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Running Into New Territory in SUSY Parameter Space

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Abstract: The LEP-II bound on the light Higgs mass rules out the vast majority of parameter space left to the Minimal Supersymmetric Standard Model (MSSM) with weak-scale soft-masses. This suggests the importance of exploring extensions of the MSSM with non-minimal Higgs physics. In this article, we explore a theory with an additional singlet superfield and an extended gauge sector. The theory has a number of novel features compared to both the MSSM and Next-to-MSSM, including easily realizing a light CP-even Higgs mass consistent with LEP-II limits, $\tan\beta \lesssim 1$, and a lightest Higgs which is charged. These features are achieved while remaining consistent with perturbative unification and without large stop-masses. Discovery modes at the Tevatron and LHC are discussed.

Keywords: afr, gut, hig, suy.

1. Introduction

The minimal supersymmetric standard model (MSSM) is the most-studied theory of new physics at the weak-scale. Nevertheless, extensions to the MSSM are profligate; considerable effort has been applied, for example, to address shortcomings in the Higgs sector $(H, \bar{H})$ of the MSSM. In particular, the MSSM allows for a superpotential mass term $\mu H\bar{H}$ and fixes the quartic-coupling to the size of the (relatively small) electroweak gauge couplings. The unexplained coincidence of scale between μ and the soft-masses is the μ-problem. The small size of the quartic predicts at tree-level a lightest CP-even Higgs state, h^0, with mass less than M_Z—a result ruled out by searches at LEP-II [1].

Additional quantum corrections from the top sector can raise m_{h^0} to $\sim$130 GeV—though only in the case of large $\tan\beta$ (the ratio of the vacuum expectation values (VEV's) of the two Higgses), TeV stop masses, and a maximal stop mixing angle [2]. Although heavy scalar tops are not precluded by any experiment, they dominate the fine-tuning in the electroweak sector through 1-loop contributions to the Higgs doublet soft-masses; the cost of opening up parameter space in the MSSM is to reintroduce fine-tuning into the Higgs potential.

A possible remedy to both of these short-comings is the Next-to-Minimal Supersymmetric Standard Model (NMSSM) [3], in which a gauge singlet superfield (S) with superpotential interaction $\lambda_S S H\bar{H}$ is added to the MSSM. The usual μ term is prevented by imposing a Z_3 symmetry. S receives a VEV proportional to its soft-mass, leading to the desired soft-scale μ term $\lambda_S\langle S\rangle H\bar{H}$. In this case, the admixture of the singlet into the electroweak symmetry breaking sector can ease experimental constraints even though light CP-even states still exist [4]. Alternatively, as noted in [5, 6], if S has a large soft-mass and is integrated out in the non-supersymmetric limit, then the effective theory contains a new quartic-coupling $|\lambda_S|^2|H\bar{H}|^2$ which independently lifts the lightest CP-even state at tree level—irrespective of the stop masses. Although one gives up using S as a solution to the μ-problem, one can pursue additional options such as the Giudice-Masiero mechanism [7].

The models described in [6] which attempt to raise the CP-even states without heavy stops are limited in their effect on the lightest Higgs mass by requiring perturbativity of all couplings up the GUT scale. The issue is that every component of the Higgs quartic coupling is infrared free. The requirement that λ_S remains perturbative to the GUT scale forces one to choose $\lambda_S \lesssim 0.6$ at the weak scale, corresponding to $m_{h^0} \lesssim 160$ GeV. Significantly larger values of λ_S (and hence m_{h^0}) can be had at the weak-scale if one gives up standard perturbative unification and simply cuts-off the (now assumed) effective theory at the scale where λ_S hits a Landau pole [8].

In [9] these perturbative constraints led us to extend the MSSM gauge sector by a non-Abelian gauge group which adds an asymptotically-free contribution to the quartic coupling. The extended gauge structure allowed us to consider significantly larger quartics at the weak scale without spoiling GUT unification. After imposing constraints from precision electroweak observables and requiring no fine-tuning in the Higgs sector, the resultant increase in the lightest CP-even mass bound was dramatic: $m_{h^0} \lesssim 350$ GeV. Similar bounds

were found recently in models which use hard effects from low-scale SUSY breaking [10], and in models which interpret the Landau pole in the quartic as a compositeness scale [11].

In this article we explore an alternative which combines the benefits of the MSSM-plus-singlet theories with those of the gauge-extension models. We raise m_{h^0} by including a singlet coupling $\lambda_S S H \bar{H}$, and as in [6], we integrate S out by giving it a relatively large soft-mass ($\sim$ 1 TeV). In the absence of further ingredients, this would leave us with the limited increase in m_{h^0} corresponding to $\lambda_S \lesssim 0.6$ from the perturbative unification bound mentioned above. The bound is removed when we include an asymptotically-free $SU(2)$ interaction which counteracts the tendency of λ_S to drive itself large at high scales. Atypical regions of supersymmetric parameter space, specifically those with small $\tan\beta$ & $m_{H^+} < m_t, m_{h^0}$, are motivated by our construction. Regions with light m_{H^+} have previously been identified in other MSSM extensions [12], and such regions might prove more generic as the exploration of models beyond-the-MSSM continues.

The resulting model (Section 2) is structurally similar to the that of Ref. [9], although we focus on the effect of the new singlet interaction and do not take advantage of the additional D-term contributions to the Higgs quartic as was done in [9]. Section 3 describes how the extended gauge sector helps to keep both λ_S and the top Yukawa interaction y_t (including regions where $\tan\beta < 1$) under perturbative control and consistent with GUT-scale gauge-coupling unification. Readers interested in the generic properties of the allowed parameter space should skip to Section 4, where we find not only that the upper bound on the lightest CP-even state is 250 GeV, but also that the lightest state in the Higgs spectrum can be $H^\pm$ and discuss the novel phenomenology. Section 5 contains our conclusions.

2. The Gauge-Extended MSSM + Singlet

Our theory has gauge group $SU(3)_c \times SU(2)_1 \times SU(2)_2 \times U(1)_Y$ with couplings $g_3, g_1, g_2,$ and g_Y, respectively. The non-universal charge assignment of the model is similar to that of Topflavor or "heavy" extended technicolor models [13] which have $SU(2)_1$ as the weak-group for the Higgs fields $(H, \bar{H})$ and third generation of fermions, but $SU(2)_2$ as the weak-group for the first two generations of fermions. To obtain the correct low energy quantum numbers after $SU(2)_1 \times SU(2)_2 \to SU(2)_D$, the third generation and Higgs fields are not charged under $SU(2)_2$, and the first two generations are not charged under $SU(2)_1$.

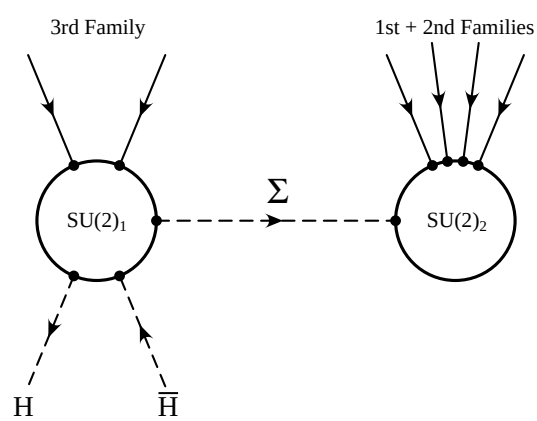

Figure 1: SU(2) structure

To break the $SU(2)_1 \times SU(2)_2$ to the diagonal subgroup we include a bi-doublet chiral field Σ which transforms as a $(2, \bar{2})$ and gets a VEV (u) from superpotential interactions with a singlet field, S_Σ. This is depicted schematically in Figure 1. In [9], it was shown that

the quartic potential for the Higgs is enhanced by the additional $SU(2)$ D-terms when the diagonal symmetry breaking occurs at a scale significantly smaller than m_Σ, the soft-mass for Σ. For this paper, we assume that $SU(2)_1 \times SU(2)_2$ breaks to the diagonal $SU(2)$ close to the supersymmetric limit, and the enhancement of [9] is negligible. Precision electroweak constraints were examined in [9] and require that the breaking scale, u, for $SU(2)_1 \times SU(2)_2 \to SU(2)_D$ must be $\gtrsim 2.5$ TeV. We will thus consider $u = 3$ TeV as an example choice for the symmetry-breaking scale for the remainder of this work. The interactions and fields needed for Yukawa couplings and diagonal $SU(5)$ unification are described below.

To avoid the LEP-II bound without a need for extremely heavy stops, we instead enhance the quartic coupling by including a singlet, S, with superpotential,

$$W = \mu H\bar{H} + \lambda_S S H\bar{H} + \frac{\mu_S}{2} S^2, \tag{2.1}$$

where we self-consistently ignore linear and trilinear terms. We also include a soft-mass for S, m_S^2. We here ignore the effects of both a soft-trilinear term $A_S S H\bar{H}$ and a Supergravity generated soft-tadpole $t_s S$. The effect below m_S of A_S would be to correct the effective $|\lambda_S|^2$ in Eq. 2.2 by $|A_S|^2/m_S^2$; the soft-tadpole would induce a VEV for S of order t_s/m_S^2 and could be controlled through low-scale SUSY mediation [14] or by the appearance of a new gauge group (under which S is charged) at an intermediate scale.

If both μ_S & $m_S \sim v = 174$ GeV, the theory is similar to the NMSSM, where S must be included in the Higgs potential and a VEV for S could possibly replace the conventional MSSM μ term. This is an interesting case, but it typically does not realize very large increases in the Higgs mass and thus is tightly constrained by the LEP-II bound (though small regions of parameter space with in which the light state is mostly singlet and thus has suppressed couplings to the Z survive) [4]. For $\mu_S \gg m_S \& v$, S decouples supersymmetrically, and the low energy theory is simply the MSSM, including the Higgs sector. We focus on the third possibility, in which S is integrated out in the non-supersymmetric limit, where $m_s \gg v$ and $\mu_S^2/(m_s^2+\mu_s^2) \ll 1$. In practice, $\mu_S \sim 100$ GeV and $m_S \sim 1$ TeV suffice. This is the non-decoupling limit described in [6] and the limit pursued throughout the rest of this paper.

In this limit there are no singlet Higgses in the low energy spectrum, and the remnant effect is in the quartic potential for H and $\bar{H}$:

$$\frac{g^2}{8}\left(H^\dagger \vec{\sigma} H - \overline{H}\vec{\sigma}\overline{H}^\dagger\right)^2 + \frac{g_Y^2}{8}\left(|\overline{H}|^2 - |H|^2\right)^2 + |\lambda_S|^2 \left|H\bar{H}\right|^2, \tag{2.2}$$

where $1/g^2 = 1/g_1^2 + 1/g_2^2$ is the low energy $SU(2)$ coupling of the standard electroweak theory. The first two terms are the ordinary MSSM D-terms for the $SU(2) \times U(1)$ gauge interactions. The last term is the effective contribution from S when $\mu_S^2/(m_s^2+\mu_s^2) \ll 1$. The Higgs scalar potential also contains the usual MSSM quadratic pieces

$$\left(|\mu|^2 + m_H^2\right)|H|^2 + \left(|\mu|^2 + m_{\bar{H}}^2\right)\left|\bar{H}\right|^2 - \left(bH\bar{H} + c.c.\right) \tag{2.3}$$

which contribute to electroweak symmetry breaking. Throughout, we define $\langle H\rangle = v_H$, $\langle \bar{H}\rangle = v_{\bar{H}}$, $v_H^2 + v_{\bar{H}}^2 = v^2$, $v = 174$ GeV, and $\tan\beta = v_{\bar{H}}/v_H$.

General CP-conserving two higgs doublets models were studied in [15]. Unlike the MSSM, this Higgs potential has no flat (or negative) D-term directions. Electroweak symmetry breaking occurs so long as $b^2 > (|\mu|^2 + m_H^2)(|\mu|^2 + m_{\bar{H}}^2)$ (which insures at least one negative mass eigenvalue in the Higgs mass matrix) and $b/\sin 2\beta + M_W^2 > v^2\lambda_S^2$ (so that $H^\pm$ do not develop VEV's and $U(1)_{EM}$ remains unbroken).

2.1 Yukawa Couplings

As noted in [9], Yukawa interactions for the third generation fermions may be written down at the renormalizable level, since the Higgses and the third generation are charged under the same $SU(2)$. Yukawa couplings for the first two generations can be generated by adding a massive Higgs-like pair of doublets $\overline{H}', H'$, that are charged under $SU(2)_2$. They couple to the first two generations via Yukawa-type couplings and mix with the regular Higgses via superpotential operators such as $\lambda'\overline{H}\Sigma H'$ with $\lambda'\langle\Sigma\rangle \sim \mu$. A supersymmetric mass $\mu_{H'}H'\bar{H}'$, with $\mu_{H'} \gtrsim \langle\Sigma\rangle$, for the new doublets allows us to integrate them out and generates Yukawa couplings for the first two generations at low energies. Mixing with the third generation (*i.e.*, V_{cb} and V_{ub}) can easily be generated since the right-handed quarks have the same $SU(3)\times U(1)$ quantum numbers. The result is that the MSSM Higgses have essentially MSSM Yukawa interactions, though low $\tan\beta$ is now accessible as explained in section 3.

2.2 Grand Unification

This model can also be made consistent with gauge coupling unification. The details are relevant because they influence the β function coefficients for the gauge couplings— important when we determine the bound on λ_S by requiring perturbativity to the GUT scale in Sec. 3 below. The full group $SU(3)_c \times SU(2)_1 \times SU(2)_2 \times U(1)_Y$ can be embedded in $SU(5) \times SU(5)$ broken by a bi-fundamental field at the GUT scale with a VEV $\langle\Xi\rangle = diag\{M, M, M, 0, 0\}$ [17]. Gauge coupling unification is predicted (with theoretical uncertainty beyond one-loop) because the standard model gauge couplings are only a function of the diagonal gauge coupling. At one loop, one can track the diagonal $SU(2)$ through its beta-function coefficient as it is the sum of those of the two $SU(2)_i$. It receives an extra -6 from the additional triplet of gauge bosons. We include a bi-doublet Σ and two triplets charged under $SU(2)_2$ which altogether contribute +6 to the diagonal beta function. We must also add an additional vector-like pair of triplets to effectively complete a 5 and $\overline{5}$ with the extra pair of Higgs-like fields (H' and $\bar{H}'$) required for Yukawa interactions for the first two generations. With these additions, both $SU(2)$ models achieve the same level of unification as in the MSSM at one loop.

2.3 Supersymmetry Breaking

Although our framework is independent of the mechanism of supersymmetry breaking, we do ask that the soft-mass for the singlet be somewhat larger than the soft-masses for the MSSM fields. In gauge mediation, supersymmetry is broken in a hidden sector and communicated to the MSSM at one-loop (for gaugino masses) or two-loops (for scalar masses) through messenger fields. The singlet can be coupled directly to SUSY breaking

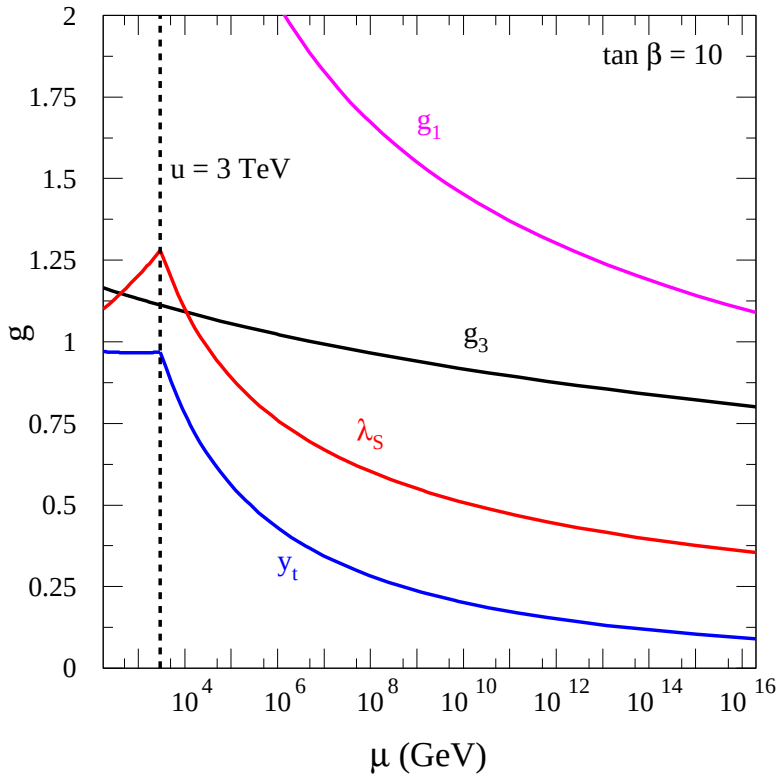

Figure 2: Renormalization group flow of couplings for the model, with $g_1(3\text{ TeV}) = 2$, $\lambda_S(200\text{ GeV}) = 0.8$ and $\tan\beta = 10$.

messengers $M, \bar{M}$ via the superpotential coupling $SM\bar{M}$. The result is a one-loop (instead of two loop) squared soft-mass for S which is roughly $4\pi/\alpha$ larger than the squared MSSM soft-masses, as desired.

Another option which makes use of extra dimensions is gaugino mediation [18]: SUSY breaks on a sequestered brane, couples directly to all bulk fields (including the gauginos), and then communicates to the MSSM matter on a visible brane through the gauginos. Gaugino masses arise from a superpotential term of the form $X\mathcal{W}_\alpha\mathcal{W}^\alpha$, while bulk scalar fields receive *squared* masses from the Kähler term $X^\dagger X S^\dagger S$. Therefore, if we put the singlet in the bulk, its mass receives an enhancement of $\sqrt{ML}$ relative to the gaugino masses. Here, L is the length of the extra-dimension, while M is some fundamental scale. Rough bounds from flavor constraints and naive dimensional analysis predict $10 \lesssim ML \lesssim 100$, resulting in the proper enhancement for the soft-mass of S.

3. Perturbativity Constraints

The enhanced quartic effect in the non-decoupling limit is strongly limited by a desire for

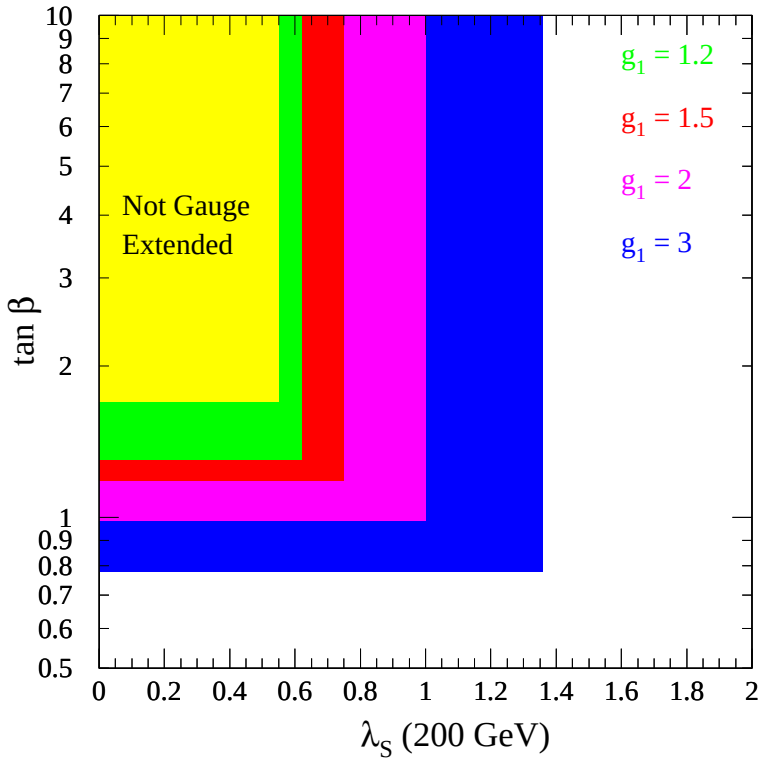

Figure 3: Allowed regions of $\tan\beta$ and $\lambda_S(200\text{ GeV})$ for given values of the new $SU(2)$ coupling $g_1(3\text{ TeV})$. For the indicated regions, all couplings in the model remain perturbative during 1-loop RG evolution up to the GUT scale. From left to right, the regions are: the non-gauge extended model of [6], and our gauge extension with $g_1(u)$ =1.2, 1.5, 2, and 3.

perturbative unification. Both couplings λ_S and y_t feed into each other's renormalization group equations (RGE's) with positive coefficients. If either λ_S or y_t is large at a low scale (required for $m_{h^0} > M_Z$, or low $\tan\beta$, respectively), non-perturbative physics is reached long before M_{GUT}.

Both of these problems are largely ameliorated by the presence of new, relatively strong, gauge interactions, which drive both y_t and λ_S down at large scales, owing to the Higgs and top quark participation in the stronger group. The dominant terms in the renormalization group equations at one loop (including the spectator matter described above, necessary for gauge coupling unification and Yukawa interactions for the first two generations) are[1],

$$\frac{dg_1}{dt} = -2\frac{g_1^3}{16\pi^2}$$

[1]These RGE's are valid above the $SU(2) \times SU(2)$ breaking scale, u. Below u, we use the RGE's appropriate for the broken phase.

$$
\begin{aligned}
\frac{dg_2}{dt} &= 4\frac{g_2^3}{16\pi^2} \\
\frac{dg_3}{dt} &= -2\frac{g_3^3}{16\pi^2} \\
\frac{d\lambda_S}{dt} &= \frac{\lambda_S}{16\pi^2}\left(3\left|y_t\right|^2 + 4\left|\lambda_S\right|^2 - 3g_1^2\right) \\
\frac{dy_t}{dt} &= \frac{y_t}{16\pi^2}\left(6\left|y_t\right|^2 + \left|\lambda_S\right|^2 - 3g_1^2 - \frac{16}{3}g_3^2\right),
\end{aligned} \tag{3.1}
$$

with sub-dominant contributions from the bottom Yukawa (particularly at low $\tan\beta$), and g_Y. For large enough g_1, the strong gauge interactions drive y_t and λ_S smaller just above the electroweak scale. The effectiveness is limited by the fact that the additional $SU(2)$ is itself strongly asymptotically-free, and thus becomes more and more irrelevant in the RGEs at higher and higher energies. However, the result is a wider region of allowed parameter space consistent with perturbativity to the GUT scale. In Figure 2 we show a sample flow of the couplings for $\tan\beta = 10$, $g_1(u = 3\text{ TeV}) = 2$, and $\lambda_S(200\text{ GeV}) = 0.8$.

In Figure 3 we show the allowed regions of $\tan\beta$ and λ_S, for fixed values of $g_1(u = 3\text{ TeV})$, by requiring that all couplings remain perturbative up to the GUT scale. It should be noted that there is in fact a minimum value of g_1 for this theory which is compatible with perturbative unification. Because g_2 is not asymptotically-free, if its value at u is too large, it will flow strong before the GUT scale. From the one-loop RGE expression above, we see that this happens when $g_1 \leq 1.2$ at $u = 3$ TeV.

4. Higgs Properties and Phenomenology

4.1 Higgs Spectrum

We include stop corrections and find that the one-loop Higgs CP-even mass matrix can be written as a function of the CP-odd mass (m_A), the μ parameter, the stop masses $m_{\tilde{t}_i}$, and the stop mixing parameter, A_t [19]:

$$
\begin{aligned}
m_{11}^2 &= m_Z^2\cos^2\beta + m_A^2\sin^2\beta - \frac{3y_t^2}{16\pi^2}\mu^2\frac{Z^2}{3} \\
m_{22}^2 &= m_Z^2\sin^2\beta + m_A^2\cos^2\beta + \frac{3y_t^2}{16\pi^2}\left(4m_t^2\log\frac{m_{\tilde{t}_1}m_{\tilde{t}_2}}{m_t^2} + A_t(2m_tZ - A_t\frac{Z^2}{3})\right) \\
m_{12}^2 &= -\frac{1}{2}(m_Z^2 + m_A^2 - 2v^2\lambda_S^2)\sin 2\beta + \frac{3y_t^2}{16\pi^2}\mu\left(m_tZ - A_t\frac{Z^2}{3}\right)
\end{aligned} \tag{4.1}
$$

where:

$$
Z = \frac{m_t(A_t + \mu\cot\beta)}{m_t^2 + \frac{1}{2}(m_Q^2 + m_U^2)} \tag{4.2}
$$

and the stop masses are defined with respect to the soft-masses for left- and right-handed stops (m_Q,m_U) as:

$$
\begin{aligned}
m_{\tilde{t}_{1,2}}^2 &= m_t^2 + \frac{1}{2}(m_Q^2 + m_U^2) \pm W \\
W^2 &= \frac{1}{4}(m_Q^2 - m_U^2)^2 + y_t^2v^2|A_t\sin\beta - \mu\cos\beta|^2
\end{aligned} \tag{4.3}
$$

The charged Higgs mass is (at one-loop):

$$m_{H^\pm}^2 = m_A^2 + m_W^2 - v^2 \left|\lambda_S\right|^2 . \tag{4.4}$$

In the above, the top Yukawa coupling is evaluated at the stop mass scale in order to take into account the leading effect from RGE-improvement [2]. We have neglected the subdominant corrections from the gauginos and superpartners of light fermions. The CP-even mass eigenstates and mixing angle α are obtained by diagonalizing this two by two matrix. The largest h^0 masses are obtained for large λ_S, the decoupling regime $m_A > m_{h^0}$, and $\tan\beta \sim 1$. For $\lambda_S \sim 1.5$, these parameters (with soft-masses for the stops of 200 GeV) result in $m_{h^0} \sim 260$ GeV, which is the largest h^0 mass that can be realized in our model consistent with perturbativity up to the GUT scale. Such a large value of λ_S starts to reintroduce fine-tuning into the Higgs soft-mass through the renormalization group equations if SUSY is broken close to the Planck scale. For $m_S \sim 1$ TeV and $\lambda_S \sim 1.5$ this contribution is less than 1 TeV.

In Fig. 4 we plot the spectrum for $m_Q = m_U = \mu = 200$ GeV, $\tan\beta = 1$ and $\lambda_S = 0.8$ (requiring $g_1(u) \sim 2$, see Fig. 3), as a function of m_A. It should be noted that this set of soft SUSY-breaking parameters in the MSSM would predict a lightest CP-even Higgs mass much below the LEP-II bound, and would thus be ruled out. We see that even for these rather small soft-masses for the stops (and no stop mixing at all), we can easily accommodate CP-even Higgs masses consistent with the LEP-II bounds. This eliminates the fine-tuning of the weak scale that is inevitable in the MSSM.

4.2 Collider Phenomenology

This theory has a rich and distinct Higgs phenomenology compared to the MSSM. For example, Figure 4 shows a sample Higgs spectrum for $\tan\beta = 1$. Figure 5 indicates the region of parameter space where the lightest state is the *charged higgs* corresponding to the case when the heavier Higgs is responsible for electroweak symmetry-breaking.

The LEP-II limits on the charged Higgs mass (from direct searches such as $e^+e^- \to H^+H^-$) is 81 GeV [20] and requires $m_A > 140$ GeV if $\lambda_S = 0.8$ and $\tan\beta = 1$. In this regime the light CP-even Higgs has small couplings to gauge bosons and may avoid LEP-II bounds even for $m_h \leq 115$ GeV. Indirect constraints on m_{H^+} from the observed rate of $b \to s\gamma$ (see [21] for a recent summary) are stringent even in models with minimal flavor violation and superpartners at the TeV scale. However, supersymmetric cancellations from loops involving light charginos and stops (see [22], for example) can reproduce the observed $b \to s\gamma$ rates even for $m_{H^+} \sim 80$ GeV and $\tan\beta \sim 1$ [23]. For $m_{H^+} \sim 80$ GeV and $\tan\beta \sim 1$, the m_{H^+} contribution is cancelled by $\sim 75\%$.

For much of the region of parameters in which it is the lightest, the charged Higgs mass is much less than the top mass. This implies (for low $\tan\beta$) that the charged Higgs will predominantly decay into charm and strange, with a much smaller fraction into τ and neutrino. For $m_{H^\pm} \sim m_t$ the three-body decay through an off-shell top into $Wb\bar{b}$ becomes interesting. The dominant production for this range of masses is through a rare top decay,

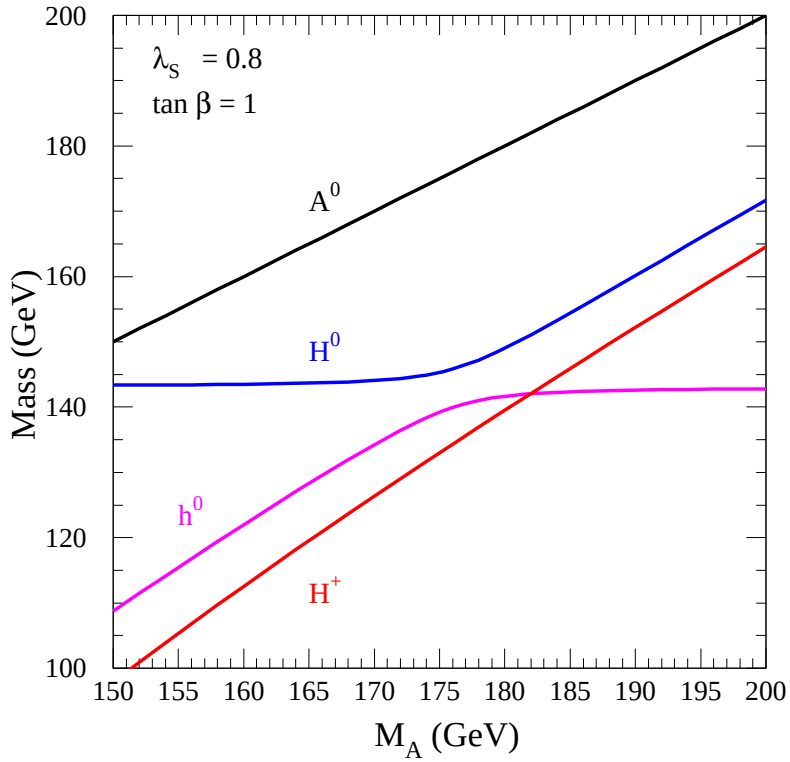

Figure 4: One loop spectrum of Higgs masses (varying m_A) for $\tan\beta = 1$ and $\lambda_S = 0.8$. The soft parameters were chosen as $m_Q = m_U = \mu = 200$ GeV and $A_t = 0$.

$t \to H^+ b$. The branching fraction is given at tree-level by,

$$\frac{\Gamma(t \to H^+ b)}{\Gamma(t \to W^+ b)} = \frac{1}{1 + 2m_W^2/m_t^2} \left(\frac{1 - m_{H^\pm}^2/m_t^2}{1 - m_W^2/m_t^2} \right)^2 \times \cot^2\beta \tag{4.5}$$

and is enhanced for $\tan\beta < 1$. One-loop SUSY QCD corrections tend to suppress the branching ratio slightly [24]. This decay can be distinguished from usual top decay first because it modifies the branching ratios into jets compared to leptons, and second because the decay products reconstruct an intermediate $H^\pm$ mass instead of the W mass. At the Tevatron run II, predictions are that the region with $\tan\beta < 1$ and $m_{H^\pm} < 120$ GeV can be discovered or excluded through top decay with 2 fb^{-1} of integrated luminosity [25]. The LHC with 100 fb^{-1} is expected to be able to to see a charged Higgs with $m_{H^\pm} < m_t - 20$ GeV for all values of $\tan\beta$ [26].

Charged Higgs masses greater than m_t are possible (continuing to assume $\lambda_S = 0.8$) if m_A is larger than 210 GeV. In this case, the dominant decay is $H^+ \to t\bar{b}$ and the dominant production at the LHC is associated production of a Higgs with a top quark, through partonic processes such as $gb \to H^- t$ [27]. This process can study charged Higgs bosons

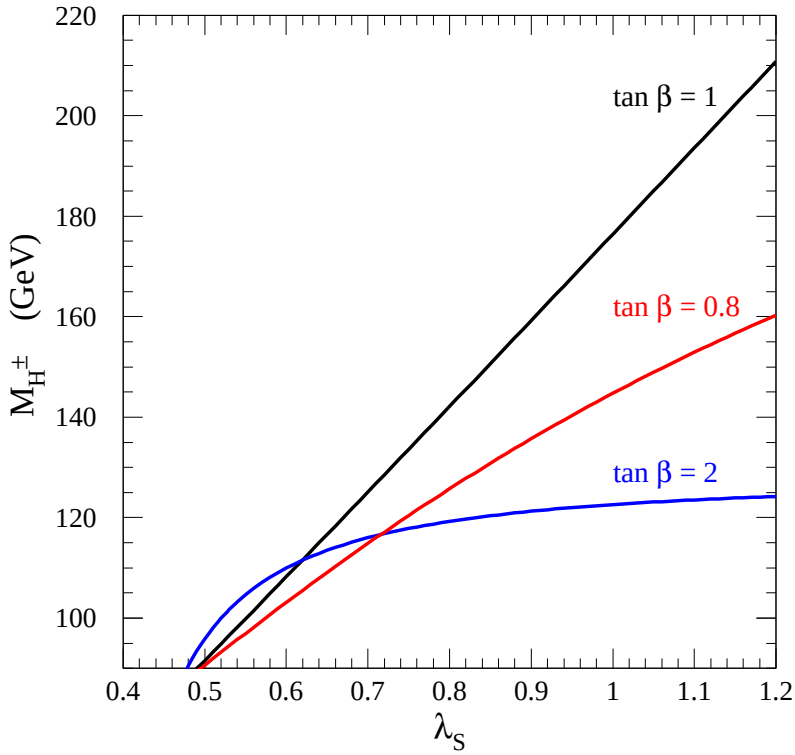

Figure 5: The charged Higgs mass versus λ_S for different values of $\tan\beta$. The area below each line represents the parameter space in which the charged Higgs is the lightest Higgs in the spectrum. The lines themselves represent the charged Higgs masses for which the charged and lightest CP-even Higgs bosons are degenerate. The soft parameters were chosen as $m_Q = m_U = \mu = 200$ GeV and $A_t = 200$ GeV, and there is moderate dependence on them. For example, if the above parameters are set at 500 GeV, the values of the lines at $\lambda = 1.2$ are (from lowest to highest $\tan\beta$) 162, 210, and 145 GeV.

with masses up to about 400 GeV in the low $\tan\beta$ region [28] with 100 fb^{-1}. Another production mechanism is through an off-shell W boson, $q\bar{q}' \to W^* \to H^+ A_0$ [29], leading to final states with $t\bar{b}b\bar{b}$. Two of the bottoms reconstruct m_{A^0}, and thus typically have much higher energies than $b\bar{b}$ from gluon splitting.

Note that despite the absence of additional weak scale Higgs bosons, this theory modifies the usual MSSM $m_{H^\pm}^2 - m_A^2 = m_W^2$ mass relation. Even when the spectrum is roughly consistent with the MSSM (say, for modest λ_S and large m_A such that m_{h^0} has mass just above the LEP bound and h^0 has largely SM-like couplings), this fact can be tested at the LHC through the associated production of $H^\pm A^0$ provided the masses are less than about 300 GeV [29]. As with any extension of the MSSM that affects the Higgs quartic, one could also combine a measurement of the light CP even Higgs mass with precision measurements

of the stop masses and mixing angle (for example, at a linear collider [30]) to show that the Higgs mass does not satisfy the MSSM relation.

The additional gauge bosons (with masses in the several TeV range) associated with the top-flavor group can be produced at the LHC. The dominant process is one in which two first generation quarks fuse into a W'. The W' coupling to light quarks is suppressed compared to third generation fermions, but this is compensated by the much larger probability to find a high-energy first generation (valence) quark inside a proton compared to the probability of finding a third generation quark. Once produced, the W' predominantly decays into $t\overline{b}$ and $\tau\overline{\nu}$, and their super-partners. The decay into $t\overline{b}$ looks like s-channel single top production [31] and has been carefully studied in Ref. [32], including NLO QCD corrections to the signal rate, estimations of detector efficiencies and backgrounds, etc. The conclusion of that study is that W' masses less than 4.5 TeV can discovered with 100 fb^{-1} of integrated luminosity.

Finally, the additional gauge fields and the Σ provide a number of heavy fermions with masses around $u \sim 3$ TeV, and the super-partner of the singlet S is typically weak scale. They will have small (but potentially important) mixings with the standard charginos and neutralinos, and strong mixing with each other. The gaugino components will also prefer to couple to the third family, and decays involving multiple bottom and top quarks (plus missing energy from the LSP) and τ leptons will dominate. The masses are probably too large for copious production at the LHC, but the influence through mixing on the lighter neutralinos and charginos could affect their production and decay in a relevant way. Precision measurements of neutralino and chargino couplings at a future e^+e^- linear collider could reveal this mixing.

5. Conclusions

The need for a Higgs mass greater than the LEP-II bound has placed the MSSM in an interesting position and motivates extensions which allow for larger Higgs masses. In this article we have presented one such model which invokes a singlet Higgs to increase the Higgs mass by effectively adding a new Higgs quartic to the superpotential. The new feature is the addition of asymptotically-free gauge interactions which tend to drive this extra quartic and the top Yukawa coupling smaller at high energies.

This allows one to explore larger quartic couplings consistent with perturbative unification than in the past [6] and predicts a new bound on the lightest CP-even state $m_{h^0} < 250$ GeV. It also opens a window of $\tan\beta < 1$ for which the top Yukawa remains perturbative all the way up to the GUT scale. The result is a theory in which the lightest CP-even Higgs may be heavier than the LEP-II bound at tree level without the need to invoke large stop masses and introduce electroweak fine-tuning. In fact, the phenomenology associated with larger quartics and lower $\tan\beta$ is more general than the specific model presented here; a theory such as the "SUSY Fat Higgs" model [11] which invokes a low scale cut-off on λ_S and the Yukawa interactions (while still remaining consistent with unification) also allows one to explore the same regions of parameter space with similar phenomenology.

Replacing the singlet with a pairs of triplets (with hypercharge ± 1) as in Ref [6] is unlikely to enhance the mass-bound further as the additional matter will make the β-function of g_1 vanish at one-loop. Further, a triplet contribution to the Higgs quartic will inevitably lead to a VEV for the triplet, which is disfavored by precision electroweak data. Instead one could imagine enhancing the bound by additionally including the mechanism of Ref. [9] by lifting the soft-mass for the gauge-breaking field to ~ 10 TeV.

The resulting phenomenology is somewhat unusual with the charged Higgs as the lightest Higgs for a range of parameters. The additional gauge bosons are also expected to be visible at the LHC and provide a tangible way in which experiments would be able to test this scenario and distinguish it from alternatives.

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