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Phenomenology of $SU(5)$ Finite Unified Theories

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Abstract. Finite Unified Theories (FUTs) are $N = 1$ supersymmetric Grand Unified Theories (GUTs) which can be made finite even to all-loop orders, leading to a large reduction in the number of free parameters. We confront the predictions of $SU(5)$ FUTs with the top and bottom quark masses, which allows us to discriminate among different models. We include further low-energy phenomenology constraints, such as $B$ physics observables, the bound on the SM Higgs mass and the cold dark matter density, and then are able to make predictions for the lightest Higgs boson mass and the sparticle spectrum.

1. Introduction

Finite Unified Theories are $N = 1$ supersymmetric GUTs which can be made finite even to all-loop orders, including the soft supersymmetry (SUSY) breaking sector, see [1–5] for details and further references. The constructed finite unified $N = 1$ supersymmetric GUTs predicted correctly from the dimensionless sector (Gauge-Yukawa unification), among others, the top quark mass [6,7]. The search for finiteness has been extended to the soft supersymmetry breaking sector (SSB) of these theories [8,9], which involves parameters of dimension one and two. Thus, it was possible to extend the predictive power to the Higgs sector and the SUSY spectrum. This, in turn, allows to make predictions for low-energy precision and astrophysical observables.

In here, we present an exhaustive search of $SU(5)$-based finite SUSY models, taking into account the restrictions resulting from the low-energy observables [10]. Finally, the predictions of the “best” model (i.e. that is still allowed after taking the phenomenological restrictions into account) for the Higgs and SUSY searches at the LHC are reviewed [10].

Finiteness can be understood by considering a chiral, anomaly free, $N = 1$ globally supersymmetric gauge theory based on a group G with gauge coupling constant $g$. The superpotential of the theory is given by

$$W = \frac{1}{2} m^{ij} \Phi_i \Phi_j + \frac{1}{6} C^{ijk} \Phi_i \Phi_j \Phi_k ,$$ (1)

where $m^{ij}$ (the mass terms) and $C^{ijk}$ (the Yukawa couplings) are gauge invariant tensors and the matter field $\Phi_i$ transforms according to the irreducible representation $R_i$ of the gauge group.
G. All the one-loop $\beta$-functions of the theory vanish if the $\beta$-function of the gauge coupling $\beta_g^{(1)}$, and the anomalous dimensions of the Yukawa couplings $\gamma_{ij}^{(1)}$, vanish, i.e.

$$
\sum_i \ell(R_i) = 3C_2(G), \quad \frac{1}{2}C_{ipq}C^{ipq} = 2\delta_i^j g^2 C_2(R_i),
$$

where $\ell(R_i)$ is the Dynkin index of $R_i$, and $C_2(G)$ is the quadratic Casimir invariant of the adjoint representation of $G$. These conditions are enough to guarantee two-loop finiteness too [11]. A striking fact is the existence of a theorem [12–14] that guarantees the vanishing of the $\beta$-functions to all-orders in perturbation theory. This requires that, in addition to the one-loop finiteness conditions (2), the Yukawa couplings are reduced in favour of the gauge coupling to all-orders (see [5] for details). Alternatively, similar results can be obtained [15–17] using an analysis of the all-loop NSVZ gauge beta-function [18, 19].

Consider the superpotential given by (1) along with the Lagrangian for SSB terms

$$
-L_{SB} = \frac{1}{6} h^{ijk} \phi_i \phi_j \phi_k + \frac{1}{2} b^{ij} \phi_i \phi_j + \frac{1}{2} \left( m_i^2 \phi^* i \phi_j + \frac{1}{2} M \lambda \lambda + \text{H.c.} \right),
$$

where the $\phi_i$ are the scalar parts of the chiral superfields $\Phi_i$, $\lambda$ are the gauginos and $M$ their unified mass. The one- and two-loop finiteness for $h^{ijk}$ can be achieved by

$$
h^{ijk} = -MC^{ijk} + \cdots = -M\rho^{ijk} g + O(g^5).
$$

Furthermore, it was found that RGI SSB scalar masses in Gauge-Yukawa unified models satisfy a universal sum rule at one-loop [20]. This result was generalized to two-loops for finite theories [4], and then to all-loops for general Gauge-Yukawa and finite unified theories [21]. Then the following soft scalar-mass sum rule is found [4]

$$
\frac{\left( m_i^2 + m_j^2 + m_k^2 \right)}{MM^1} = 1 + \frac{g^2}{16\pi^2} \Delta^{(2)} + O(g^4)
$$

for $i, j, k$ with $\rho^{ijk} \neq 0$, where $m_{i,j,k}^2$ are the scalar masses and $\Delta^{(1)}$ is the two-loop correction

$$
\Delta^{(2)} = -2 \sum_i \left[ (m_i^2/MM^1) - (1/3) \right] \ell(R_i),
$$

which vanishes for the universal choice, i.e. when all the soft scalar masses are the same at the unification point. Surprisingly, it vanishes in the models considered here too.

2. $SU(5)$ Finite Unified Theories

A realistic two-loop finite $SU(5)$ model was presented in [22], and shortly afterwards the conditions for finiteness in the soft susy breaking sector at one-loop [11] were given. Since these finite models have usually an extended Higgs sector, in order to make them viable a rotation of the Higgs sector was proposed [23]. The first all-loop finite theory was studied in [6, 7], without taking into account the soft breaking terms. Naturally, the concept of finiteness was extended to the soft breaking sector, where also one-loop finiteness implies two-loop finiteness [24], and then finiteness to all-loops in the soft sector of realistic models was studied [25, 26], although the universality of the soft breaking terms lead to a charged LSP. This fact was also noticed in [27], where the inclusion of an extra parameter in the Higgs sector was introduced to alleviate it. With the derivation of the sum-rule in the soft supersymmetry breaking sector and the proof
that it can be made all-loop finite the construction of all-loop phenomenologically viable finite models was made possible [4, 21].

Here we will examine such all-loop Finite Unified theories with $SU(5)$ gauge group, where the reduction of couplings has been applied to the third generation of quarks and leptons. An extension to three families, and the generation of quark mixing angles and masses in Finite Unified Theories has been addressed in [28], where several examples are given. These extensions are not considered here. Realistic Finite Unified Theories based on product gauge groups, where the finiteness implies three generations of matter, have also been studied [29].

The particle content of the models we will study consists of the following supermultiplets: three ($\Phi + \Omega$), needed for each of the three generations of quarks and leptons, four ($\Phi + \Phi$) and one $24$ considered as Higgs supermultiplets. When the gauge group of the finite GUT is broken the theory is no longer finite, and we will assume that we are left with the MSSM.

Thus, a predictive Gauge-Yukawa unified $SU(5)$ model which is finite to all orders, in addition to the requirements mentioned already, should also have the following properties:

(i) One-loop anomalous dimensions are diagonal, i.e., $\gamma^{(1)}_{i,j} \propto \delta_{i,j}$.
(ii) Three fermion generations, in the irreducible representations $\Phi_i, \Omega_i$ ($i = 1, 2, 3$), which obviously should not couple to the adjoint $24$.
(iii) The two Higgs doublets of the MSSM should mostly be made out of a pair of Higgs quintet and anti-quintet, which couple to the third generation.

In the following we discuss two versions of the all-order finite model. The model of ref. [6, 7], which will be labeled A, and a slight variation of this model (labeled B), which can also be obtained from the class of the models suggested in ref. [25] with a modification to suppress non-diagonal anomalous dimensions.

The superpotential which describes the two models takes the form [4, 6, 7]

$$W = \sum_{i=1}^{3} \left[ \frac{1}{2} g_{u}^{i} \Phi_{i} \Omega_{i} + g_{d}^{i} \Phi_{i} \bar{\Omega}_{i} \right] + g_{23}^{d} \Phi_{2} \Omega_{3} + g_{23}^{d} \bar{\Phi}_{2} \bar{\Omega}_{3} + g_{32}^{d} \Phi_{3} \bar{\Omega}_{4} + g_{32}^{d} \bar{\Phi}_{3} \bar{\Omega}_{4} + \frac{4}{3} g_{a}^{f} H_{a} 24 \bar{H}_{a} + \frac{g_{3}^{\lambda}}{3} (24)^{3}, \quad (7)$$

where $H_{a}$ and $\bar{H}_{a}$ ($a = 1, \ldots, 4$) stand for the Higgs quintets and anti-quintets.

We will investigate two realizations of the model, labelled A and B. The main difference between model A and model B is that two pairs of Higgs quintets and anti-quintets couple to the $24$ in B, so that it is not necessary to mix them with $H_{4}$ and $\bar{H}_{4}$ in order to achieve the triplet-doublet splitting after the symmetry breaking of $SU(5)$ [4]. Thus, although the particle content is the same, the solutions to the finiteness equations and the sum rules are different, which will reflect in the phenomenology, as we will see.

2.1. FUTA

After the reduction of couplings the symmetry of the superpotential $W$ (7) is enhanced. For model A one finds that the superpotential has the $Z_{7} \times Z_{3} \times Z_{2}$ discrete symmetry with the charge assignment as shown in Table 1, and with the following superpotential

$$W = \sum_{i=1}^{3} \left[ \frac{1}{2} g_{u}^{i} \Phi_{i} \Omega_{i} + g_{d}^{i} \Phi_{i} \bar{\Omega}_{i} \right] + g_{4}^{f} H_{4} 24 \bar{H}_{4} + \frac{g_{3}^{\lambda}}{3} (24)^{3}, \quad (8)$$

3
The non-degenerate and isolated solutions to $\gamma_i^{(1)} = 0$ for model FUTA, which are the boundary conditions for the Yukawa couplings at the GUT scale, are:

\[
(g_1^u)^2 = \frac{8}{5} g^2, \quad (g_1^d)^2 = \frac{6}{5} g^2, \quad (g_2^u)^2 = (g_3^u)^2 = \frac{8}{5} g^2, \\
(g_2^d)^2 = (g_3^d)^2 = \frac{6}{5} g^2, \quad (g_2^u)^2 = 0, \quad (g_2^d)^2 = (g_3^d)^2 = 0, \\
(g_1^l)^2 = \frac{15}{7} g^2, \quad (g_1^l)^2 = (g_3^l)^2 = 0, \quad (g_1^l)^2 = g^2.
\]

In the dimensionful sector, the sum rule gives us the following boundary conditions at the GUT scale for this model [4,22,23]:

\[
m_{H_u}^2 + 2m_{10}^2 = m_{H_d}^2 + m_{\overline{5}}^2 + m_{10}^2 = M^2,
\]

and thus we are left with only three free parameters, namely $m_{\overline{5}} \equiv m_{\overline{5}_3}$, $m_{10} \equiv m_{10_3}$ and $M$.

### 2.2. FUTB

Also in the case of FUTB the symmetry is enhanced after the reduction of couplings. The superpotential has now a $Z_4 \times Z_4 \times Z_4$ symmetry with charges as shown in Table 2 and with the following superpotential

\[
W = \sum_{i=1}^{3} \left[ \frac{1}{2} g_i^u 10_i 10_i H_i + g_i^d 10_i \overline{5}_i \overline{H}_i \right] + g_{23}^u 10_2 10_3 H_4 \\
+ g_{23}^d 10_2 \overline{5}_3 \overline{H}_4 + g_{32}^d 10_3 \overline{5}_2 \overline{H}_4 + g_{12}^f H_2 24 \overline{H}_2 + g_{12}^f H_3 24 \overline{H}_3 + \frac{g_{3}^l}{3} (24)^3.
\]

For this model the non-degenerate and isolated solutions to $\gamma_i^{(1)} = 0$ give us:

\[
(g_1^u)^2 = \frac{8}{5} g^2, \quad (g_1^d)^2 = \frac{6}{5} g^2, \quad (g_2^u)^2 = (g_3^u)^2 = \frac{4}{5} g^2, \\
(g_2^d)^2 = (g_3^d)^2 = \frac{3}{5} g^2, \quad (g_2^u)^2 = (g_3^u)^2 = \frac{4}{5} g^2, \quad (g_2^d)^2 = (g_3^d)^2 = \frac{3}{5} g^2, \\
(g_1^l)^2 = \frac{15}{7} g^2, \quad (g_2^l)^2 = (g_3^l)^2 = \frac{1}{2} g^2, \quad (g_1^l)^2 = 0, \quad (g_1^l)^2 = 0,
\]

and from the sum rule we obtain:

\[
m_{H_u}^2 + 2m_{10}^2 = M^2, \quad m_{H_d}^2 - 2m_{10}^2 = -\frac{M^2}{3}, \\
m_{\overline{5}}^2 + 3m_{10}^2 = \frac{4M^2}{3}.
\]
i.e., in this case we have only two free parameters \( m_{10} \equiv m_{10_3} \) and \( M \) for the dimensionful sector.

### Table 2. Charges of the \( Z_4 \times Z_4 \times Z_4 \) symmetry for Model FUTB.

| 5_1 | 5_2 | 5_3 | 10_1 | 10_2 | 10_3 | H_1 | H_2 | H_3 | H_4 | \( \overline{H}_1 \) | \( \overline{H}_2 \) | \( \overline{H}_3 \) | \( \overline{H}_4 \) | 24 |
|------|------|------|------|------|------|------|------|------|------|------------|----------|----------|----------|------|
| Z_4  | 1    | 0    | 0    | 1    | 0    | 0    | 2    | 0    | 0    | 0          | -2       | 0        | 0        | 0      |
| Z_4  | 0    | 1    | 0    | 0    | 1    | 0    | 2    | 0    | 3    | 0          | -2       | -3       | 0        | 0      |
| Z_4  | 0    | 0    | 1    | 0    | 0    | 1    | 0    | 0    | 2    | 3          | 0        | -2       | -3       | 0      |

As already mentioned, after the \( SU(5) \) gauge symmetry breaking we assume we have the MSSM, i.e. only two Higgs doublets. This can be achieved by introducing appropriate mass terms that allow to perform a rotation of the Higgs sector \([6, 7, 22, 23, 30]\), in such a way that only one pair of Higgs doublets, coupled mostly to the third family, remains light and acquire vacuum expectation values. To avoid fast proton decay the usual fine tuning to achieve doublet-triplet splitting is performed. Notice that, although similar, the mechanism is not identical to minimal \( SU(5) \), since we have an extended Higgs sector.

Thus, after the gauge symmetry of the GUT theory is broken we are left with the MSSM, with the boundary conditions for the third family given by the finiteness conditions, while the other two families are not restricted.

We will now examine the phenomenology of such all-loop Finite Unified theories with \( SU(5) \) gauge group and, for the reasons expressed above, we will concentrate only on the third generation of quarks and leptons. An extension to three families, and the generation of quark mixing angles and masses in Finite Unified Theories has been addressed in \([28]\), where several examples are given. These extensions are not considered here. Realistic Finite Unified Theories based on product gauge groups, where the finiteness implies three generations of matter, have also been studied \([29]\).

### 3. Predictions of Low Energy Parameters

Since the gauge symmetry is spontaneously broken below \( M_{\text{GUT}} \), the finiteness conditions do not restrict the renormalization properties at low energies, and all it remains are boundary conditions on the gauge and Yukawa couplings \((9)\) or \((12)\), the \( h = -MC \) \((4)\) relation, and the soft scalar-mass sum rule at \( M_{\text{GUT}} \), as applied in the two models, Eq. \((10)\) or \((13)\). Thus we examine the evolution of these parameters according to their RGEs up to two-loops for dimensionless parameters and at one-loop for dimensionful ones with the relevant boundary conditions. Below \( M_{\text{GUT}} \) their evolution is assumed to be governed by the MSSM. We further assume a unique supersymmetry breaking scale \( M_s \) (which we define as the geometric mean of the stop masses) and therefore below that scale the effective theory is just the SM.

We now present the comparison of the predictions of the two models (\textbf{FUTA}, \textbf{FUTB}) with the experimental data, starting with the heavy quark masses see ref. \([10]\) for more details. For the top quark pole mass we used the experimental value \( M_{\text{top}}^{\exp} = (170.9 \pm 1.8) \) GeV \([31]\). Although the current experimental value is \( 171.2 \pm 2.1 \) \([32]\), it does not affect this analysis and our conclusions. For the bottom quark mass we used the running mass evaluated at \( M_Z \) \( m_{\text{bot}}(M_Z) = 2.82 \pm 0.07 \) \([33]\) to avoid the uncertainties from the running of \( M_Z \) to the \( m_b \) pole mass, which are not related to the predictions of the FUT models.
In fig.1 we show the **FUTA** and **FUTB** predictions for $M_{\text{top}}$ and $m_{\text{bot}}(M_Z)$ as a function of the unified gaugino mass $M$, for the two cases $\mu < 0$ and $\mu > 0$. In the value of the bottom mass $m_{\text{bot}}$, we have included the corrections coming from bottom squark-gluino loops and top squark-chargino loops [34], known usually as the $\Delta f$ effects. The bounds on the $m_{\text{bot}}(M_Z)$ and the $M_{\text{top}}$ mass clearly single out **FUTB** with $\mu < 0$, as the solution most compatible with this experimental constraints. Although $\mu < 0$ is already challenged by present data of the anomalous magnetic moment of the muon $a_\mu$, a heavy SUSY spectrum as the one we have here gives results for $a_\mu$ very close to the SM result, and thus cannot be excluded on this fact alone.

In addition the value of $\tan \beta$ is found to be $\tan \beta \sim 54$ and $\sim 48$ for models A and B, respectively. Thus the comparison of the model predictions with the experimental data is survived only by **FUTB** with $\mu < 0$.

We now analyze the impact of further low-energy observables on the model **FUTB** with $\mu < 0$. As additional constraints we consider the following observables: the rare $b$ decays $\text{BR}(b \to s\gamma)$ and $\text{BR}(B_s \to \mu^+ \mu^-)$, the lightest Higgs boson mass as well as a loose CDM constraint, assuming

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**Figure 1.** The bottom quark mass at the $Z$ boson scale (upper) and top quark pole mass (lower plot) are shown as function of $M$ for both models.
Figure 2. The lightest Higgs mass, $M_h$, as function of $M$ for the model FUTB with $\mu < 0$, see text.

it consists mainly of neutralinos. More details and a complete set of references can be found in ref. [10].

For the branching ratio $BR(b \to s\gamma)$, we take a value given by the Heavy Flavour Averaging Group (HFAG) is [35]

$$BR(b \to s\gamma) = (3.55 \pm 0.24^{+0.09}_{-0.10} \pm 0.03) \times 10^{-4}. \quad (14)$$

For the branching ratio $BR(B_s \to \mu^+\mu^-)$, the SM prediction is at the level of $10^{-9}$, while the present experimental upper limit from the Tevatron is $5.8 \times 10^{-8}$ at the 95\% C.L. [36], providing the possibility for the MSSM to dominate the SM contribution.

Concerning the lightest Higgs boson mass, $M_h$, the SM bound of 114.4 GeV [37] can be used. For the prediction we use the code FeynHiggs [38–41].

The lightest supersymmetric particle (LSP) is an excellent candidate for cold dark matter (CDM) [42], with a density that falls naturally within the range

$$0.094 < \Omega_{\text{CDM}} h^2 < 0.129 \quad (15)$$

favoured by a joint analysis of WMAP and other astrophysical and cosmological data [43]. Assuming that the cold dark matter is composed predominantly of LSPs, the determination of $\Omega_{\text{CDM}} h^2$ imposes very strong constraints on the MSSM parameter space, and we find that no FUT model points fulfill the strict bound of (15). On the other hand, many model parameters would yield a very large value of $\Omega_{\text{CDM}}$. It should be kept in mind that somewhat larger values might be allowed due to possible uncertainties in the determination of the SUSY spectrum (as they might arise at large $\tan \beta$, see below). Therefore, in order to get an impression of the possible impact of the CDM abundance on the collider phenomenology in our model, we will analyze the case that the LSP does contribute to the CDM density, and apply a more loose bound of

$$\Omega_{\text{CDM}} h^2 < 0.3. \quad (16)$$

Notice that lower values than the ones permitted by (15) are naturally allowed if another particle than the lightest neutralino constitutes CDM. For our evaluation we have used the code micrOMEGAs [44].

The prediction for $M_h$ of FUTB with $\mu < 0$ is shown in Fig. 2. The constraints from the two $B$ physics observables are taken into account. In addition the CDM constraint (evaluated with
Figure 3. The lightest Higgs mass, $M_h$, plotted against $M$ and the LSP, which can be the neutralino $\chi^0$ (red crosses) or the stau $\tilde{\tau}$ (blue squares), for the model FUTB with $\mu < 0$, see text.

microOMEGAs [44]) is fulfilled for the darker (red) points in the plot, see ref. [10] for details. The lightest Higgs mass ranges in

$$M_h \sim 121 - 126 \text{ GeV},$$

where the uncertainty comes from variations of the soft scalar masses, and from finite (i.e. not logarithmically divergent) corrections in changing renormalization scheme. To this value one has to add ±3 GeV coming from unknown higher order corrections [40]. We have also included a small variation, due to threshold corrections at the GUT scale, of up to 5% of the FUT boundary conditions. Thus, taking into account the $B$ physics constraints (and possibly the CDM constraints) results naturally in a light Higgs boson that fulfills the LEP bounds [37].

In Fig. 3 we present the Higgs mass for FUTB for the case when the LSP is the neutralino $\chi^0$ (red crosses) and when it is the stau $\tilde{\tau}$ (blue squares), for the range of values of the gaugino mass $M$ where the loose CDM constraint is fulfilled (left part of Fig. 2). From Fig. 3 it is clear that the prediction for the Higgs mass lies in the same range for both cases. Notice that in case the LSP is the stau it can decay by introducing bilinear R-parity violating terms, which respect the finiteness conditions. This R-parity violating terms allow us to introduce neutrino masses [45], which make them an attractive possibility. R-parity violation would have a small impact on the collider phenomenology presented here, but would remove the CDM bound (15) completely and the LSP would not be the CDM candidate, but a gravitino or an axino [46–48] could play that role.

In the same way the whole SUSY particle spectrum can be derived. The resulting SUSY masses for FUTB with $\mu < 0$ are rather large. The lightest SUSY particle starts around 500 GeV, with the rest of the spectrum being very heavy.

In Fig. 4 we plot the mass of the lightest observable SUSY particle (LOSP) as function of $M$, that comply with the $B$ physics constraints, as explained above. The darker (red) points fulfill in addition the loose CDM constraint 16. The LOSP is either the light scalar $\tau$ or the second lightest neutralino (which is close in mass with the lightest chargino).

In Fig. 5 we show the masses of various colored particles: $m_{\tilde{t}_1}$, $m_{\tilde{b}_1}$ and $m_{\tilde{g}}$. The masses show a nearly linear dependence on $M$. Assuming a discovery reach of $\sim 2.5$ TeV yields a coverage up to $M \sim 2$ TeV. This corresponds to the largest part of the CDM favored parameter space. The
Figure 4. The mass of the LOSP is presented as a function of $M$. Shown are only points that fulfill the $B$ physics constraints. The dark (red) dots in addition also satisfy the loose CDM constraint of 16.

Figure 5. The mass of various colored particles are presented as a function of $M$. Shown are only points that fulfill the $B$ physics constraints, the black ones satisfy also the loose CDM constraint.

observation of SUSY particles at the LHC or the ILC will only be possible in very favorable parts of the parameter space. For most parameter combinations only a SM-like light Higgs boson in the range of eq. (17) can be observed.

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