Pushing the SUSY Higgs mass towards 125 GeV with a color adjoint

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Abstract

We show that inclusion of a TeV scale chiral superfield transforming in the adjoint representation of the color SU(3) to the MSSM particle content modifies the renormalization group running of some parameters in such a way that a 125 GeV mass of the light Higgs boson is accommodated more comfortably than in cMSSM / mSUGRA. Put differently, the introduction of a color adjoint TeV scale superfield helps resurrecting lighter choices for the stop and gluino which are otherwise disfavored in cMSSM / mSUGRA.

Introduction: The discovery of the Higgs boson, or more appropriately, a Higgs-like boson at the CERN Large Hadron Collider (LHC) [1,2] has pushed one of the most advertised class of supersymmetric (SUSY) models, namely, the constrained minimal supersymmetric standard model (cMSSM), or equivalently, minimal supergravity (mSUGRA), into an uncomfortable corner. The disappointment arising from the so far unsuccessful attempts to underpin any supersymmetric relics via direct searches at the LHC [3,4] got further aggravated by the news that the Higgs boson is as heavy as 125 GeV. This is so because the third generation squark masses and the associated soft trilinear scalar coupling, on which LHC cannot put as stringent direct constraints as on their first two generation counterparts, are now pushed to one to few TeV to ensure that the lightest Higgs mass receives sufficient radiative enhancement. It is in this context that we write this short note. We propose that a simple augmentation of the MSSM particle content with a chiral superfield transforming in the adjoint representation of the SU(3) color group proves to be useful in easing part of this difficulty and thus resurrecting some of the lost parameter space. Put briefly, our scenario is the following: add a colored chiral superfield whose scalar and fermionic components behave like a scalar gluon and a gluino, respectively, with no other non-vanishing gauge quantum numbers, each weighing around a TeV. Their presence would modify the renormalization group (RG) running of the QCD coupling $g_3$ (more specifically, would add positive contributions to its beta function), which would in turn feed into the running of the top Yukawa coupling $y_t$, and the trilinear scalar coupling $A_t$ entering the stop mixing matrix. These modifications give us a few territorial advantages over cMSSM: (i) the gluino and the lighter stop can be lighter than what they should weigh in cMSSM for generating the 125 GeV mass of the light Higgs; (ii) the rather large (possibly maximal) stop-mixing, which facilitates reaching out to 125 GeV mass of the Higgs, does not compel $|A_0|$ any more to be as large as what cMSSM requires it to be. The motivation for adding an adjoint superfield may come from string theory, more specifically, the intersecting D-brane models [5]. However, such color adjoint fields in four dimensional context need not be seen only as spies from extra dimension or string theory, they may very well have more mundane ancestry. The main upshot of our analysis is that the introduction of a color adjoint superfield at the TeV scale improves the fine-tuning.

Other options for creating more room for accommodating a 125 GeV SUSY Higgs also exist. Next-to-minimal supersymmetry (NMSSM) is already known to possess an improved fine-tuning as its gauge singlet superfield, coupled to the two Higgs doublets in the superpotential, provides a tree level mass to the lightest CP-even Higgs [6,7]. A recent numerical study of the trilinear scalar couplings $A_\lambda$ and $A_\kappa$ in the conventional scale invariant version of NMSSM, however, shows that the 125 GeV mass of the Higgs boson is compatible only in some well-separated islands of the parameter space [8]. Further reduction in fine-tuning in NMSSM has more recently achieved by introducing extra matter descending from $E_6$ origin in a scenario which also possesses a discrete $R$-symmetry solving the domain wall problem and enhancing proton stability [9]. A bottom-up approach for addressing the fine-tuning problem, which goes by the name of 'natural SUSY', has also gained attention where the third generation sfermions and the Higgsino are kept light, while the rest...
of the superpartners are considered heavy [10–12]. Additional matter fields transforming as grand unified theory (GUT) multiplets have also been employed to better realize the 125 GeV mass of the Higgs, improving consistency with the muon $(g - 2)$ measurement at the same time [13, 14]. Our method of comfortably achieving the 125 GeV mass of the Higgs relies on adding a color adjoint state that does not directly couple to the Higgs sector but its effect filters through to the Higgs mass only via modifications of RG running of various couplings. Adjoint representation states have been employed for different purposes so far. It was shown that such a colored chiral superfield, appearing e.g. in the context of a 4-dimensional realization of $N = 2$ supersymmetry in a 5-dimensional theory [15, 16], helps the gluino acquire a large Dirac mass in a class of super-soft SUSY breaking models [17] actually helps to improve the fine-tuning [18] of parameters. A very recent study [19] aiming to improve the fine-tuning has employed the 24-plet SU(5)-adjoint superfield, the vacuum expectation value of whose singlet component helps to enhance the tree level Higgs mass in the NMSSM-style. As a result, the 125 GeV Higgs mass is reached with a lighter stop and smaller mixing, the colored states being used to keep the running of the top Yukawa coupling under control.

Figure 1: RG running of the top Yukawa coupling in cMSSM+ (red solid line) and in cMSSM (blue dashed line).

Figure 2: RG running of the trilinear scalar parameter $A_t$ in cMSSM+ (red solid line) and in cMSSM (blue dashed line). Here, $M_1/2 = 500$ GeV and $\tan \beta = 10$.

Formalism: We refer to the present scenario as ‘cMSSM+’ for repeated use in the subsequent text, which implies the MSSM particle content plus a single $SU(3)_C$ adjoint chiral superfield. We now demonstrate how its introduction induces a drastic modification to RG evolution of several parameters. All we need to calculate are the contributions of the fermion and scalar components of the adjoint superfield to the QCD gauge beta function. These are given by

$$\Delta b_3 = \frac{4}{3} \cdot \frac{3}{2} = 2,$$

and

$$\Delta b_3^f = \frac{1}{3} \cdot 3 = 1,$$

(1)

where the factor 3 represents color, and the factor $(1/2)$ in the fermionic contribution comes from its Majorana nature. Hence $\Delta b_3^f = \Delta b_3^f + \Delta b_3^s = 3$. We assume that the new fermion and the scalar weigh around a TeV. So, as soon as this energy is crossed, the new states are sparked into life, and the above increment in the beta function changes the slope of the running of $g_3$ keeping it flat at its weak scale value all along (up to one loop precision). This constitutes the primary effect and the rest is simply its consequence, as we explain now step by step. We recall that the gauge beta functions are given by ($t = \ln Q/(1$ GeV)),

$$\beta_{g_3} \equiv \frac{d}{dt} g_3 = \frac{b_a}{16\pi^2} g_3^3,$$

(2)

where $(b_1, b_2, b_3) = (33/5, 1, -3)$ for MSSM at one loop [20]. Since, only $b_3$ receives an increment in cMSSM+, as shown in Eq. (1), the cMSSM+ set reads: $(b_1, b_2, b_3) = (33/5, 1, 0)$. For our purpose, one loop estimate of beta functions is enough [1]. Since in cMSSM+ $g_3$ hardly runs beyond the TeV scale, the gluino mass also remains stationary at the leading order. Admittedly, gauge couplings do not unify in this model since only the slope of $g_3$ running is modified, although the value of $g_3$ remains perturbative all the way up to the high scale [1]. In our subsequent numerical discussions on cMSSM+

1 Two loop RG evolution in Dirac Gaugino context has been discussed in [21].

2 Additional chiral multiplets suitably charged under $SU(2)_L$ and $U(1)_Y$ may be added to ensure that the RG curves of $\alpha_2$ and $\alpha_3$ are also bent appropriately to reinstate gauge coupling unification at a value higher (still perturbative) than in MSSM. But this is not our main focus and we do not pursue the unification issue any further. For a discussion on gauge coupling unification in F-theory GUT models with Dirac gauginos, see [22].

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we treat $M_G = 2 \cdot 10^{16}$ GeV as a high scale reference point (for comparison of various running vis-à-vis MSSM), and assume that the common gaugino and scalar supersymmetry breaking mass parameters, $M_{1/2}$ and $m_{0}$, respectively, refer to that point. We now look at the RG running of the top Yukawa coupling, where for illustration we display only the dominant terms:

$$\beta_{y_t} \equiv \frac{d}{dt} y_t \simeq \frac{y_t}{16\pi^2} \left[ 6y_t^2 y_t - \frac{16}{3} g_3^2 \right].$$  \hspace{1cm} (3)

Since $g_3$ stays at the large weak scale value even at the high scale, the RG trajectory of the top Yukawa coupling is bent to lower values compared to MSSM at the high scale – see Fig. 1. This will help us understand the evolution pattern of the trilinear coupling $A_t$. Again, we display the dominant terms for providing intuition:

$$16\pi^2 \frac{d}{dt} A_t \simeq A_t \left[ 18y_t^2 y_t - \frac{16}{3} g_3^2 \right] + \frac{32}{3} y_t g_3^2 M_3.$$  \hspace{1cm} (4)

An interplay of Eqs. (2), (3) and (4) provides the insight that starting from a given negative high scale value $A_0$, the weak scale value $A_t$ is more negative in cMSSM+ compared to cMSSM – see Fig. 2 (drawn for $M_{1/2} = 500$ GeV and $\tan \beta = 10$). It is now known that a negative $A_t$ of larger magnitude is more helpful for reaching out to 125 GeV mass of the Higgs (see e.g. [23][24] for recent studies). This transpires from

$$m_h^2 = M_Z^2 \cos^2 2\beta + \frac{3}{4\pi^2} v^2 \left[ \log \frac{M_S^2}{m_t^2} + \frac{X_t^2}{M_S^2} \left( 1 - \frac{X_t^2}{12M_S^2} \right) \right],$$  \hspace{1cm} (5)

where $v = \sqrt{v_u^2 + v_d^2} = 174$ GeV, $\tan \beta = v_u/v_d$, $X_t \equiv A_t - \mu \cot \beta$, and $M_S \equiv \sqrt{m_{\tilde{t}_1} m_{\tilde{t}_2}}$ is the geometric mean of the two stop masses. Although Eq. (5) does not care about the sign of $A_t$, but given the slope of its RG trajectory a large negative value of $A_t$ is easier achieved than a positive value of the same magnitude.

Results and other implications: In Fig. 3 we show the scatter plot of the Higgs mass for different choices of the ratio of high scale parameters $A_0$ and $m_0$ in cMSSM+ (red points) and in cMSSM (blue points). We scan over the following ranges: $m_0 = [0, 2]$ TeV, $M_{1/2} = [0, 2]$ TeV and $A_0 = [0, -2]$ TeV, keeping $\tan \beta = 10$ and $\mu > 0$ (preferred by $(g-2)$ of muon). What is significant is that for the above parameter choices, especially, the magnitude of negative $A_0$ not exceeding 2 TeV, cMSSM struggles to give the Higgs a mass of 125 GeV [23][24], but for the same ranges of parameters cMSSM+ offers a 3 to 4 GeV enhancement to the Higgs mass which is enough to bring it into consistency with the CMS and ATLAS measurement. This is simply a consequence of a more negative value of $A_t$ that is attainable in cMSSM+ compared to what is possible in cMSSM, i.e. $|A_t (\text{cMSSM+})| > |A_t (\text{cMSSM})|$, starting from a given (negative) $A_0$ at the high scale. This happens because the running of $A_t$ has a steeper slope in cMSSM+ due to the tweaking of its RG evolution by the adjoint contribution. In Fig. 4 we choose the same high scale parameters, except that now $A_0 = [0, -4]$ TeV, so that cMSSM can accommodate a 125 GeV mass of the Higgs. The shaded regions in the plane of the gluino and the lighter stop masses correspond to points for which the light Higgs weighs between 123 and 127 GeV. What we demonstrate in Fig. 4 is that there is a significant recovery of the lighter spectrum in the cMSSM+ compared to cMSSM. We have made use of two packages, SuSpect [26] and micromegas [27], during the implementation of these plots, and we have ensured that the predictions for some low scale observables, e.g. $B_s$ decays to $\mu^+\mu^-$, are consistent with their experimental observations in the shaded regions.

We now give a quantitative estimate of how much we gain in terms of fine-tuning. Using the Barbieri-Giudice criterion [28], a rather crude estimate of the amount of fine-tuning as a function of the stop masses is given by

$$\Delta \approx \frac{10}{33} \cdot \frac{m_{\tilde{t}_1} m_{\tilde{t}_2}}{(650 \text{ GeV})^2} \cdot \ln \left( \frac{M_G}{M_Z} \right),$$  \hspace{1cm} (6)

where it is to be noted that $\ln(M_G/M_Z) \approx 33$. Using a fixed set of high scale input parameters, we present two sample points for stop masses at the weak scale in GeV: $(m_{\tilde{t}_1}, m_{\tilde{t}_2}) = (1186, 1950)$ for cMSSM and $(512, 902)$ for cMSSM+. Both these points correspond to $m_h = 123$ GeV. Using Eq. (6), we obtain $\Delta \approx 54$ for cMSSM and $\Delta \approx 11$ for cMSSM+. Thus we gain roughly a factor of 5 in terms of fine-tuning, though this rough estimate which we presented for illustration is only for a specific set of sample points.

We briefly mention some phenomenological implications of the color adjoint superfield. As shown in [5], its scalar component (assumed heavier) can decay into its fermionic component and gluino much before the time of nucleosynthesis
to avoid any cosmological problems. The fermionic state can decay into characteristic 4t+ neutralino missing energy through nonrenormalizable interactions, which has very low standard background. If the fermionic component is long-lived, there are ways to ensure that its relic density is insignificant. In the LHC context it has been shown that the experimental bounds on such models with Dirac gauginos are weaker than on similar MSSM-type models \[13\]. Several other authors have studied different hadronic decay signatures of color-octet scalars and pseudo-scalars \[30,32\]. The outcome of our analysis in the context of the Higgs mass may provide further motivation for the LHC studies of the adjoint states.

Conclusion and outlook: Na"ively, one might think that the presence of any colored matter would induce a similar shift in the Higgs mass. Indeed so, but not always in the preferred direction. We tried also with diquarks (\(\Phi\)) appearing through a superpotential \(W_{\text{DQ}} = y_{\Phi}^a u_i^c u_j^c \Phi + \mu_{\Phi} \bar{\Phi} \Phi\), where for illustration we considered coupling with up-type singlets. Here \(\Phi = (3, 1, 4/3)\) and \(\bar{\Phi} = (3, 1, -4/3)\) (see, e.g. Ref. \[33\] for a list of possibilities for different diquark representations). It is easy to check that \(\Delta b_3 = 1\) in this case, which is to be contrasted with \(\Delta b_3 = 3\) for the adjoint – see Eq. (1). This is one of the reasons behind the much smaller shift in the Higgs mass that one can get with a diquark. Indeed in the diquark case, \(\Delta h_2\) and \(\Delta b_1\) would be non-vanishing, but these are numerically not so relevant in this context. Crucially, the diquark Yukawa coupling \(y_{\Phi}\) contributes in the ‘wrong’ direction to the running of \(y_t\) and \(A_t\), and its magnitude has to be kept under control as otherwise \(y_t\) would blow up pretty fast. We relegate a more detailed study of different types of diquarks in this context to a future publication. This last observation of ours, i.e. Yukawa running in this context is indeed a tricky issue, is in accord with a recent study claiming that additional chiral fermions at the GUT scale with large Yukawa couplings modify \(A_t\) in a way that the light Higgs mass is actually reduced for the same stop and gluino masses \[34\].

Herein lies the reason as to why the color adjoint extension offers the most promising scenario in the present context.

We advance three distinct features which make our scenario a worthy competitor of the alternative avenues for providing a few extra GeV to the Higgs mass: (i) The existence of a chiral superfield in the adjoint representation is well motivated as arising from higher dimensional supersymmetric theory. (ii) The colored fermionic component may provide a large Dirac mass of the gluino, which offers many interesting features, including an improved naturalness. We have also shown using the Barbieri-Giudice criterion that fine-tuning in cMSSM+ is lessened by a factor \(~5\) with respect to cMSSM . (iii) Even if the adjoint scalar and fermion states are much heavier than 1 TeV, e.g. if we assume them to weigh around 10 TeV, our main conclusion remains unaffected. The central issue is what is the value of the strong coupling when it stops running (at one loop level). If the adjoint scalar and fermion masses are about 1 TeV, then \(\alpha_s^{-1}(1\ \text{TeV}) \simeq 9.7\) is where the strong coupling freezes and stays unmoved for higher energies up to the GUT scale. On the other hand, if the adjoint scalar and fermion masses are around 10 TeV, then the slope of the strong coupling keeps changing up to the energy scale.
of 10 TeV, and then as soon as the adjoint states are excited the coupling gets frozen at \( \alpha_\gamma^{-1}(10 \text{ TeV}) \simeq 10.8 \), which is much closer to its value at 1 TeV and quite far away from the GUT scale value (in cMSSM, i.e. without the adjoint state) \( \alpha_\gamma^{-1} \simeq 23.4 \). As a result, even for a 10 TeV adjoint state, a few GeV enhancement to the Higgs mass would not be a problem. Indeed, the collider phenomenology of a 10 TeV state is less exciting, which is why we assumed the adjoint states to weigh around 1 TeV. In this sense our scenario offers a rather robust mechanism for the incremental shift in the Higgs mass.

We reiterate that by no means one can say that cMSSM (or, equivalently, mSUGRA) is already disfavored. All we observe is that the light Higgs mass in cMSSM struggles to reach out to the last few rungs of its experimental range. This can be considerably eased if, instead of holding ourselves hostage to the conventional particle content of the MSSM, we add a new TeV scale color adjoint superfield, which has enough motivations to exist and which offers rich phenomenology to be explored at the LHC.

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