Challenging the minimal supersymmetric SU(5) model

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Abstract. We review the main constraints on the parameter space of the minimal renormalizable supersymmetric SU(5) grand unified theory. They consist of the Higgs mass, proton decay, electroweak symmetry breaking and fermion masses. Superpartner masses are constrained both from below and from above, giving hope for confirming or definitely ruling out the theory in the future. This contribution is based on Ref. [1].

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INTRODUCTION

It is often claimed that the minimal supersymmetric SU(5) grand unified theory is ruled out. This statement is based on the analysis of Ref. [2]: gauge coupling unification constrains the colour Higgs triplet to be lighter than \( 3.6 \times 10^{15} \) GeV, while the non-observation of proton decay requires it to be significantly heavier. If two massive colour triplets were available, like in the model with an extra vectorlike fundamental representation [3], both constraints would be compatible, but in the minimal supersymmetric SU(5) model the same triplet cannot be (relatively) light and heavy at the same time.

This analysis relies on two implicit assumptions:

- the predictions of the model are not affected by the presence of non-renormalizable operators;
- the soft supersymmetry breaking masses do not exceed the TeV scale, apart from the first two generations of sfermions, whose masses can be as large as 10 TeV.

How much does the conclusion of the above analysis depend on these assumptions? Allowing non-renormalizable operators can turn the theory back to life [4]: on one side the masses of the colour octet and weak triplet inside the adjoint Higgs field can now differ, which relaxes the unification constraint on the colour triplet mass [5, 6]; on the other side the new terms allow for a much more flexible flavour structure in both the sfermion [7] and the fermion sectors [8, 9], thus relaxing the naive bound on the triplet mass from proton decay [10, 11, 12]. Of course this is possible at the expense of a large number of unknown parameters, which makes the theory less predictable.

There is however an argument against the presence of such non-renormalizable terms. We know that the Planck-suppressed operator

\[
\frac{\kappa}{M_{\text{Planck}}} \propto 10_F^1 \times 10_F^1 \times 10_F^2 \times 5_F^1,
\]

where \( 1, 1, 2, I \) are generation indices, must be further suppressed by \( \kappa \lesssim 10^{-7} \) in order to satisfy the strong proton decay bounds. This may tell us that all Planck-suppressed operators enjoy some extra suppression. After all, we do not really know how gravity contributes to the effective field theory below \( M_{\text{Planck}} \). It may well be that these operators are exponentially suppressed or even vanish exactly. This we will assume throughout this talk.

We are thus left with the second option, namely considering larger soft superpartner masses, which by the way is also favoured by flavour physics constraints and by the mass of the recently discovered Higgs boson. This is what this talk is about: we will require the superpartner masses to fit all the experimental constraints in the minimal

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renormalizable supersymmetric SU(5) model. All we assume are SU(5)-invariant boundary conditions for the soft terms, i.e. supersymmetry must be broken above $M_{GUT}$, as in supergravity. We will see that this is not a very constraining assumption: in fact flavour-blind supersymmetry breaking mechanisms like gauge mediation will not be available, since satisfying the various experimental constraints requires a large splitting between sfermion generations. As we will show, the third generation (as well as the higgsinos and the heavy Higgs bosons) must be very heavy ($\sim 10^{2-3}$ TeV) in order not to destabilize the electroweak vacuum, while the first two generations of sfermions (and the gauginos) need to be lighter ($\sim 10$ TeV) in order to be able to correct the SU(5) fermion mass relations.

**THE MINIMAL RENORMALIZABLE SUPERSYMMETRIC SU(5) MODEL**

Let us start with a short description of the minimal renormalizable supersymmetric SU(5) model. The Higgs sector consists of the adjoint $24_H$, whose vacuum expectation value

$$\langle 24_H \rangle = V \text{ Diag}(2,2,2,-3,-3)$$

is responsible for the spontaneous gauge symmetry breaking $SU(5) \rightarrow SU(3)_C \times SU(2)_L \times U(1)_Y$. This vev is found as a minimum of the scalar potential defined by the superpotential:

$$W_{24} = \frac{\mu}{2} \text{Tr}(24_H^2) + \frac{\lambda}{3} \text{Tr}(24_H^3),$$

with

$$V = \frac{\mu}{\lambda}. \quad (4)$$

Expanding the above superpotential around this minimum, the masses of the weak triplet and colour octet components of the adjoint Higgs superfield are found to be:

$$m_3 = m_8 = 5\lambda V. \quad (5)$$

The MSSM Higgs doublet pair lives in the $5_H \oplus \bar{5}_H$ representation, whose superpotential couplings are:

$$W_5 = \bar{5}_H (m + \eta 24_H) 5_H. \quad (6)$$

After breaking of the SU(5) gauge symmetry, the colour triplet and electroweak doublet components of the Higgs fields $5_H \oplus \bar{5}_H$ acquire the following masses:

$$M_T = m + 2\eta V, \quad (7)$$

$$m_D = m - 3\eta V. \quad (8)$$

In order to recover the MSSM at low energy and to prevent too fast proton decay, the doublet-triplet splitting is achieved by means of the fine-tuning:

$$m = 3\eta V, \quad (9)$$

so that the mass of the heavy colour triplet is:

$$M_T = 5\eta V. \quad (10)$$

Next, the Yukawa sector of the model is given by:

$$W_{\text{Yukawa}} = Y_{ij}^{10} 10_F^i 10_F^j 5_H + Y_{ij}^{5\bar{5}} 10_F^i \bar{5}_H, \quad (11)$$

where $i, j$ are generation indices running from 1 to 3. This simple structure leads to the following predictions at the GUT scale:

$$M_U = M_U^f, \quad (12)$$

$$M_D = M_D^f. \quad (13)$$

The equality between the GUT-scale down quark and charged lepton masses leads to particularly bad predictions for the first two generations and has to be corrected. We will do that with the only resource we still have, i.e. supersymmetric threshold corrections (for a recent analysis see for example Ref. [13] and references therein).
Finally, the kinetic sector gives the heavy gauge bosons a mass:

\[ M_V = 5g_5 V. \]  

(14)

Comparing Eq. (10) with Eq. (14) and assuming perturbativity (i.e. \( \eta \lesssim 1 \)), we thus obtain

\[ M_T \lesssim M_V, \]  

(15)

since \( g_5 \approx 0.7 \).

**OUR ASSUMPTIONS**

It is time now to describe in detail the assumptions and inputs that we will use in our analysis:

- we consider SU(5)-invariant, but otherwise arbitrary soft supersymmetry breaking terms;
- we insist on our electroweak vacuum to be the global minimum of the scalar potential. We will in particular impose the (strictly speaking not necessary) condition that

\[ m_{\tilde{H}_u}^2 > 0 \]  

(16)

at all energies between \( M_{\text{GUT}} \) and the electroweak symmetry breaking scale

\[ M_{\text{EWSB}} \equiv \sqrt{\tilde{m}_1(M_{\text{EWSB}})\tilde{m}_2(M_{\text{EWSB}})}, \]  

(17)

where the matching between the MSSM and the SM is done;

- the fermion masses receive corrections dominantly from 1-loop gluino exchange;
- the \( d = 5 \) proton decay operator is dominated by 1-loop wino exchange;
- in order to minimize the dangerous supersymmetric contributions to flavour-changing processes, we assume the masses of the first two generations of sfermions to be approximately equal:

\[ \tilde{m}_1 \approx \tilde{m}_2 \left( \equiv \tilde{m}_{1,2} \right). \]  

(18)

**FERMION MASS CORRECTIONS**

As already mentioned, in the minimal renormalizable supersymmetric SU(5) model, the GUT-scale mass relations are only corrected by 1-loop finite supersymmetric threshold corrections (for a recent work see for example [13] and references therein). We will only consider the corrections to down quark masses, which are enhanced (compared with the corrections to charged lepton masses) by the strong coupling constant:

\[ \delta m_i = -\frac{2\alpha_3}{3\pi} \frac{m_e^2}{\tilde{m}_{1,2}} \left( A_i \cos \beta - \mu h_i \sin \beta \right) I_1 \left( \frac{m_e^2}{\tilde{m}_{1,2}} \right), \]  

(19)

where \( m_i \) and \( A_i \) are the down quark masses and the associated A-terms, and the loop function \( f \) is given by:

\[ I_1(x) \equiv \frac{1 - x + x \log x}{(1 - x)^2}. \]  

(20)

Comparing the SU(5) predictions for fermion masses (obtained by running the GUT-scale values down to low energy) with the experimental values one finds that these corrections need to be positive for the down quark and negative for the strange quark. This means that the term proportional to \( \mu \) cannot dominate Eq. (19).

So the dominant part comes from the \( A_i \)’s, which however are bounded by vacuum stability constraints [14]:

\[ A_i = a_i h_i \sqrt{3} \left( 2\tilde{m}_{1,2}^2 + \tilde{m}_{1,2}^4 \right)^{1/2}, \quad |a_i| \leq 1. \]  

(21)
Since the required corrections to the down and strange quark masses are larger than the masses themselves, large $A_i$ terms are needed. Then the only way to maintain vacuum stability while avoiding a suppression of the corrections (19) by large squark or gluino masses is to assume
\[ m_{H_d} \gg \tilde{m}_{1,2}, m_{\neq}. \tag{22} \]

Such a heavy $H_d$ gives a large 1-loop contribution to the hypercharge D-term (by contrast it only receives small contributions from the sfermions, due to the SU(5) boundary conditions on soft masses), which yields tachyons in the superpartner spectrum unless it is compensated for by the contribution of $H_u$. We will enforce this solution by assuming
\[ m_{H_u}^2(M_{\text{GUT}}) = m_{H_d}^2(M_{\text{GUT}}) \equiv m_{H_d}^2. \tag{23} \]

The large value of $m_{H_u}$ in turn tends to dominate the renormalization group equations of the stop masses and drive them negative at low energy. To avoid tachyons, the soft mass parameter $m_{10_i}$ (and to a lesser extent $m_{3i}$) must be large, hence the third generation sfermions must be heavier than the first two generation sfermions (and gauginos, according to Eq. (22)).

For a fixed value of $m_{H_u}$, the values of the $A$-terms are bounded by Eq. (21), which together with the requirement that Eq. (19) accounts for the measured down quark masses implies an upper limit on the soft masses $\tilde{m}_{1,2} \sim m_{\neq}$.

**PROTON DECAY**

Here the tendency for the soft masses is opposite than in the previous section: the higher the supersymmetry breaking scale, the longer the proton lifetime. Our estimate of the proton decay lifetime in the minimal renormalizable supersymmetric SU(5) model is (using the hadronic matrix elements of Ref. [15]):
\[ \tau(p \to K^+\bar{\nu}) \approx 2 \times 10^{32} \text{ yrs} \left( \frac{\tilde{m}_{1,2}}{10 \text{ TeV}} \right)^2 \left( \frac{1/3 I_1/(1/9)}{m_{\tilde{u}}/\tilde{m}_{1,2} I_1(m_{\tilde{u}}/\tilde{m}_{1,2})} \right)^2 \left( \frac{2 \tan \beta}{1 + \tan^2 \beta} \right)^2 \left( \frac{M_T}{10^{17} \text{ GeV}} \right)^2. \tag{24} \]

This is to be compared with the experimental constraint $\tau(p \to K^+\bar{\nu}) > 2.3 \times 10^{33}$ yrs (90% C.L.) [16]. Heavy colour triplet and first two generation sfermions are favoured, as well as a small value of $\tan \beta$. With $M_T = 10^{17}$ TeV and $\tan \beta = 1.7$, the experimental bound is saturated by Eq. (24) for $\tilde{m}_{1,2} \approx 30$ TeV.

**GAUGE COUPLING UNIFICATION CONSTRAINTS**

The colour triplet and superpartner masses are actually not independent of each other when gauge coupling unification constraints are taken into account. One can solve numerically the 2-loop renormalization group equations (RGEs) for gauge couplings with the top quark Yukawa coupling evolved at 1-loop only (the other Yukawa couplings are neglected, including the bottom quark one, since the proton decay constraint forces $\tan \beta$ to be small). The matching between the SM and the MSSM is performed at the scale $M_{\text{EWSB}}$, while the one between the MSSM and the SU(5) theory is done at $M_{\text{GUT}}$. The mass splittings in the MSSM and in the GUT spectra are taken into account with 1-loop threshold corrections. Putting everything together, one arrives at the following constraints:
\[ \frac{M_T}{M_{\text{GUT}}} = \exp \left[ \frac{5\pi}{6} (-\alpha_i^{-1} + 3\alpha_2^{-1} - 2\alpha_3^{-1}) \right] \left( \frac{m_3}{m_{\neq}} \right)^{5/2} \]
\[ \times \left( \frac{m_{\tilde{u}}}{m_{\tilde{q}}} \right)^{5/3} \prod_{i=1}^{3} \left( \frac{m_{Q_i}^4}{m_{Q_i}^2 m_{\tilde{e}_i} m_{\tilde{e}_i}} \right)^{1/12} \left( \frac{m_{A_i}^2}{M_{\text{EWSB}}^2} \right)^{1/6}, \tag{25} \]
\[ \frac{[M_{\text{GUT}}^{2}(m_3 m_{\neq})]^{1/2}}{M_{\text{GUT}}} = \exp \left[ \frac{\pi}{18} (5\alpha_i^{-1} - 3\alpha_2^{-1} - 2\alpha_3^{-1}) \right] \left( \frac{m_3}{m_{\neq}} \right)^{5/12} \]
\[ \times \left( \frac{M_{\text{EWSB}}^2}{m_{\tilde{u}} m_{\tilde{q}}} \right)^{1/3} \prod_{i=1}^{3} \left( \frac{m_{\tilde{e}_i} m_{\tilde{e}_i}}{m_{\tilde{Q}_i}^2} \right)^{1/36}. \tag{26} \]
where we kept the dependence on \( m_3 \) and \( m_8 \) separately, although \( m_3 = m_8 \) holds in the minimal renormalizable SU(5) model. The \((\alpha_i)_{2-loop}(M_{\text{GUT}})\) are calculated using only SM/MSSM RGEs without thresholds, so they do not unify. Notice that the matching scales \( M_{\text{EWSB}} \) and \( M_{\text{GUT}} \) drop out at leading order as they must \(^3\).

In the approximation where all superpartners lie at a single scale \( M_{\text{EWSB}} \) and gauge couplings are evolved with 1-loop RGEs, the colour triplet and heavy gauge boson masses scale as \[^{12}\] :

\[
M_T \propto M_{\text{EWSB}}^{5/6},
\]

\[
M_V \propto M_{\text{EWSB}}^{2/9},
\]

This is good to memorize. It tells us that the partial proton lifetime due to \( d = 5 \) operators goes essentially as

\[
\tau_p(d = 5) \propto M_T^{22/5},
\]

i.e., it increases even faster with the mass of the heavy mediator than the usual gauge exchange mode of non-supersymmetric theories. On the other hand, the dependence of \( M_V \) on \( M_{\text{EWSB}} \) is rather weak, and according to Eq. (26), the perturbativity constraint (15) can be satisfied by decreasing the colour octet and weak triplet masses, i.e., by decreasing the parameter \( \lambda \) in Eq. (5).

Reintroducing the mass splittings in the superpartner spectrum in Eqs. (25) and (26), one finds that the most important contribution to \( M_T \) (resp. \( M_V \)) comes from the higgsino mass (resp. from gaugino masses) \[^{12}\].

THE HIGGS BOSON MASS

Since the third generation sfermions are constrained to be very heavy, the electroweak symmetry breaking scale \( M_{\text{EWSB}} \) defined by Eq. (17) is necessarily much higher than the electroweak scale \( v = 174 \) GeV. In this case, the standard MSSM Higgs mass formula cannot be applied: one must decouple the heavy superpartners at the scale \( M_{\text{EWSB}} \) and evolve the Higgs quartic coupling \( \lambda \) with the Standard Model RGE down to the electroweak scale, where the Higgs mass is given at leading order by the relation \( m_h = \sqrt{2}\lambda v^2 \).

Neglecting the splittings between superpartner masses, the Higgs mass determines a unique relation between \( \tan \beta \) and \( M_{\text{EWSB}} \) through the boundary condition \( \lambda (M_{\text{EWSB}}) = (g^2 + g'^2)(M_{\text{EWSB}}) \cos^2 2\beta / 4 \). Solving the 1-loop Standard Model RGE for the Higgs quartic coupling, one obtains

\[
\tan \beta = 1.7 \Rightarrow M_{\text{EWSB}} = 411 \text{ TeV}.
\]

These are the values we will consider in the following.

THE GUT AND HEAVY SUPERPARTNER SPECTRUM

We already learnt that \( m_{H_u} \) must be large in order for the needed threshold corrections to down quark masses to be consistent with the absence of dangerous charge and colour breaking minima. We have also seen that imposing the equality of \( m_{H_u} \) and \( m_{H_d} \) at \( M_{\text{GUT}} \) prevents the generation of a large hypercharge D-term that would induce tachyons in the spectrum. It follows that both \( m_{H_u} \) and \( m_{H_d} \) are large over an important energy range (in fact, only \( m_{H_u} \) runs sizably between \( M_{\text{GUT}} \) and \( M_{\text{EWSB}} \)), which has an important impact on the running of the third generation sfermion and Higgs soft masses. Since we consider small values of \( \tan \beta \), for which \( \lambda_b \ll \lambda_t \), this impact is most significant for the stop and \( H_u \) soft masses:

\[
m^2_{H_u}(m) = m^2_H - \frac{2\tilde{m}^2_{103} + m^2_H}{2} \left[ 1 - \left( \frac{m}{M_{\text{GUT}}} \right)^{3\lambda^2_t(m)/4\pi^2} \right],
\]

\[
\tilde{m}^2_{23}(m) = \tilde{m}^2_{103} - \frac{2\tilde{m}^2_{103} + m^2_H}{3} \left[ 1 - \left( \frac{m}{M_{\text{GUT}}} \right)^{3\lambda^2_t(m)/4\pi^2} \right],
\]

\(^3\) In particular, the combination \( M_{\text{GUT}} \exp \left[ \frac{\pi}{6} (\alpha^{-1} + 3\alpha_2^{-1} - 2\alpha_3^{-1})_{2-loop}(M_{\text{GUT}}) \right] \) does not depend on \( M_{\text{GUT}} \) at the 1-loop level.
$\langle \tilde{m}_{10}, m_H \rangle$ plane corresponds to the intersection of the black curve with the hatched area (here for $\tan \beta = 1.7$, $M_{\text{EWSB}} = 411$ TeV and $M_{\text{GUT}} = 10^{17}$ GeV). See text for details.

\[ \tilde{m}^2_{Q_3}(m) = \tilde{m}^2_{10_{1b}} - \frac{2\tilde{m}^2_{10_{1b}} + m^2_H}{6} \left[ 1 - \left( \frac{m}{M_{\text{GUT}}} \right)^{\frac{3\lambda^2(m)/4\pi^2}{\log(M_{\text{GUT}}/m)}} \right], \]  

in which the subleading Yukawa and gauge contributions have been neglected\(^4\), and

\[ \tilde{\lambda}^2(m) = \int_{\log m}^{\log M_{\text{GUT}}} \lambda^2(m') d\log m' \left( \log(M_{\text{GUT}}/m) \right). \]  

Now we can look for the region of the parameter space satisfying the constraints:

\[ m^2_{H_u} > 0, \]  
\[ \tilde{m}^2_{Q_3}, \tilde{m}^2_{u_3} > 0, \]  
\[ |\mu|^2 = \frac{m^2_{H_d} - m^2_{H_u} \tan^2 \beta}{\tan^2 \beta - 1} > 0, \]

at the electroweak symmetry breaking scale (17), which in the absence of a significant stop mixing is well approximated by:

\[ M_{\text{EWSB}} = \sqrt{\tilde{m}_{Q_3}(M_{\text{EWSB}}) \tilde{m}_{u_3}(M_{\text{EWSB}})}. \]

Condition (37) is equivalent to saying that electroweak symmetry breaking is possible, while condition (35) ensures the absence of charge and colour breaking minima deeper than our electroweak vacuum. Since, for a given value of $\tan \beta$, $M_{\text{EWSB}}$ is determined by the Higgs boson mass, Eq. (38) implies a unique relation between $m_H\equiv m_{H_u}(M_{\text{GUT}}) = m_{H_d}(M_{\text{GUT}})$ and $\tilde{m}_{10_{1b}} \equiv \tilde{m}_{Q_3}(M_{\text{GUT}}) = \tilde{m}_{u_3}(M_{\text{GUT}})$, which is represented by the black curve on Fig. 1. The light blue region in the top left, the magenta and grey regions in the lower right satisfy the inequalities (35), (36) and (37), respectively, while the hatched region is the portion of the parameter space allowed by all three constraints. Notice that relaxing the inequality (35) would just allow the points along the black curve below the shaded region, towards smaller values of $m_H$, not enlarging significantly the parameter space.

Thus, at the level of approximation described above, the region of the parameter space of the minimal renormalizable supersymmetric SU(5) model compatible with all experimental constraints depends on a single parameter, say $\tilde{m}_{10_{1b}}$. Taking for example:

\[ \tilde{m}_{10_{1b}} = 2000 \text{ TeV}, \]  

\(^4\) Let us recall that $m_{\tilde{g}}, m_{\tilde{t}} \ll m_{H_d}$, or in terms of GUT-scale parameters $M_{1/2} \ll m_H$. \n
FIGURE 2. The allowed region in the \((\tilde{m}_{1,2}, m_{\tilde{g}})\) plane corresponds to the hatched area (here for \(\tan \beta = 1.7, M_{\text{FWSB}} = 411\) TeV, \(M_{\text{GUT}} = 10^{17}\) GeV and \(\tilde{m}_{10_3} = 2000\) TeV). See text for details.

we find:

\[
\begin{align*}
    m_h &= \mu = 677\ \text{TeV}, \\
    m_A &= \sqrt{(\mu^2 + m_{H_d}^2)(1 + 1/\tan^2 \beta)} = 3508\ \text{TeV},
\end{align*}
\]

where \(m_A\) is the mass of the (approximately degenerate) heavy MSSM Higgs bosons. Plugging these parameters into Eq. (25) and neglecting the mass splittings within the first two generations of sfermions, and using the approximate relation \(m_{\tilde{u}}/m_{\tilde{g}} \simeq (\alpha_2/\alpha_3)(m_{\tilde{g}})\), we obtain for \(M_{\text{GUT}} = 10^{17}\) GeV and \(\tilde{m}_{1,2} = m_{\tilde{g}} = 11\) TeV:

\[
\begin{align*}
    M_T &= 4.0 \times 10^{17}\ \text{GeV}, \\
    [M_T^2(m_3 m_{\tilde{g}})^{1/2}]^{1/3} &= 8.5 \times 10^{15}\ \text{GeV}.
\end{align*}
\]

The exact value of the sfermion and gaugino masses turn out to have a small effect on the colour triplet mass, while the combination of masses \([M_T^2(m_3 m_{\tilde{g}})^{1/2}]^{1/3}\) is only weakly dependent on sfermion masses.

THE "LIGHT" SUPERPARTNER SPECTRUM

We are now in a position to determine also the ranges of the gaugino and first two generation sfermion masses for which the predicted proton lifetime and the fermion masses are consistent with experiment. The four different regions in the \((\tilde{m}_{1,2}, m_{\tilde{g}})\) plane on Fig. 2 represent from left to right the following constraints: the purple satisfies both \(|a_d| < 1\) and \(|\alpha_3| < 1\), the hatched corresponds to the allowed region (all three constraints satisfied), the brown means both \(|\alpha_3| < 1\) and \(\tau_p > \tau_p^{\text{exp}}\), while the extreme right region satisfies only the proton decay bound.

CONCLUSIONS

We have identified a region in the parameter space of the minimal renormalizable supersymmetric SU(5) model that is consistent with all experimental and theoretical constraints: gauge coupling unification, the measured charged fermion and Higgs boson masses, the absence of charged and colour breaking vacua and the experimental limit on the proton lifetime. The analysis has been simplified by making suitable approximations. Some single points in the allowed parameter space have been studied with a better precision, and the results are compatible with the estimates presented in this talk.
Let us finish with a short comment about neutrino masses and dark matter. In the absence of additional multiplets like SU(5) singlets, the only possibility here seems to include bilinear R-parity violating terms (for a review see for example Ref. [17]). After a generalized doublet-triplet splitting, there are enough parameters to fit the neutrino masses and mixings. In this context, the only dark matter candidate is a light, order GeV or less [18] gravitino.

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