\end{align}

where we have used tree-level Higgs potentials. We compute $m_{H_u(d)}^2$ using the boundary conditions in eqs. (2), (3) and (4) and running down from $M$ to the stop mass $m_{\tilde{t}}$ by one-loop renormalization group (RG) equations\

\[5.\] This is enough for the purpose of studying naturalness. We use $\alpha_3(M_Z) = 0.117$ and the top and bottom quark running masses at the top mass scale as 165 and 2.8 GeV respectively.

Now we are all set to discuss the fine-tuning criteria. In general, if a quantity $X = X(a_1, a_2, \ldots)$ is determined as a function of many input parameters $a_i$, then a measure of how much the parameters $a_i$ are fine-tuned in the determination of the quantity $X$ can be expressed through the parameter

\[\Delta X_{a_i} = \left| \frac{a_i \partial X}{X \partial a_i} \right|, \quad (7)\]

for each $a_i$. $\Delta$ measures the relative variation of $X$ against the relative variation of $a_i$. In our case, a large fine-tuning (as defined in eq. (7)) corresponds to large cancellations among independent terms.

Let us first examine the amount of fine-tuning that goes in the determination of $\sin 2\beta$ through eq. (6). For large $\tan \beta$, $\sin 2\beta$ is small and a cancellation is possible only if $B = B_0 + B_{\text{RG}}$ (where $B_{\text{RG}}$ is the running effect of $B$ and that does not depend on $B_0$) is much smaller than $B_0$. The quantity $|B_{\text{RG}}/B|$, which is in fact the same as $\Delta_{B_0}^{\sin 2\beta}$, becomes larger if $M$ increases (because in this case $B_{\text{RG}}$ increases) or if $\tan \beta$ becomes larger (because in this case $B \propto \sin 2\beta$ decreases). For example, for $\tan \beta \simeq 50$, $M = 10^5 \text{TeV}$, $\Delta_{B_0}^{\sin 2\beta} \simeq 10$.

Now let us concentrate on the determination of $M_Z$ through eq. (5). $M_Z$ can be expressed as a function of two sets of parameters: $\Lambda$, $M$ and $R$ in one set and $\mu$, $\square_u^2$ and $\square_d^2$, which depend on unknown hidden sector couplings, in another. Neglecting the running of $\square_u(d)$ (which are small), one can rewrite eq. (5) as

\[M_Z^2 = 2 \frac{\tilde{m}_{H_d}^2 - \tilde{m}_{H_u}^2 \tan^2 \beta}{\tan^2 \beta - 1} - 2\tilde{\mu}^2, \quad (8)\]

where, $\tilde{m}_{H_u(d)}^2 = m_{H_u(d)}^2 (\square_u^2 = \square_d^2 = 0)$ and

\[\tilde{\mu}^2 = \mu^2 - \frac{\square_d^2 - \square_u^2 \tan^2 \beta}{\tan^2 \beta - 1}. \quad (9)\]
From eq. (9), it is clear that all the hidden sector couplings could effectively be expressed just in terms of one parameter through a redefinition of $\mu^2 \rightarrow \tilde{\mu}^2$ and, therefore, there is no loss of generality even if we set $\Box^2_t = \Box^2_d = 0$ in our subsequent analysis.

Fine-tuning, in our case, corresponds to large cancellations between the two *uncorrelated* quantities in the right hand side (RHS) of eq. (8). A quantitative measure of the cancellation between these two terms is given by

$$\Delta = \frac{2\mu^2}{M_Z^2} = \left| 2 \frac{m_{H_d}^2 - m_{H_u}^2 \tan^2 \beta}{(\tan^2 \beta - 1)M_Z^2} - 1 \right| = \Delta_{\mu M_Z}^2. \tag{10}$$

Owing to our ignorance of the hidden sector couplings, $\Delta = \Delta_{\mu M_Z}^2$ is the best available measure of fine-tuning in this set-up. The rest of our paper concerns a qualitative and a quantitative analysis of $\Delta$.

First, we observe that for large $\tan \beta$, $\Delta \simeq 2|m_{H_u}^2/M_Z^2|$. In this case, the squared top Yukawa coupling ($Y_t$) is insensitive to $\tan \beta$, while the bottom- and tau-Yukawa couplings do not contribute to $m_{H_u}^2$ renormalization. Therefore, for fixed $\Lambda$ and for $\tan \beta \geq 6$, $\Delta$ does not vary with $\tan \beta$. Actually, for large $\tan \beta$, $m_{\tilde{t}_R} \lesssim m_{\tilde{u}_R}$ owing to nonnegligible $Y_t$-radiative corrections, so that, if we fix $m_{\tilde{t}_R}$ instead of $\Lambda$, a small dependence of $\Delta$ on $\tan \beta$ creeps in. In any case, we hereafter confine ourselves to moderate $\tan \beta$ region.

Apart from $\tan \beta$, $\Delta$ depends also on $M$, $R$ and $\Lambda$ (or $m_{\tilde{t}_R}$). Before demonstrating their exact quantitative effects, let us examine their qualitative dependences by varying them one at a time while keeping the others fixed. From eqs. (2) and (3) and the RG evolution it follows:

$$m_{H_d}^2 \simeq \frac{3}{2} \left[ \tilde{\alpha}_2^2(M) + R^2 \left( \tilde{\alpha}_2^2(M) - \tilde{\alpha}_2^2(m_t) \right) \right] \Lambda^2, \tag{11}$$

$$m_{H_u}^2 \simeq m_{H_d}^2 - 3 \int_0^{t(m_t)} dt \tilde{Y}_t (m_{\tilde{Q}_3}^2 + m_{\tilde{t}_R}^2 + m_{H_u}^2), \tag{12}$$

where, $t(Q) = 2 \ln(M/Q)$ and $\tilde{Y}_t = Y_t/4\pi$. In the RHS of eq. (12) the second term is roughly proportional to $\alpha_3^2$, and, neglecting the Yukawa radiative corrections (just for qualitative understanding), the dependences on $\alpha_3^2$ itself turn out to be

$$m_{\tilde{Q}_3}^2 \simeq \left[ \frac{8}{3} \tilde{\alpha}_3^2(M) + \frac{8}{9} R^2 \left( \tilde{\alpha}_3^2(m_t) - \tilde{\alpha}_3^2(M) \right) \right] \Lambda^2. \tag{13}$$

From eq. (10) it follows,

$$\Delta = \left( \frac{\Lambda}{M_Z} \right)^2 H(M, R, \tan \beta) - 1, \tag{14}$$

where, actually, $H$ has a mild logarithmic dependence on $\Lambda$ through $m_{\tilde{t}}$, as is apparent from

$$H \approx \frac{6 \tan^2 \beta}{\tan^2 \beta - 1} \int_0^{t(m_t)} dt \tilde{Y}_t \left[ \frac{16}{9} \tilde{\alpha}_3^2(M) + \frac{16}{9} R^2 \left( \tilde{\alpha}_3^2 - \tilde{\alpha}_3^2(M) \right) \right]$$

$$- 3 \left[ \tilde{\alpha}_2^2(M) + R^2 \left( \tilde{\alpha}_2^2(M) - \tilde{\alpha}_2^2(m_t) \right) \right]. \tag{15}$$

Owing to the $\alpha_3^2$-dependence, the first term on the RHS of eq. (15) is the leading one (also, this term is responsible for electroweak symmetry breaking) and, consequently, $H$ is positive.
The prefactor \( (\Lambda/M_Z)^2 \) in eq. (14) can be expressed in terms of the right-handed selectron mass as

\[
\left( \frac{\Lambda}{M_Z} \right)^2 \simeq \frac{5}{6\alpha_1^2(M)} \left[ \frac{m_{\tilde{e}_R}^2}{M_Z^2} - \sin^2 \theta_W \frac{\tan^2 \beta - 1}{\tan^2 \beta + 1} \right],
\]

where, for illustration, we do not explicitly show the small U(1) RG correction. In the above equation, the term containing \( \sin^2 \theta_W \) corresponds to the tree-level U(1) D-term contribution. For given values of \( \tan \beta, M \) and \( R, H \) is determined and this could be translated into an upper bound of right-handed selectron mass. For example, for \( \tan \beta = 6, M = 100 \) TeV and \( R = 1, m_{\tilde{e}_R} < 83 \) (280) GeV for \( \Delta < 10 \) (100) (these numbers follow from exact numerical analysis). We emphasize that for a light selectron, the tree-level D-term relaxes the constraint. We also notice that the experimental lower limit on \( m_{\tilde{e}_R} > 45 \) (70) GeV from LEP1 (LEP1.5) can be translated through eq. (14) to a lower limit of \( \Lambda \) depending on \( M \); for example, for \( M \sim \Lambda, \Lambda \gtrsim 10 \) (30) TeV and, on the other extreme, for \( M \sim M_{\text{GUT}}, \Lambda \gtrsim 4 \) (15) TeV.

In the moderate \( \tan \beta \) region and for given values of \( M, m_{\tilde{e}_R} \) and \( R \), a lower \( \tan \beta \) increases \( \Delta \) implying a larger fine-tuning and hence a stronger constraint. Decreasing \( \tan \beta \) increases \( \Delta \) because the leading term in \( H \) contains a factor \( \tan^2 \beta/(\tan^2 \beta - 1), Y_t(m_t) \propto (\tan^2 \beta + 1)/\tan^2 \beta \) and finally for a given \( m_{\tilde{e}_R}, (\Lambda/M_Z)^2 \) increases (see eq. (13)).

Other parameters remaining fixed, if the U(1)\(_R\) symmetry breaks at a scale lighter than \( M, R < 1 \) in the minimal GMSB model (\( n_5 = 1, n_{10} = 0 \)), implying a lower \( \Delta \) and hence a weaker constraint. If U(1)\(_R\) breaks at \( M \), then having more messengers increases \( R \) (from 1 to at most 2 remaining consistent with the requirement of perturbative unification) and the constraint becomes stronger.

The dependence of \( \Delta \) on \( M \) is not apparent unlike in the cases of other parameters owing to several counteracting contributions. \( \Delta \) depends in fact on \( M \) through the range of integration, the boundary conditions on gaugino and scalar masses and the \( \alpha_1^{-2}(M) \) dependence of \( \Delta \) in eq. (14) in conjunction with eq. (10). It turns out that \( \Delta \) decreases when \( M \rightarrow M_{\text{GUT}} \).

Now we concentrate on exact numerical analysis. The dependences of \( \Delta \) on \( M \) and \( R \) are plotted in Fig. 1a while in Fig. 1b we exhibit the limits on \( m_{\tilde{e}_R} \) as functions of \( \tan \beta \) for conservative choices of \( M \) and \( R \) extracted from Fig. 1a. Combining eqs. (14) and (16) yields

\[
\Delta \simeq \left( \frac{m_{\tilde{e}_R}^2}{M_Z^2} - \sin^2 \theta_W \frac{\tan^2 \beta - 1}{\tan^2 \beta + 1} \right) \frac{5}{6\alpha_1^2(M)} H(M, R, \tan \beta) - 1
\]

(17)

For studying \( R \) and \( M \) dependences, we fix \( \tan \beta = 6 \). First we choose a representative value \( m_{\tilde{e}_R} = 100 \) GeV and denote the fine-tuning parameter as \( \Delta_{100} \). Then we plot isocurves of \( \Delta_{100} (= 3, 10, 20) \) in the \( \log_{10}(M/\text{GeV})-R \) plane. The fine-tuning for an arbitrary value of \( m_{\tilde{e}_R} \) can be read from Fig. 1a, using

\[
\Delta \simeq \left( \frac{m_{\tilde{e}_R}^2}{M_Z^2} - 0.21 \right) \Delta_{100},
\]

(18)

which is a good approximation unless \( m_{\tilde{e}_R} \) is in a physically uninteresting region, being much heavier than 100 GeV. It is apparent from Fig. 1a that in the region \( M \sim M_{\text{GUT}} \) a right-handed selectron as heavy as 100 GeV does not suffer from fine-tuning problems. Also lowering \( R \) reduces the fine-tuning.
Figure 1: (a): Contour plots of $\Delta_{100}$ (the fine-tuning parameter for $m_{\tilde{e}_R} = 100$ GeV) in the $\log_{10}(M/{\text{GeV}})$–$R$ plane for $\tan \beta = 6$; (b) naturalness upper limits on the right-handed selectron mass at fixed $M = 10^5$ TeV and $R = 1$.

In Fig. 1b, we show the naturalness upper limit on $m_{\tilde{e}_R}$ as a function of $\tan \beta$ in the moderate $\tan \beta$ regime for $M = 10^5$ TeV (which corresponds to $m_{\tilde{G}} \sim 100$ eV for $\Lambda = 10$ TeV) and $R = 1$.

To conclude, we have studied the naturalness constraints in the GMSB models. The requirement of a natural determination of $M_Z$ in terms of the parameters of the model yields rather strong upper limits on the mass of the right-handed selectron. For example, for $M = 10^5$ TeV, $R = 1$ and $\tan \beta = 6$, $m_{\tilde{e}_R} < 80$ (230) GeV for $\Delta = 10$ (100). If we choose $\Delta < 10$ (i.e. a fine-tuning not larger than one order of magnitude), the upper limits are quite interesting for LEP2. We note in passing that an explanation of the CDF $e^+e^-\gamma\gamma$ + missing energy event requires a cancellation of one order of magnitude in a large range of $M$. The quoted upper limits are more stringent than in the Minimal Supersymmetric Standard Model. In fact, in the latter scenario, the universal boundary conditions on the scalar masses do not permit the lightest scalar to be much lighter than the (up-type) Higgs mass. On the other hand, in the GMSB case, owing to the proportionality of the scalar masses to different gauge couplings (squared) at $M$, the lightest scalar becomes light enough compared to the Higgs mass to pose a strong fine-tuning constraint.

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After completing our work, we have become aware of a related work in ref. [12].

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6 As observed before, $\Delta$ is flat with respect to $\tan \beta$ for $\tan \beta > 6$. 

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