Bottom–up Approach in Supersymmetric Models*

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Abstract

We present a bottom-up approach to the question of supersymmetry breaking in the MSSM. Starting with the experimentally measurable low energy supersymmetry breaking parameters which can take any values consistent with present experimental constraints, we evolve them up to an arbitrary high energy scale. Approximate analytical expressions for such an evolution, valid for low and moderate values of tan β, are presented. We then discuss qualitative properties of the high energy parameter space and in particular, identify the conditions on the low energy spectrum which are necessary for the parameters at high energy scale to satisfy simple regular pattern such as universality or partial universality.

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1 Introduction

In supersymmetric theories the mechanism of supersymmetry breaking remains a fundamental open problem. Its low energy manifestation is the supersymmetric spectrum. Therefore, one may hope to get some insight into this problem from experiment. Currently, the most popular view on supersymmetry breaking is that the parameters of the low energy effective theory have their origin in the GUT (or string) scale physics [1]. Supersymmetry is spontaneously broken in an invisible sector and this effect is transferred to our sector through supergravity interactions at the scale $M_{Pl}$. Other models have also been proposed [2] in which the supersymmetry breaking is communicated to the electroweak sector through some other messengers at the energy scale $M \ll M_{Pl}$. In the Minimal Supersymmetric Standard Model (MSSM) the soft supersymmetry breaking parameters at the large scale are connected to their low energy values via the renormalization group equations (RGE) which do not contain any new unknown parameters. With any simple theoretical Ansatz for the pattern of soft supersymmetry breaking terms at the scale where supersymmetry breaking is transmitted to the observable sector one can study superpartner spectra in the top-down approach. It is clear that any particular Ansatz for parameters at high energy scale gives only a small subspace of all low energy parameter space. In order to have a broad overview of the low energy - high energy parameter mapping it is of interest to supplement the top-down approach with a bottom-up one where we can learn how certain qualitative features of the low energy spectra reflect themselves on the qualitative pattern of soft terms at different energy scales. Of course, eventual measurement of the superpartner and Higgs boson spectra and various mixing angles in the sfermion sector will permit (in the framework of MSSM) a complete bottom-up mapping. We shall know then the pattern of the soft parameters at any hypothetical scale $M$ of supersymmetry breaking and this will have major impact on our ideas on its origin.

This talk is based on results obtained in reference [3]. We present here the main features of the bottom-up mapping for the set of parameters $\mu$, $M_2$, $m_{H_1}^2$, $m_{H_2}^2$, $B$ and the third generation squark mass parameters $m_Q^2$, $m_U^2$, $m_D^2$ and $A_t$. To a very good approximation this is a closed set of parameters, whose RG running decouples from the remaining parameters [4]. We shall concentrate on the region of small to moderate values of $\tan \beta$ in which, for $m_t = 175 \pm 6$ GeV, the bottom quark Yukawa coupling effects may be neglected. For this case we present analytic expressions (at one loop level) for the values of our set at $M_{GUT}$ and at any $M < M_{GUT}$ in terms of its values at $M_Z$ scale. The equations are valid for arbitrary boundary values of the parameters at the scale $M$.

In absence of direct experimental measurement of the low energy parameters, we disscuss the mapping of the region characterized by light chargino and right handed stop with masses of the order of $M_Z$ which is consistent with all existing experimental constraints and is of interest for LEP2.

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1. In the small or moderate $\tan \beta$ regime, the dependence of the RG running of these parameters on slepton (and the two first generation sfermion) masses comes only through the small hypercharge D-term contributions which we neglect here. Their contribution can be found in [3].
The bottom–up mapping is presented in section 2. Section 3 contains some general results obtained with this method and its application to the low energy region with light chargino and stop. More detailed presentation of the bottom–up mapping and more results can be found in reference \[3\].

## 2 Solutions to the renormalization group equations

We write down first the approximate solutions to the RG equations in the low and moderate \(\tan \beta\) regime (i.e. neglecting the effects of Yukawa couplings other than the top quark one) for the relevant parameters. For our purpose, it is useful to give the expression of the high energy parameters as a function of the low energy ones. They read:

\[
\begin{align*}
\tilde{m}^2(0) &= \frac{m^2(t)}{1-y} + y \left(A_t^2(0) - 2\hat{\xi} A_t(0) M_{1/2}^2\right) + \frac{y\hat{\eta} - y^2\hat{\xi}^2 - \eta}{1-y} M_{1/2}^2 \quad (1) \\
\tilde{m}^2_K(0) &= m^2_K(t) + \frac{c_K}{1-y} \tilde{m}^2(t) + \frac{c_K}{1-y} \left(A_t^2(0) - 2\hat{\xi} A_t(0) M_{1/2}^2\right) \\
&- \left[\eta_K + \frac{c_K}{1-y} \left(\hat{\eta} - \hat{\xi} + y\hat{\xi}^2\right)\right] M_{1/2}^2 \quad (2) \\
A_t(0) &= \frac{A_t(t) + \left(\xi_u - y\hat{\xi}\right) M_{1/2}}{1-y} \quad (3) \\
B(0) &= B(t) + \frac{c_B}{1-y} A_t(t) + \left[\xi_B + \frac{c_B}{1-y} \left(\xi_u - \hat{\xi}\right)\right] M_{1/2} \quad (4) \\
\mu^2(0) &= \frac{\mu^2(t)}{(1-y)^{c_\mu/c}} \prod_{i=1,2,3} \left(\frac{\alpha_i(t)}{\alpha_i(0)}\right)^{a_{i,u}/b} \quad (5)
\end{align*}
\]

where \(\tilde{m}^2(t) \equiv m^2_{Q}(t) + m^2_{U}(t) + m^2_{H_2}(t), \quad \eta \equiv \eta_Q + \eta_U + \eta_{H_2}, \quad t \equiv \frac{1}{2\pi} \log \frac{M}{Q}, \quad M = M_{GUT} \) or any intermediate scale, and \(m^2_K\), with \(K = H_i, Q, U, D\), denotes the Higgs, left handed squark, right handed up-type squark and right handed down-type squark soft supersymmetry breaking mass parameters, respectively. Quantities at \(t = 0\) are the initial values of the parameters at the scale \(M\). The coefficients \(c\) and \(c_K\) read: \(c = 6, \quad c_B = 3, \quad c_Q = 1, \quad c_U = 2, \quad c_{H_2} = 3, \quad c_D = c_{H_1} = 0; \quad c_\mu = 3\). The function \(y \equiv y(t)\) is defined as

\[
y(t) \equiv \frac{Y_t(t)}{Y_f(t)}. \quad (6)
\]

where \(Y_t = h_t^2/4\pi\ (h_t \) is the top quark Yukawa coupling) have the well known form:

\[
Y_t(t) = \frac{Y_t(0) E(t)}{1 + cY_t(0) F(t)} \quad (7)
\]
and its infrared fixed point value, \( Y_f(t) \), is given by

\[
Y_f(t) = \frac{E(t)}{cF(t)}
\]  

with

\[
E(t) = \prod_{i=1,2,3} \left( \frac{\alpha_i(0)}{\alpha_i(t)} \right)^{(a_i/b_i)}
\]

\[
F(t) = \int_0^t E(t') dt'.
\]

Functions \( \xi_j(t) \), \( \eta_K(t) \), \( \hat{\xi}(t) \), \( \hat{\eta}(t) \) are defined by

\[
\xi_j(t) = I \left[ \sum_i \frac{M_i(0)}{M_{1/2}} \frac{a_i^j a_i^2(t)}{\alpha_i(0)} \right]
\]

\[
\eta_K(t) = I \left[ \sum_i \frac{M_i^2(0)}{M_{1/2}^2} \frac{d_K a_i^2(t)}{\alpha_i^2(0)} \right]
\]

\[
\hat{\xi}(t) = H \left[ \sum_i \frac{M_i(0)}{M_{1/2}} \frac{a_i^j a_i^2(t)}{\alpha_i(0)} \right]
\]

\[
\hat{\eta}(t) = H \left[ \sum_i \frac{M_i^2(0)}{M_{1/2}^2} \frac{\overline{d} \alpha_i^2(t)}{\alpha_i^2(0)} \right] + H \left[ \sum_i \frac{M_i(0)}{M_{1/2}} \frac{\overline{d} \alpha_i^2(t)}{\alpha_i(0)} \xi_u(t) \right]
\]

where the coefficients \( a_i^j \) and \( d_i^j \) read: \( a_i^u = (13/15, 3, 16/3) \), \( a_i^d = (7/15, 3, 16/3) \), \( a_i^b = (3/5, 3, 0) \); \( a_i^\nu = (5/3, 3, 0) \); \( d_Q = (1/15, 3, 16/3) \), \( d_L = (16/15, 0, 16/3) \), \( d_D = (4/15, 0, 16/3) \), \( d_{H_1} = d_{H_2} = (3/5, 3, 0) \), \( \overline{d} \equiv d_Q + d_L + d_{H_2} \) while \( H \) and \( I \) are defined in the following way:

\[
H[f(t)] = \int_0^t f(t') dt' - \frac{1}{F(t)} \int_0^t F(t') f(t') dt'
\]

\[
I[f(t)] = \int_0^t f(t') dt'
\]

Factors \( M_i(0)/M_{1/2} \neq 1 \) appear because we do not assume exact gauge coupling unification and \( M_{1/2} \) can be chosen by convention e.g. as \( M_{1/2} \equiv M_3(0) \). Parameters \( \xi_j \) and \( \eta_K \) can be computed analytically but \( \hat{\xi} \) and \( \hat{\eta} \) require numerical integration. We give their typical values in table 1. More values of all relevant parameters can be found in reference 3.
Table 1: Typical values of $\hat{\xi}$ and $\hat{\eta}$.

| $\xi$ | $M$ [GeV] | 2 · 10^{16} | 1 · 10^{10} | 1 · 10^4 | 1 · 10^2 |
|-------|-----------|-------------|-------------|----------|----------|
| $\alpha_3$ | .115 | 2.16 | 1.07 | .640 | .375 |
| | .120 | 2.23 | 1.11 | .664 | .389 |
| | .125 | 2.30 | 1.15 | .687 | .403 |

| $\hat{\eta}$ | $M$ [GeV] | 2 · 10^{16} | 1 · 10^{10} | 1 · 10^4 | 1 · 10^2 |
|--------------|-----------|-------------|-------------|----------|----------|
| $\alpha_3$ | .115 | 12.2 | 4.16 | 2.00 | .988 |
| | .120 | 12.8 | 4.40 | 2.11 | 1.04 |
| | .125 | 13.5 | 4.64 | 2.22 | 1.09 |

The soft SUSY breaking parameters are expressed in terms of the physical parameters according to

\begin{align}
\langle m^2_{H_1}(t) &= \sin^2 \beta M_A^2 + \frac{t_\beta}{2} M_Z^2 - \mu^2 \\
\langle m^2_{H_2}(t) &= \cos^2 \beta M_A^2 - \frac{t_\beta}{2} M_Z^2 - \mu^2 \\
\mu(t)B(t) &= \sin \beta \cos \beta M_A^2 \\
m^2_\tilde{Q}(t) &= M_{t_1}^2 \cos^2 \theta_i + M_{t_2}^2 \sin^2 \theta_i - m_t^2 - \frac{t_\beta}{6} (M_Z^2 - 4 M_W^2) \\
m^2_\tilde{T}(t) &= M_{t_1}^2 \sin^2 \theta_i + M_{t_2}^2 \cos^2 \theta_i - m_t^2 + \frac{2}{3} t_\beta (M_Z^2 - M_W^2) \\
A_t(t) &= \frac{M_{t_1}^2 - M_{t_2}^2}{m_t} \sin \theta_i \cos \theta_i + \mu(t) \cot \beta
\end{align}

where $M_{t_1}^2$, $M_{t_2}^2$ and $\theta_i$ are respectively the heavier and lighter physical top squark masses and their mixing angle while $t_\beta \equiv (\tan^2 \beta - 1)/(\tan^2 \beta + 1)$. One should stress that eqs. (1–5) are valid for general, non-universal values of the soft SUSY breaking mass parameters at the scale $M$ and that unification assumptions for the gauge couplings and for gaugino masses have not been used ($M_{1/2}$ is by convention equal to the gluino mass $M_3$ at the scale $M$). Moreover, the functions $y(t)$ and $Y_f(t)$ defined by eqs. (6,8) are auxiliary functions defined for any scale $M$. The function $y(t)$ is very convenient in presenting the solutions to the RG equations. The whole dependence of the results on the large top Yukawa coupling can be easily expressed in terms of $y$ (see eqs. (1–5)). For $M = M_{\text{GUT}}$ a consistent perturbative treatment of the theory can only be performed if

\[ y_{\text{GUT}} \equiv y \left( t = \frac{1}{2\pi} \log \frac{M_{\text{GUT}}}{M_Z} \right) < 1, \]
where $y_{\text{GUT}} \approx 1$ defines the quasi infrared fixed point solution. For scenarios with supersymmetry broken at lower scales $M \ll M_{\text{GUT}}$, the same values of $m_t$ and $\tan \beta$ (or equivalently of $y_{\text{GUT}}$) give obviously much lower values for the auxiliary function $y(t)$, where $t$ is defined at the scale $Q = M_Z$, ($t = \frac{1}{2\pi} \log \frac{M}{M_Z}$). One should also remember that, in general, for $M < M_{\text{GUT}}$, new matter multiplets are expected to contribute to the running of the gauge couplings above the scale $M$. New complete $SU(5)$ or $SO(10)$ multiplets do not destroy the unification of gauge couplings but their value at the unification scale becomes larger and, correspondingly, the perturbativity bound, $y_{\text{GUT}} < 1$, allows for larger values of the top quark Yukawa coupling (smaller $\tan \beta$) at the $M_Z$ scale. We shall use the auxiliary function $y(t)$ when presenting our results.

3 Results of bottom–up mapping in MSSM.

We recall that our theoretical intuition about the superpartner spectrum has been to a large extent developed on the top-down approach in the minimal supergravity model (with universal soft terms) and on minimal models of low energy supersymmetry breaking. For instance, it is well known that in the minimal supergravity model for low $\tan \beta$ the renormalization group running gives

$$M_Z^2 \approx -2\mu^2 + \mathcal{O}(3)m_0^2 + \mathcal{O}(12)\frac{M_{1/2}^2}{2}$$

(23)

where $M_{1/2}$ and $m_0$ are the universal gaugino and scalar mass parameters (at the high energy scale) respectively. The large coefficient in front of $M_{1/2}^2$ is the well known source of fine tuning for $M_{1/2} > M_Z$ and therefore one naturally expects light neutralino and chargino both gaugino-like since, from (23), $\mu > M_{1/2}$.

For moderate values of the parameters at the GUT scale, the RG evolution of the squark masses gives

$$m_Q^2 = \mathcal{O}(6)M_{1/2}^2 + \mathcal{O}(0.5)m_0^2 + \ldots$$

$$m_U^2 = \mathcal{O}(4)M_{1/2}^2 + \ldots$$

(24)

where ellipsis stand for terms proportional to $A_0$ and $A_0 M_{1/2}$ with coefficients going to zero in the limit $y \to 1$. We see that the hierarchy $m_Q^2 \gg m_U^2$ can be generated provided $\mathcal{O}(1\text{TeV}) \sim m_0 \gg M_{1/2} \sim \mathcal{O}(M_Z)$, i.e. consistently with naturality criterion. For large enough values of the universal trilinear soft terms, $A_0$, one may even obtain $m_U^2 < 0$ and in consequence $M_{\tilde{t}_1} \gg M_{\tilde{t}_2} \sim \mathcal{O}(M_Z)$, $m_{\tilde{c}^+_1} \sim \mathcal{O}(M_Z)$ and gaugino-like.

Of course, Ans"atze for soft terms at high energy scale such as full universality or partial universality selects only a small subset of the low energy parameter space, even if the latter is assumed to contain light chargino and stop and to be characterized by the hierarchy $M_{\tilde{t}_1} \gg M_{\tilde{t}_2}$. Since the sparticle spectrum with $m_{\tilde{\chi}^+_i} \sim m_{\tilde{N}_i^0} \sim M_{\tilde{t}_2} \sim \mathcal{O}(M_Z)$ and $M_{\tilde{t}_1} \gg M_{\tilde{t}_2}$ is important for LEP2 and at the same time consistent with the precision

\footnote{The coefficients in eq. (23) are for $\tan \beta = 1.6$; they decrease with increasing $\tan \beta$. The exact values of the coefficients in eqs. (23,24) depend slightly on the values of $\alpha_s$, $\sin^2 \theta_W$ and $M_{\text{GUT}}$ chosen.}
electroweak data \(^3\), it is of interest to study the mapping of such a spectrum to high energies in a general way and to understand better its consistency (or inconsistency) with various simple patterns.

We shall discuss now some general features of the behaviour of the soft supersymmetry breaking parameters, which may be extracted from eqs. (1-5). We consider first the case \( M = M_{GUT} \) and the limit \( y_{GUT} \to 1 \) as a useful reference frame. One should be aware, however, that already for \( y_{GUT} \simeq 0.8-0.9 \) large qualitative departures from the results associated with this limit are possible. In the limit \( y_{GUT} \to 1 \), from eqs. (1) and (3), we see that if the high energy parameters are of the same order of magnitude as the low energy ones, the following relations must be fulfilled:

\[
\begin{align*}
\Delta m^2 &\equiv m^2(t) - \left(\eta - y\hat{\eta} + y^2\hat{\xi}^2\right) M^2_{1/2} \to 0 \\
\Delta_A &\equiv A_t(t) - \left(y\hat{\xi} - \xi_u\right) M^2_{1/2} \to 0,
\end{align*}
\]

irrespectively of their initial values (IR fixed point). This is a well known prediction, which, for \( y_{GUT} \neq 1 \), remains valid in the gaugino dominated supersymmetry breaking scenario in the minimal supergravity model.

Let us now suppose that the relations (25) are strongly violated by experiment, namely the scalar masses (or the soft supersymmetry breaking parameter \( A_t \)) are much different than predicted by (25), for the values of \( m_t \) and \( \tan \beta \) corresponding to \( y_{GUT} \) close to 1. Clearly, this means that

\[
\begin{align*}
A_t(0) &\sim \mathcal{O}\left(\frac{\Delta_A}{1 - y}\right) \\
\overline{m}^2(0) &\sim \mathcal{O}\left(\frac{\Delta m^2}{1 - y}\right) + \mathcal{O}(A_t^2(0))
\end{align*}
\]

i.e. supersymmetry breaking must be driven by very large initial scalar masses and the magnitude of the effect depends on the departure from the fixed point relations for \( A_t \) and \( \overline{m}^2 \), eq. (25), and the proximity to the top quark mass infrared fixed point solution measured by \( 1 - y_{GUT} \). Moreover, for not very small values of \( |\Delta_A| \) and/or \( |\Delta m^2| \) it follows from eq. (2) that in the limit \( y_{GUT} \to 1 \), independently of the actual values of the masses at the scale \( M_Z \), their initial values are correlated such that

\[
m_Q^2(0) : m_U^2(0) : m_{H_2}^2(0) \simeq 1 : 2 : 3
\]

and \( m_{H_1}^2(0) \) is much smaller than the above three soft masses. In other words, the values at \( M_Z \) are obtained (via RGEs) by a very high degree of fine tuning between the initial values \( m_{K_R}^2(0) \). We conclude that eqs. (25) are necessary (but of course not sufficient) conditions for large departures from the prediction (27) for the soft scalar masses in the limit \( y \to 1 \). In particular, they are necessary conditions for the spectrum consistent with fully (or partially) universal initial values of the third generation squark and Higgs boson soft masses.
Figure 1: Soft supersymmetry breaking parameters at the scale \( M = M_{\text{GUT}} = 2 \cdot 10^{16} \) GeV, for \( y = 0.98 \) obtained by mapping the low energy parameter space specified in the text.

When eqs. (25) are violated, there are essentially two ways of departing from the prediction (27). One is to increase the value of \( \tan \beta \) (i.e. to decrease the top quark Yukawa coupling for fixed \( m_t \)). The other way is to lower the scale at which supersymmetry breaking is transferred to the observable sector, since for \( M \ll M_{\text{GUT}} \) the soft supersymmetry breaking parameters do not feel the strong rise of the top quark Yukawa coupling at scales close to \( M_{\text{GUT}} \). In both cases, \( y \) takes values smaller than one and we can depart from (27) even when (25) are not satisfied.

We shall illustrate the above general considerations by numerical bottom–up mapping of the low energy parameter space characterized as follows: \( m_{\chi^+} = 90 \) GeV, \( M_{\tilde{t}} = 60 \) GeV (the pattern of soft supersymmetry breaking parameters does not depend strongly on the exact value of the light chargino and stop masses), \( 0.1 < M_2/|\mu| < 10 \) for both signs of \( \mu \). To reduce the parameter space, in our numerical analysis we make the assumption that \( M_1/\alpha_1 = M_2/\alpha_2 = M_3/\alpha_3 \) at any scale. For \( m_t = 175 \) GeV, \( \alpha_s(M_Z) = 0.118 \) we consider two examples: \( M = M_{\text{GUT}} = 2 \times 10^{16} \) GeV, \( y_{\text{GUT}} = 0.98 \) (corresponding to \( \tan \beta \approx 1.6 \)) and \( M = 10^7 \) GeV, \( y = 0.75 \) (corresponding to \( \tan \beta \approx 1.25 \)).

We scan over the remaining relevant parameters \( M_A, M_{\tilde{t}_1} \) < 1 TeV, and the whole range of \( \theta_i \). The accepted low energy parameter space is then defined as a subspace in which \( M_h > 60 \) GeV, \( 0.98 \times 10^{-4} < BR(b \rightarrow s\gamma) < 3.66 \times 10^{-4} \), and we also require that \( \chi^2 \leq \chi^2_{\text{min, SM}} + 2 \) where \( \chi^2 \) is for the fit to the electroweak observables which do not involve heavy flavors.

In Figure 1 we show the dependence of \( m_U/m_Q \) on the value of \( M_{1/2} \) and the behaviour of the soft supersymmetry breaking Higgs mass parameters in the infrared limit: \( y = 0.98 \) (masses are defined by \( m_K = m_K^2/\sqrt{|m_K^2|} \)). A clear concentration of solutions around \( m_U/m_Q = \sqrt{2} \) and \( m_{H_2}/m_Q = \sqrt{3} \) appears. The above ratios are
Figure 2: Same as Figure 1, but for $M = 10^7$ GeV, $y=0.75$.

Figure 2 shows the solutions obtained for low supersymmetry breaking scale $M = 10^7$ GeV and for $y = 0.75$. The value of $y$ is now much smaller than 1 and the violation of the sum rules, eq. (25), does not imply an enhancement of the values of the parameters at the scale $M$. From the analysis of our equations, for these low values of $M$ and $y(t)$, it follows that the initial values of soft terms at the scale $M$ tend to reflect the pattern observed at the $M_Z$ scale. Due to the small renormalization group running effects and the fact that $\eta_Q \simeq 0.85$, the ratio $m_Q/M_{1/2}$ is just a reflection of the chosen values of the left handed stop parameter and the gluino mass at low energies, increasing for lower values of $M_{1/2}$. More interesting is the behaviour of $m_U/m_Q$. The concentration of solutions around $m_U/m_Q \simeq \sqrt{2}$ disappears for this low values of $y$ and, instead, the ratio $m_U/m_Q$ tends, in general, to be lower. Hence, a light stop is not necessarily in conflict with models of supersymmetry breaking in which $m_U/m_Q \simeq 1$ at the scale $M$. Figure 2 shows that for SUSY breaking scale of order $10^7$ GeV, the three soft masses $m_Q$, $m_U$ and $m_{H_2}$ can have similar values for many solutions, especially for light gauginos. The second soft Higgs mass, $m_{H_1}$, tends to be rather light but the full universality can also be easily obtained.

governed by the size of the gaugino masses. For low values of $M_{1/2}$, the maximum value of $m_U/m_Q$ is given by $\sqrt{2}$, while for solutions with large values of $M_{1/2} = O(1 \text{ TeV})$, $m_U/m_Q > \sqrt{2}$ is possible. Similarly, values of $m_{H_2}/m_Q \geq \sqrt{3}$ may only appear for large values of the gaugino masses, while low values of the gaugino masses always lead to $m_{H_2}/m_Q \leq \sqrt{3}$. For $y \simeq 1$, a strong concentration of solutions around the boundary values is observed. Such behaviour of the solutions is just a reflection of the properties discussed above for $y$ close to 1 and shows the global tendency in the mapping of the low energy region selected by our criteria.
4 Conclusions

In this paper we have discussed the mapping to high energies of the low energy parameters of the MSSM. We investigated the specific region in low energy parameter space with light chargino and right-handed stop which can be interested for LEP2. For heavy top quark and small or moderate tan $\beta$ the global pattern of the mapping is determined by the proximity of the top quark Yukawa coupling to its IR fixed point value and by the assumed scale at which supersymmetry breaking is transmitted to the observable sector. The general pattern of this mapping is the dominance of the scalar masses (over the gaugino mass) in the supersymmetry breaking. We have also identified the conditions which are necessary for the spectrum to be consistent with a simple Ansatz like universality (or partial universality) of the high energy values for the scalar masses. In particular, the $SO(10)$—type initial conditions, with universal squark masses but with non-universal Higgs masses are compatible with an interesting subregion of the considered parameter space.

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